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Sviluppo di un apparato sperimentale per la realizzazione di un gas quantistico dipolare di atomi di Disprosio

Development of an experimental apparatus for the realization of dipolar quantum gases of Dysprosium atoms

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Chapter 1

Introduction

In this thesis I will describe the birth and the first steps of a new experiment for the study of a quantum degenerate dipolar gas of Dysprosium. In chapter 1 I will explain the state of the art of quantum gases experiments and how this experiment will bring a significant improvement. I will also introduce some relevant data about the atomic species we use and finally I will provide an overview of our experimental setup.

In chapter 2 I will discuss the theory of light-matter interaction to better understand the working principles of the cooling stages I designed and implemented: Zeeman slower, optical molasses and magneto optical trap (MOT). I will also introduce a few principles of non linear optics exploited in second harmonic generation (SHG). Indeed, we built a SHG cavity to generate laser light at the broadest atomic transition. I will also describe the locking techniques we use for the locking of the cavity and the locking of the laser to the atomic transition.

In chapter 3 I will describe the optical system, in particular the SHG cavity and the vacuum setup. Then, I will report the first measurements to characterize both the laser system and the atomic beam, also I will report the first signature of slowed atoms.

In the last part of the thesis I will speak about the next stages of the experiment, in particular the MOT, which has not yet been implemented.

1.1 Why dipolar quantum gases?

Most physical systems, notably solids, are governed by long-range interactions such as the Coulomb interaction between electrons and the dipolar interaction between magnetic spins. Besides the standard phases of matter, these interactions have been predicted to give rise to a wealth of exotic quantum phases, encompassing supersolids, quasicrystals, frustrated crystals and self-assembled structures. Observing these phenomena in standard condensed-matter systems is however very hard, because of the difficulties in eliminating unwanted effects (lattice phonons, disorder, etc.) and in controlling precisely the systems parameters. In the last 20 years quantum gases have been successfully used to study quantum phenomena due to the high degree of control that can be achieved with laser light and other electromagnetic fields. Quantum gases are allowing the realization of quantum simulators of paradigmatic problems, such as quantum phases or strongly correlated matter dynamics, that cannot be solved with classical computers or studied in standard condensed-matter systems.

This great opportunity has however been so far exploited mostly in the case of atomic species, such as the alkalies, interacting via van der Waals forces, which have substantially a contact nature. This prevents the possibility of studying long-range interaction effects. Some progress has been achieved using quantum degenerate gases of molecules, but the experimental challenges in the realization of stable samples of ultracold molecules have somehow limited the proceedings of this approach, even if this remains an open path.

The use of atoms with high magnetic dipole like Dysprosium or Erbium can be a breakthrough in this topic. These atoms, besides the van der Waals interaction, also interact via dipole-dipole interaction, which is both long-range and anisotropic as explained in the section 1.3. The combination of these two ingredients leads to the appearance of a number of peculiar quantum phenomena so far unexplored. Dramatically different phases are expected to appear depending on the geometry of the system and on the relative strength of kinetic energy, interaction energies and chemical potential. Notably, all these parameters can be controlled and tuned in the experiment.

1.2 Ultracold quantum gases

Quantum gases are very dilute systems if compared with condensed matter ones; nevertheless most of their properties are governed by interactions between particles. Indeed, if atoms are confined in tight trapping potentials to reduce the system dimensionality or in lattices, they can become stronglycorrelated particles whose behavior is dominated by interactions. In the recent years, strongly-correlated ultracold atoms have demonstrated to represent extraordinarily useful tools to study condensed matter-like physics in a clean and controllable environment and they have been employed as simulator of quantum Hamiltonians [1].

Quantum effects can be observed in a systems with density n when the spacing between particles, $n^{-1/3}$, is comparable or smaller than the de Broglie wavelength $\lambda_T = \sqrt{2\pi\hbar/mk_BT}$, where m is the mass of the particle and T is the temperature of the system [2]. At first sight all the efforts to reach ultracold temperatures may seem useless since employing higher densities would allow to observe quantum effects at higher temperatures. So

why all the ultracold atoms experiment use very low densities, of the order of 10^{-19}m^{-3} ? The reason can be found observing that almost all elements are solid at temperatures below the K. At those temperatures, measurements must then be performed before the gas "realizes" that it should be a solid. The lifetime of this metastable condition is given by the rate of three body loss, that is the process by which three atoms are converted into a bond dimer and a free atom that carries away the binding energy. Which is basically the fist step to form a solid from an atomic gas. The rate of this losses can be expressed as $R = Ln^2$ where L is a parameter that depends on the atomic species involved and is of the order of 10^{-40} m⁶/s, thus for a lifetime of the order of 1s we must have $n < 10^{-20} \,\mathrm{m}^{-3}$. On the other side, elastic scattering, which is responsible for the thermalization of the gas, bounds the density from below. Indeed, the scattering rate can be expressed as $\gamma = n\sigma v_T$ where $v_T = \sqrt{8k_BT/\pi m}$ and σ is the scattering section and is of the order of 10^{-15} m². Imposing the quantum regime $\lambda_T n = 1$ we find that $\gamma \propto n^{4/3}$ thus the densities for experiments with quantum gases are bound to $10^{-19} \sim 10^{-20} \text{m}^{-3}$.

[2] [3] [4] The interparticle spacing of a quantum gas at this densities is of the order of $2 \cdot 10^4 a_0$ where $a_0 = 5.3 \cdot 10^{-11}$ m is the Bohr radius. Thus the average separation is by far greater than the dimension of the electronic cloud surrounding the atomic nucleus. This makes interaction to be pairwise scattering events, since it is very unlikely to find three or more atoms in range of their own electronic clouds. Since atoms are neutral in charge the first non vanishing order interaction is the so called van der Waals interaction due to dipole-induced dipole interaction, whose potential is proportional to $1/R^6$, where R is the distance between the nucleus.

This potential is classified as a short range interaction [3]. In a system of particles interacting via short-range interactions, the energy is intensive in the thermodynamic limit. On the contrary, in systems with long-range interactions, the energy per particle does not depend only on the density, but also on the total number of particles. It is easy to see that a necessary condition for obtaining an extensive energy is that the integral of the interaction potential U(r) converges at large distances.

$$\int_{r_0}^{\infty} U(r) d^D r \neq \infty \quad \Rightarrow \quad U(r) \propto 1/r^n \text{ with } n > D$$
(1.1)

where r_0 is a short distance limitation to avoid divergence of the integral and D is the dimensionality of the system. From this point of view van der Waals interaction is short range in any dimensionality. It is also important to consider the angular momentum l = 0, 1, 2, ... in the scattering process; for identical bosons or fermions only even or odd values of l are allowed, respectively. The effective scattering potential contains a centrifugal barrier of the order of [5]:

$$E_{cent} \sim \frac{\hbar^2}{m_r a_S^2} l(l+1) \tag{1.2}$$

Where a_S is the van der Walls scattering length described below and m_r is the reduced mass of the colliding atoms. This barrier is usually of the order of mK for l = 1, so that for an ultracold gas, whose temperatures are of the order of some μ K, scattering is possible only for bosons in s-wave (l = 0) because in this case the barrier vanishes. Fermions instead can not collide because the s-wave channel is not allowed.

In consideration of the above, this interaction can be schematized as a contact interaction described by a single parameter a_S , called the scattering length:

$$U_{\rm eff} = \frac{4\pi\hbar^2 a_S}{m}\delta(r) \equiv g\delta(r) \tag{1.3}$$

Typical values of a_S for the alkali are of the order of $100a_0$. This potential, despite its simplicity, leads to very interesting phenomena. A prominent example is the superfluid-Mott insulator transition driven by the competition between repulsive contact interaction and kinetic energy in a lattice [6]. More complex Hamiltonians have also been investigated with such systems, revealing the interplay between different elements intrinsically present in real matter, although hardly controllable. This is the case for example of disordered interacting systems [7].

Another important feature of the contact potential in quantum gases, which is one of the great advantages in the use of quantum gases as quantum simulator, is that this potential is tunable. Indeed, taking advantage of magnetic Feshbach resonances [8] it is possible to change the value of a_S and also its sign by varying the value of an external magnetic field [5]:

$$a(B) = a_{bg} \left(1 - \frac{\Delta}{B - B_0} \right) \tag{1.4}$$

where Δ is the width of the Feshbach resonance and B_0 its position; a_{bg} is the background (non resonant) scattering length.

1.3 Dipolar interaction

In the previous section I spoke about ultracold atoms and van der Waals short range interactions and how useful they are to perform quantum simulations. However some neutral atoms have a permanent magnetic dipolar moment in their ground state, thus they interact also through a dipole-dipole potential:

$$U_{dd}(\mathbf{r}) = \frac{C_{dd}}{4\pi} \frac{(\mathbf{e}_1 \cdot \mathbf{e}_2)r^2 - 3(\mathbf{e}_1 \cdot \mathbf{r})(\mathbf{e}_2 \cdot \mathbf{r})}{r^5}, \qquad C_{dd} = \begin{cases} \mu_0 \mu^2 \\ d^2/\epsilon_0 \end{cases}$$
(1.5)



Figure 1.1: Two particles interacting via dipole-dipole potential. [3]

Where \mathbf{e}_1 and \mathbf{e}_1 are the versors indicating the dipoles orientations and \mathbf{r} is the distance between the two dipoles as drawn in figure 1.1. The expression for C_{dd} changes if magnetic or electric dipoles are involved: μ_0 and ϵ_0 are the vacuum permeability and permittivity and μ and d are the magnetic or electric dipole moment. Assuming that the sample is polarized, i.e. all the dipoles are oriented in the same directions the dipole-dipole can be expressed as:

$$U_{dd}(r) = \frac{C_{dd}}{4\pi} \frac{1 - 3\cos(\theta)}{r^3}$$
(1.6)

Where θ is the angle between the dipole orientations. This potential is strictly speaking long range only in 3D since it does not meet the criterion n > D, and thus the interaction depends not only on the density but also on the total number of particles. However, also in 2D and 1D the range of interaction is sensibly longer than the contact van der Waals potential and particles do not need to be superimposed in order to interact. In this sense we can consider the dipole force a long-ranged force in any dimensionality.

Another fundamental feature of this potential is that it is anisotropic and both attractive and repulsive interactions are possible depending on the relative orientation of the two dipoles. This leads to a series of new interesting phenomena, beyond those that can be simulated using the simple contact interaction. Dipolar quantum gases pave the way to the study of exotic quantum insulating "solid" phases, such as checkerboard phase or superfluid phases such as the supersolid phase [9].

Let us take a closer look at the dipole-dipole interaction. The potential has the angular symmetry of the Legendre polynomial of second order (dwave, $\Delta l = 2$) $P_2(\cos\theta) \propto 1 - 3\cos\theta$. As θ varies, the coefficient P_2 goes from 1 (repulsive) for dipoles sitting side by side to -2 (attractive) for dipoles in the head to tail configuration.

In order to realize a quantum gas with significant dipole-dipole interaction, one can use particles having either an electric dipole moment d, or a magnetic dipole moment μ . Usually the dipolar coupling is much higher in the electric case. Indeed, the typical order of magnitude of d for an atomic or molecular system is $d \sim ea_0$ where e is the charge of the electron, while $\mu \sim \mu_B$. The ration between the two potentials can be expressed as

$$\frac{\mu_0 \mu^2}{d^2 / \epsilon_0} \sim \alpha^2 \sim 10^{-4} \tag{1.7}$$



Figure 1.2: Two polarized dipoles side by side repel each other while in a 'head to tail' configuration attract each other. [3]

where α is the fine structure constant.

To discern if a system is governed by the dipolar potential it is important to compare the dipole-dipole interaction with the van der Waals interaction. For this purpose it is useful to define a scattering length also for the dipolar potential:

$$a_{\rm dd} \equiv \frac{C_{\rm dd}m}{12\pi\hbar^2} \tag{1.8}$$

The behavior of a dipolar system depends on the ratio $\epsilon_{\rm dd} = a_{\rm dd}/a_S$. If $\epsilon_{\rm dd} \ll 1$ the dipolar interaction is only a perturbation in the dynamics of the system, while for $\epsilon_{\rm dd} \gg 1$ it is possible to observe dipolar effects.

There are different candidates to realize a quantum degenerate gas of dipolar particles. In order to achieve the strongest interaction possible an electric dipole is the best candidate, however there are some technical issues that makes very difficult to use those systems.

Polar molecules are ideal candidates to study the dipole-dipole interaction due to their high electric dipole. Unfortunately, only heteronuclear molecules show permanent dipole moment. Also the molecule has to be in a low rotovibrational state because the dipole moment scales as R^{-7} and R increases with the rotovibrational quantum numbers. Finally an external field is required to polarize the sample, once the ground state has J = 0and the dipole moment averages to 0 in this case [10]. Molecules are expected to exhibit ϵ_{dd} of the order of 100 ~ 1000 however, up to the present, molecules have never been cooled directly down to quantum degeneracy. A working approach to degeneracy is to create heteronuclear molecules from ultracold atomic mixtures using Feshbach resonances. Feshbach association leads to the most excited bound state, molecules have the to be transferred to their ground state and this can be done using a STIRAP (STImulated Rapid Adiabatic Passage) process [11]. However this technique was to date successful only assembling KRb molecules which are not chemically stable (the state Rb_2+K_2 is energetically favourite against 2KRb), for this reason this molecules have to be assembled in optical lattices with one atom per specie in each lattice site. Na-K molecules have been recently proposed for

Species	Dipole moment	$a_{\rm dd}$
$^{87}\mathrm{Rb}$	$1\mu_B$	$0.7a_0$
$^{52}\mathrm{Cr}$	$6\mu_B$	$16a_0$
$^{170}\mathrm{Er}$	$7\mu_B$	$67a_{0}$
$^{164}\mathrm{Dy}$	$10\mu_B$	$133a_{0}$
KRb	0.6D	$2000a_{0}$

Table 1.1: Dipole moment and dipolar scattering length a_{dd} of some candidates for the study of dipole-dipole interaction [3].

this kind of experiments because they are chemically stable [12].

Extraordinarily large electric dipole moments can be obtained for highly excited Rydberg atoms. As the Kepler radius, and thus the dipole moment, scales with n^2 , where n is the main quantum number, the dipolar interaction energy can in principle scale like n^4 . However the lifetime of a Rydberg atom in a dense gas is very short to perform experiments on the dipolar interaction. Also the blockade effect prevents the excitation of other Rydberg atoms in a large sphere surrounding and existing Rydberg atom. Therefore it is possible to excite only few atoms in the typical dimension of a BEC thus it is difficult to study many body physic with those particles. However Rydberg atoms in a BEC are currently used to investigate collective behaviour in the excitation dynamics and the excitation blockade effect due to dipolar interactions [13].

So far we have seen that despite their interaction is stronger for electric dipoles, there are several technical issues in creating samples suitable to be studied. Another possible way to realize dipolar interacting samples is to use magnetic dipoles. Indeed, many atoms have a magnetic moment of the order of μ_B in their ground state, and it is much easier to prepare a quantum degenerate atomic gas if compared to molecules. Some atoms, like Chromium, Erbium, Dysprosium, and others, have a large magnetic moment of several Bohr magnetons in their ground state, and thus experience significant magnetic dipole-dipole interaction as reported in table 1.1.

Assuming a_S of the order of $100 a_0$ it is clear that the most magnetic atoms can be used to study the dipole-dipole interaction since $\epsilon_{\rm dd} \sim 1$. Furthermore it is, in principle, possible to use Feshbach resonances to reduce the scattering length down to the order of a_0 to make the dipolar interaction dominant [14].

In this experiment we intend to study the effects of the dipolar interactions in bosonic Dy in reduced dimensionality. In particular we would like to study the consequences of the interplay between the dipolar interaction and a periodic lattice in one and two dimensions. The Hamiltonian for such a system is the Bose -Hubbard extended Hamiltonian that takes into account

Isotope	Abundance	
160 Dy	2.34%	boson
$^{161}\mathrm{Dy}$	18.9%	fermion
$^{162}\mathrm{Dy}$	25.5%	boson
$^{163}\mathrm{Dy}$	24.9%	fermion
$^{164}\mathrm{Dy}$	28.2%	boson

Table 1.2: Natural (numeric) abundance of Dy isotopes[18]. Other isotopes are present in nature with a negligible abundance.

the dipolar interaction between atoms sitting in the next-neighboring lattice sites.

$$H = -J\sum_{\langle ij\rangle} \hat{a}_{i}^{\dagger} \hat{a}_{j} + \sum_{i} \left[\frac{U}{2} n_{i}(n_{i}-1) - \mu n_{i} \right] + \sum_{\mathbf{l}} \sum_{\langle ij\rangle_{\mathbf{l}}} \frac{U_{\mathbf{l}}}{2} n_{i} n_{j} \qquad (1.9)$$

where i, j counts the lattice sites, \mathbf{l} is the distance between near lattice sites. J is the tunneling energy, U is the on site interaction that depends both on the contact potential and the dipolar interaction, while the last part is the dipolar interaction where $U_{\mathbf{l}} = U_{dd}(\mathbf{l})$. All the parameters are controllable. The kinetic energy can be controlled varying the intensity of the optical lattice, contact interaction can be tuned via Feshbach resonances, which have been demonstrated to exist for Dy atoms at small magnetic fields [14], whereas dipolar interaction can be possibly varied by a time modulation of the orientation of the external magnetic field [15]. Finally, the chemical potential can be varied changing the atom number or the trapping confinement.

Another interesting line of research is the sudy of spinor dipolar condensates in which the orientation of the dipoles is not fixed by an external magnetic field but the dipoles are free together. This should led to selfassembled structures that have not yet been observed. This kind of experiments is technically very challenging because external field have to be reduced below μ G to avoid polarization in the sample. [16]

1.4 Dysprosium

Dysprosium is a rare-earth metal of the series of lantanides. Its atomic number is 66 and it is present in nature in various isotopes both fermionic and bosonic as listed in table 1.2. Bosons have nuclear spin I = 0 and fermions have I = 5/2. The fermionic isotopes therefore show a hyperfine structure of the fundamental level shown in Figure 1.4.

The melting point of Dy is $1412 \,^{\circ}$ C and it's vapor pressure is therefore very low until the temperature reaches values near the melting point. In



Figure 1.3: Dy energy level diagram adapted from [17].



Figure 1.4: Iperfine structure of the fermions ground state.



Figure 1.5: Dysprosium vapor pressure

figure 1.5 it is plot the vapor pressure in function of the temperature. This parameter is very important for the design of the experiment as we will discuss in the next section. To have a sufficient flow out of the oven we decided to work at temperatures of the order of 1200 °C.

The reason why we chose Dysprosium for this experiment is its magnetic moment: $10 \mu_B$. It is the most magnetic stable atom in its fundamental state. The high magnetic moment can be explained observing its electronic configuration: [Xe]4f¹⁰6s². The partially filled f orbital has a high quantum number ℓ , thus there are unpaired electrons in an high angular momentum state.

It's important to remember that the dipolar interaction scales like the squared value of the magnetic dipoles involved, as expressed in equation 1.6 so, if we compare Dysprosium with one of its neighbor in the periodic table, Erbium $(7 \mu_B)$ we can see that the interaction strength is twice as large in Dysprosium.

In figure 1.3 the principal transition of Dy are highlighted. For our experiment we are interested in two of them.

The 421 nm transition from the fundamental to level $[Xe]4f^{10}6s6p$. where an 6s electron jumps to level 6p, is the most intense and has a short lifetime of 4.94 ns: we use it for the first cooling stages. The 626 nm transition that links the fundamental with a level $[Xe]4f^95d6s^2$ where a 4f electron jumps to 5d level has instead a narrow linewidth of 136 kHz and it is suitable for further cooling to lower temperatures, thus we use this transition to realize the MOT. The main features of this lines are listed in table 1.3 and table 1.4.

symbol	description	definition	value
λ	wavelength (air)		$421.1714\mathrm{nm}$
k	wave vector	$2\pi/\lambda$	$23743{\rm cm}^{-1}$
Γ	spontaneous decay rate	1/ au	$2.02 \cdot 10^8 \mathrm{s}^{-1}$
au	life time	$1/\Gamma$	$4.94\mathrm{ns}$
$\Delta \nu_{\rm nat}$	natural linewidth	$\Gamma/2\pi$	$32.2\mathrm{MHz}$
I_{sat}	saturation intensity	$\pi hc\Gamma/3\lambda^3$	$56.4\mathrm{mW/cm^2}$
v_c	capture velocity	Γ/k	$14\mathrm{m/s}$
T_D	Doppler temperature	$\hbar\Gamma/2k_B$	$774\mu{ m k}$
a_{\max}	maximum acceleration	$\hbar k \Gamma/2m$	$5.79 \cdot 10^5 \mathrm{m/s^2}$
$T_{\rm rec}$	recoil temperature	$\hbar^2 k^2/m$	$660\mathrm{nK}$

Table 1.3: Main characteristics of the 421 nm transition

symbol	description	definition	value
λ	wavelength (air)		626.XXX nm
k	wave vector	$2\pi/\lambda$	$1.0 \cdot 10^5 {\rm cm}^{-1}$
Г	spontaneous decay rate	1/ au	$8.54 \cdot 10^5 { m s}^{-1}$
τ	life time	$1/\Gamma$	$1.17\mathrm{us}$
$\Delta \nu_{\rm nat}$	natural linewidth	$\Gamma/2\pi$	$136\mathrm{kHz}$
$I_{\rm sat}$	saturation intensity	$\pi hc\Gamma/3\lambda^3$	$72.5\mu\mathrm{W/cm^2}$
v_c	capture velocity	Γ/k	$85\mathrm{mm/s}$
T_D	Doppler temperature	$\hbar\Gamma/2k_B$	$3.26\mu{ m k}$
a_{\max}	maximum acceleration	$\hbar k\Gamma/2m$	$1.68 \cdot 10^3 \mathrm{m/s^2}$
$T_{\rm rec}$	recoil temperature	$\hbar^2 k^2/m$	$300\mathrm{nK}$

Table 1.4: Main characteristics of the 626 nm transition



Figure 1.6: Spettroscopy of the 421 nm transition of Dysprosium [19]. On each line are reported the isotope and the m_j of the corresponding iperfine transition



Figure 1.7: SolidWorks design of the final configuration of the vacuum apparatus

1.5 Experimental setup outline

In a quantum gas experiment a hot atomic vapor has to be cooled down to quantum degeneracy following a series of steps that use different cooling techniques.

In the specific case of Dysprosium one of the most challenging technical issues is the high melting temperature and the consequently low vapor pressure at room temperature. To reach a suitable vapor pressure Dysprosium must be heated up to about $1200 \,^{\circ}$ C. To reach such temperatures an oven, namely an effusive cell, is typically used. It consists of an electrically heated crucible with a hole in the front to let the high temperature atomic gas flow out. The drawback of heating the atomic source is the high thermal speed of the atomic beam effused by the cell. Such velocity, of the order of $400 \, \text{m/s}$, prevents the possibility of directly cooling the atoms in a Magneto Optical Trap (MOT) as usually done in standard alkali atoms experiments. A MOT is in fact characterized by a capture velocity of the order of tens of m/s, as we will see in section 3.3.1. A preliminary slowing step is therefore necessary.

To slow the atomic beam we use a Zeeman slower. It is composed by a series of electric coils winded around a tube. It uses the Zeeman effect combined to the light scattering force to slow the atoms while they travel through it. In front of the Zeeman slower we set up a transverse cooling stage to collimate the beam in order to maximize the flux that arrives at the end of the slower. After the Zeeman slower, atoms will have a speed low enough to be captured in the MOT.

To effectively slow the atoms we use the 421 nm transition because its short lifetime allows reaching a high scattering rate, thus a fast cooling process. This broad transition is however not suitable to reach very low temperatures because of its high doppler temperature, so we decided to use the 626 nm transition to realize the MOT.

We intend to study different dipolar phenomena with this setup, and we therefore aim to build a flexible and versatile system. In particular we are interested in the study of reduced dimensionality phenomena. A method to create 1D and 2D samples is to use optical lattices to separate the sample in a set of non interacting tubes (1D) or pancakes (2D) but this requires to illuminate the sample with laser beams in addition to the ones used for cooling. To guarantee a wide optical access we plan to use an additional glass cell attached on a side of the MOT cell. This kind of cell allow to use investigation methods such as high resolution imaging, that requires to position a microscopy objective very close to the sample, or bragg spectroscopy, that requires to send light to the sample at high angles. Another important feature of this solution, especially for a highly magnetic atom like Dysposium, is that glass is a non magnetic material. This allow to precisely control the magnetic field on the sample without residual magnetization from the vacuum setup.

Before transferring the atoms from the MOT cell to the glass cell we will catch them from the MOT in an Optical Dipole Trap (ODT) and perform a first evaporative cooling stage. This trap is based on the dipole force that a laser beam applies to atoms attracting them to the high intensity region as explained in section 2.1.2. By shifting the focus of an ODT beam it is possible to move the atoms in the glass cell where further evaporation will take place to bring the atoms to quantum degeneracy.

Chapter 2

Theoretical background

2.1 Light-matter interation

A detailed analysis of the mechanisms of interaction between electromagnetic fields and atoms is out of the aim of this thesis. However, in this section, I will review the main findings regarding the effects of a laser beam on the mechanics of atoms. The interaction between light and matter is at the basis of the cooling and trapping techniques used in quantum gases experiments.

Laser light is commonly used to investigate energy differences between atomic levels, however, other than internal energy an atom exchanges also momentum with the electromagnetic field, hence it experiences a force affecting the dynamics of its center-of-mass. This interaction can be described, in a quantum framework, as an exchange of momentum $\hbar k$ due to absorption and emission of photons. In the classical description, the radiation induces a dipole moment on the atom that interacts with the radiation itself. From light-matter interaction two forces arise: scattering force and dipole force. Let us now analyze both of them in detail. [20]

2.1.1 Scattering force

Momentum conservation implies that when an atom absorbs or emits a photon its momentum changes by $-\hbar \vec{k}$ where k is the wave vector of the photon. When a atom is illumined by a resonant laser source, every absorbed photon pushes the atom in the propagation direction. Of course after a certain time, of the order of the transition lifetime τ , the atom returns to its ground state emitting a photon of the same momentum. However, since the spontaneous emission direction is isotropic, the resulting momentum change due to the dis-excitation averages to 0 over an high number of scattered photons.

The resulting force can be calculated as the average momentum transferred times the scattering rate R_{scatt} that, for a laser of intensity I, frequency ω , detuned from the atomic transition ω_0 by $\delta = \omega - \omega_0$, can be calculated as

$$R_{\text{scatt}} = \frac{\Gamma}{2} \frac{I/I_{\text{sat}}}{1 + I/I_{\text{sat}} + 4\delta^2/\Gamma^2}$$
(2.1)

and therefore

$$\vec{F} = \hbar \vec{k} \frac{\Gamma}{2} \frac{I/I_{\text{sat}}}{1 + I/I_{\text{sat}} + 4\delta^2/\Gamma^2}$$
(2.2)

 $I_{\text{sat}} = \pi h c \Gamma / 3 \lambda^3$ is the saturation intensity of the transition of interest and $\Gamma = 1/\tau$ is the transition linewidth.

As we can see the scattering force grows as the detuning decreases and it is directly proportional to the radiation intensity. It is important to notice that the force is also proportional to k so it will be higher for shorter wavelength.

For resonant light, $\delta = 0$, and high light intensity, $I \gg I_{sat}$, we find the maximum scattering force

$$F_{\max} = \hbar k \frac{\Gamma}{2} \tag{2.3}$$

This can be still be interpreted as $F = \Delta p / \Delta t$ where the atom receives a momentum $\hbar k$ every two mean lifetime of the atomic transition i.e. the time needed to absorb a photon and then return to the fundamental state¹.

It is important to note that this force decreases quadratically with the detuning and it is correlated to the imaginary part of the refractive index of the atomic cloud which is responsible for absorption. The real part of the refractive index is instead related to the dipole force.

2.1.2 Dipole force

So far we described the scattering force, which is classically related to the absorption of light by a medium. However, there exists another kind of radiation force not related to absorption, arising from the deflection of light by a dispersive medium. In fact the deviation also implies a variation in the photons momentum direction that is transferred to the medium.

The interaction of a two level atom and a laser field in this case is called dipole force and it is due to a gradient of the light intensity. It can be demonstrated that this force takes the following expression:

$$\vec{F}_{\rm dip} = -\frac{1}{2I_{\rm sat}} \vec{\nabla} I \frac{\hbar \delta}{1 + (2\delta/\Gamma)^2 + (I/I_{\rm sat})}$$
(2.4)

Far from the resonance $\delta \gg \Gamma$ this equation can be approximated as

$$\vec{F}_{\rm dip} \simeq -\hbar \frac{\Gamma^2}{8\delta} \vec{\nabla} \left(\frac{I}{I_{\rm sat}} \right)$$
 (2.5)

 $^{^{1}}$ We are considering a closed transition so the decay always leads to the fundamental state. For the so called open transitions it is necessary to pump the atoms back to the fundamental level

It is important to note that the direction of the force depends on the sign of the detuning: it is attractive to high intensity zones if the laser is red detuned. In any real atom there are a lot of transitions, which can be both red or blue detuned with respect to a given laser frequency. In that case we have to sum over all the transition weighting each one with their own detuning and linewidth in order to calculate the dipole force. This average is usually summarized in the parameter $\alpha(\omega)$. In this case the dipole force can be expressed as

$$\vec{F}_{\rm dip} \simeq -\frac{1}{2} \alpha(\omega) \vec{\nabla} \left(\frac{I}{I_{\rm sat}}\right)$$
 (2.6)

To conclude the dissertation about the manipulation of atoms using laser light let us compare the two mechanism described above. It is important to see that dipole force is vanishing for zero detuning where the scattering force is maximum. While for high detuning $\delta \gg \Gamma$ dipole force usually dominates because of its detuning dependence

$$F_{
m scatt} \propto rac{1}{\delta^2} \qquad \qquad F_{
m dipole} \propto rac{1}{\delta}$$

Of course the condition to have a dipole force is to have an intensity gradient. This is the reason why large, collimated beams near the resonance are used for exploiting the scattering force, while far detuned, highly focused beams are used for dipole force applications. Also this force can be neglected for large collimated beams where the intensity varies slowly.

2.2 Line broadening

In atomic physics every transition is characterized not only by its energy but also by its linewidth[20]. In fact only the ground state can be considered to have a precisely defined energy, because it is a stable state, while all the excited levels have a defined lifetime. In fact time and energy are complementary variables from a quantum point of view, and the uncertainty principle states that it is impossible to know exactly the energy of an excited state since $\Delta E \geq \hbar/\Delta t$. Assuming the life time τ as the uncertainty on the time and using $E = \hbar \omega = \hbar 2\pi \nu$ is possible to define the natural linewidth as

$$\Delta \nu_{\rm nat} \equiv \frac{1}{2\pi\tau} \equiv \frac{\Gamma}{2\pi} \tag{2.7}$$

This linewidth is a direct consequence of quantum mechanics and can not be sharpened with any experimental technique. The intensity of spontaneous emitted radiation has a Lorentzian profile of FWHM $\Gamma/2\pi$ in terms of frequency (Γ in terms of pulsation)

$$I(\delta) = I_0 \frac{(\Gamma/2)^2}{\delta^2 + (\Gamma/2)^2}$$
(2.8)

This linewidth is difficult to observe because there are a series of phenomena that contribute to widen the measured linewidth. Out of many line broadening effects, in the following paragraphs I introduce Doppler broadening and power broadening. These are the two main effects that give contribution to the linewidth in the measurements of this thesis.

2.2.1 Doppler broadening

Doppler broadening is a consequence of atomic motion and Doppler effect, which is a direct consequence of relativity. An atom in motion in the laboratory frame emits and absorbs, for any given transition, light with different frequency with respect to the one measured when the particle is at rest. This frequency is

$$\omega' = \omega_0 + \vec{k} \cdot \vec{v} \tag{2.9}$$

where \vec{k} is the wavevector of the photon and \vec{v} is the velocity of the atom emitting or absorbing the light, measured in the laboratory frame.

If the atom is moving in the same direction of the photon, the product $\vec{k} \cdot \vec{v}$ is positive thus ω' will be shifted towards higher frequencies. The exact opposite happens when the photon and the atom are moving in opposite directions.

Let us now consider a laser light illuminating an atom. If the atom moves in the propagation direction of the laser, the light needs to be blue shifted with respect to the transition frequency in order to be in resonance with the atom. On the contrary if the atom moves in the opposite direction, the laser light has to be red shifted.

In most practical cases we deal with a gas of atoms, with mass m, at low densities so that it can be treat as an ideal gas. The speed distribution along one direction at a given temperature T is described by the Maxwell Boltzmann distribution

$$f(v_x) = \frac{1}{u\sqrt{\pi}}e^{-\frac{v_x^2}{u^2}}$$
(2.10)

where $u = \sqrt{2k_BT/m}$ is the most probable speed. Integrating this over all directions we find

$$f(v) = \frac{1}{u\sqrt{\pi}} 4\pi v^2 e^{-\frac{v^2}{u^2}}$$
(2.11)

The velocity distribution of a 3D gas is therefore a Gaussian multiplied by v^2 . The transitions frequencies for atoms in every velocity class are shifted according to the Doppler effect. To obtain the frequency distribution along the propagation direction of a laser light we just replace the expression for v obtained from equation 2.9 in 2.10. The result is again a gaussian distibution centered on ω_0

$$g_D(\omega) = \frac{c}{u\omega_0\sqrt{\pi}} e^{-\left[\frac{c^2}{u^2} \left(\frac{\omega-\omega_0}{\omega_0}\right)^2\right]}$$
(2.12)

The doppler broadening is defined as the full width half maximum of the distrubution

$$\frac{\Delta\omega_D}{\omega_0} = 2\sqrt{\ln(2)}\frac{u}{c} = 2\sqrt{\frac{2\ln(2)k_BT}{m}}$$
(2.13)

The doppler broadening has a gaussian shape with the width the scales as \sqrt{T} . Th effect of the gas velocity distribution is therefore to broaden the linewidth of the transition. The resulting line has a Voigt distributions, given by the convolution between the natural Lorentzian profile with the Gaussian Doppler broadening.

2.2.2 Power broadening

Power broadening is an effect that arises when an atomic sample is illuminated by a light with intensity comparable or above the saturation intensity of the interested transition. The line results to be broadened because the saturation reduces the absorption near the resonance. Far from resonance instead, the transition is not saturated because the absorption probability is smaller, so the absorption signal does not change significantly compared to the one for low intensities. This makes the absorption peak flattened at the center and unchanged on the tails. The line shape is still Lorentzian but the FWHM result increased. The expression for the power broadening can be derived [20] and takes the following form:

$$\Delta\omega_{\rm FWHM} = \Gamma \left(1 + \frac{I}{I_{\rm sat}} \right). \tag{2.14}$$

2.3 Optical cooling and trapping techniques

Let us now consider how the force produced by light matter interaction described in section 2.1 can be used to manipulate an atomic gas. I am going to introduce the theory behind all the techniques used in my thesis work or designed but not yet implemented in the experiment. First I will describe the optical molasses: they are an efficient cooling method that can lower the temperature of an atomic gas of several orders of magnitude. Their functioning can be explained using only the scattering force and the Doppler effect described in equation 2.9. I will then speak about the Zeeman slower, which is an effective way to slow an atomic beam from very high speed, it uses a combination of scattering force, Doppler effect and Zeeman effect. The last method I will describe takes advantage of the scattering force is the Magneto Optical Trap (MOT), which is a sort of completion of the optical molasses that allows not only cooling but also trapping of the atoms. Finally I will explain how the dipole force can be used to trap atoms in a non dissipative potential.

2.3.1 Optical Molasses

Optical molasses are a basic application of the scattering force to slow down atoms. As suggested by its name, the effect of this technique is to apply a force to the atoms proportional to their velocity i.e. a viscous force.

$$\vec{F} = -\alpha \vec{v} \tag{2.15}$$

This force reduces the velocity of atoms since it is always directed against the direction of motion so, from a thermal point of view, it cools the atomic gas. This technique has no confinement mechanism because there is no position dependence in the applied force. An atom can therefore escape the molasses region while travelling at low speed.

Let us derive how to produce a force like the one proposed in equation 2.15 in one direction. The extension in two and three directions is then easy. We consider an atom and a laser beam moving along the same direction. Due to the Doppler effect the transition is shifted according to the atom velocity. The total detuning of the laser from the transition can be expressed as $\delta = \delta_0 - \vec{k}\vec{v} = \delta_0 \mp kv$ where $\delta_0 = \omega - \omega_0$ is the detuning in the rest frame. The sign -(+) is considered if the atom is moving with the same (opposite) orientation of the laser beam. If we substitute this expression for δ in equation 2.2 we obtain

$$\vec{F} = \hbar \vec{k} \frac{\Gamma}{2} \frac{I/I_{sat}}{1 + I/I_{sat} + 4(\delta_0 \mp kv)^2 / \Gamma^2}$$
(2.16)

An optical molasses is basically composed by two counter propagating laser beams red detuned from the atomic transition. At a first glance, one might think that this symmetrical arrangement has no effect on an atom since there are equal and opposite forces on it. However, the radiation forces from the two laser beams balance each other only for a stationary atom, which is what we want to achieve. For a moving atom the Doppler effect works in the opposite way for the two beams and this leads to an imbalance in the forces. Let us consider what happens in the reference frame of an atom moving towards the right. In this frame the Doppler effect leads to an increase in the frequency of the laser beam propagating towards the left. This Doppler shift brings the light closer to resonance with the atom and therefore increases the rate of absorption from this beam. The laser beam propagating to the right results instead even more detuned than in the rest frame, and so the scattering rate from this laser is further reduced. This leads to a resultant force that slows the atom down:

$$\vec{F}_{\text{molass}} = \hbar \vec{k} \frac{\Gamma}{2} \frac{I}{I_{\text{sat}}} \left[\frac{1}{1 + \frac{I}{I_{\text{sat}}} + \frac{4(\delta_0 + kv)^2}{\Gamma^2}} - \frac{1}{1 + \frac{I}{I_{\text{sat}}} + \frac{4(\delta_0 - kv)^2}{\Gamma^2}} \right]$$
(2.17)

Equation 2.17 can be quite cumbersome so let us first take a simplified approach. To understand how the molasses work we can consider the particu-

lar case of low laser intensities $I \ll I_{\text{sat}}$ and small atom velocities $kv \ll \Gamma$. Under this approximation it is possible to write equation 2.16 as

$$F \simeq F_0 \mp \alpha v \tag{2.18}$$

where F_0 is the component independent on velocity

$$F_0 = \hbar k \frac{\Gamma}{2} \frac{I/I_{\text{sat}}}{1 + (2\delta_0/\Gamma)^2}$$
(2.19)

and α is

$$\alpha = 4 \frac{\hbar k^2}{\Gamma} \frac{I}{I_{\text{sat}}} \frac{-\delta_0}{\left[1 + (2\delta_0/\Gamma)^2\right]^2}$$
(2.20)

The force resulting from the two counter propagating laser beams is

$$F_{\text{molass}} = (F_0 - \alpha v) - (F_0 + \alpha v) = -2\alpha v \qquad (2.21)$$

which is a viscous force, like we expected, with viscosity $k = 2\alpha$. The viscosity depends on the detuning of the two beams and it is positive for a red detuned light. Also the viscosity is proportional to intensity but it is important to remember that this expression is valid only in the low intensities approximation. Let us now consider the full expression for the optical molasses force 2.17 plotted in figure 2.1. The first thing we note is that the force remains proportional to velocity only for small velocities. This can be explained considering that the Doppler effect for a very fast atom shift the transition frequency far from resonance for both laser beams, thus the resulting total force drops back to zero over a certain range of velocities. We can also observe that in the first plot that for small detuning the maximum force is quite low and the range of linearity is very small because the two beams are both close to the resonance for low speed. On the other side a bigger detuning makes the two beams interacting with totally different velocity classes, both at high speed. In this situation an atom at low speed is not resonating with any laser beam so the viscosity for low speed decreases. Finally, as we can note observing the plot on the bottom, increasing the intensity is an effective way to increase the viscous force only until it is close to the saturation intensity. Further increments of the intensity however allow to use a higher detuning of the laser beams because power broadening of the transition occurs. This can be useful to achieve a constant viscosity over a wider range of velocities.

To realize 2D or 3D molasses it is enough to use 2 or 3 pairs of counterpropagating laser beams orthogonal to each other.

To conclude this part let us consider which is the minimum temperature we can expect to reach using this method. So far we have considered the average result of an high number of photon absorption and emission, however fluctuations have been neglected. The spontaneous emission causes the atom



Figure 2.1: Force of an optical molasses for the 421 nm transition ($\Gamma = 2\pi 32 \text{ MHz}$) in function of the velocity. Above: the force is plotted as a function on v for $I = 10I_{sat}$ and for different light detuning: $\Gamma/2$ in orange, Γ in red, 2Γ in green and 4Γ in blue. Below: the force is plotted for a fixed detuning $\delta = \Gamma$ and for different intensities: $I = 0.3I_{sat}$ in orange, $I = I_{sat}$ in red, $I = 3I_{sat}$ in green, $I = 30I_{sat}$ in blue.

to recoil of $v_{\rm rec} = \hbar k/m$ in random directions, and these recoil kicks lead to a random walk in the space of velocities that produces a rms velocity of

$$\langle v^2 \rangle = \langle \cos^2 \phi \rangle v_{\rm rec}^2 N(t) = (1/3) v_{\rm rec}^2 R_{\rm scatt} t \qquad (2.22)$$

where N(t) is the number of photons scattered in time t; and $\langle \cos^2 \phi \rangle$ comes from the randomness of scattering directions identified with the angle ϕ . If we include also the effect from the uncertainty in the number of absorbed photons and we consider 3 pairs of beams, for a 3D molasses we find [20, ch. 9.3.1]

$$\langle v^2 \rangle = 4 v_{\rm rec}^2 R_{\rm scatt} t$$

This is the sum of all the heating contributions while the viscous cooling is acting. We can write the variation of kinetic energy as the work done by the viscous force.

$$\frac{d}{dt}\left(\frac{1}{2}m\langle v^2\rangle\right) = -\alpha\langle v^2\rangle + \frac{1}{2}m4v_{\rm rec}^2R \tag{2.23}$$

Setting the time derivative to 0 we get the stationary situation. Using the approximated expression for α of equation 2.20, and converting $\langle v^2 \rangle$ in a temperature $(T = m \langle v^2 \rangle / 2k_B)$ we find:

$$T = \frac{\hbar\Gamma}{4k_B} \frac{1 + (2\delta_0/\Gamma)^2}{-2\delta_0/\Gamma}$$
(2.24)

The minimum temperature is reached for $\delta_0 = -\Gamma/2$ ad it is defined as the Doppler cooling limit:

$$T_D = \frac{\hbar \Gamma}{2k_B} \tag{2.25}$$

It gives the lowest temperature reachable with the optical molasses technique. As we can see it depends only on the linewidth of the used transition. These temperatures for the Dysprosium transitions we use are recalled in tables 1.3 and 1.4.

2.3.2 Zeeman Slower

As pointed out in the previous section, optical molasses work only for low velocities. The plots of equation 2.17 in figure 2.1 clearly show how the molasses are effective for velocities of the order of $v_c = \Gamma/k$ that in the case of the 421 nm transition is approximately 14 m/s. As the velocity exceeds this value the effect of the molasses vanishes because the Doppler effect brings the light completely out of resonance with the atomic transition. In our experiment we have to slow an atomic beam with a typical velocity of $v_{\rm mp} = \sqrt{3k_BT/m} \simeq 490 \,\mathrm{m/s}$ corresponding to a Doppler shift of $\Delta \nu_{\rm dopp} \sim 1 \,\mathrm{GHz}$. Optical molasses would therefore have no effect on the atomic beam.

One could think about shifting the laser frequency to compensate for the Doppler effect, anyway this approach would result to be useless for two reasons. First we want to slow a wide distribution of speed, and so a laser with a fixed frequency would interact with just a class of speed, leaving the other part of the distribution imperturbated. Second, the light would be out of resonance as soon as the velocity has decreased by $\Delta v = \Gamma/k$.

A standard method to overcome this problem in slowing atomic beams is a so called Zeeman slower which, as suggested by the name, effectively exploits the Zeeman effect. The basic idea is to start cooling the atoms with a given velocity and, as their speed decreases, to use the Zeeman effect to compensate exactly the decreasing of the Doppler shift. In this way it is possible to keep the atoms in resonance with the laser light until they are almost stopped.

Let us now enter more in detail in the Zeeman slower mechanism. The energy shift caused by a magnetic field B is

$$\Delta E = \mu_B B (g_f m_{J_f} - g_i m_{J_i}) \tag{2.26}$$

where the subscript *i* and *f* identify the initial and final atomic states; $g_{i,f}$ is the landé factor; $m_{J_{i,f}}$ is the projection of the total angular momentum along the magnetic field direction. For the sake of simplicity in the following we consider the simple example of a transition $(J = 0 \rightarrow J = 1)$ and positive $g = g_i = g_f$, the generalization to more complex cases is direct. We can write the total atom-light detuning as

$$\delta = \delta_0 + kv \mp g \frac{\mu_B B m_J}{\hbar} \tag{2.27}$$

Where now we take only the positive sign for the Doppler shift because the atoms are always moving towards the laser. The sign for the Zeeman shift depends on the light polarization. The upper state is split in 3 Zeeman levels $m_J = 0, \pm 1$. The transition resonant with a σ^+ polarized beam increases its energy as the field increase, and so the light will result red detuned. in this case the sign minus has to be considered in equation 2.27. The sign plus has instead to be considered for a σ^- polarized beam.

There are two ways to build a Zeeman slower. We could set the laser frequency to be resonant with slow atoms ($\delta_0 \sim 0$) and use a magnetic field to compensate the doppler shift for faster atoms. This configuration is called σ^+ because this is the polarization that makes the Zeeman shift opposite to the Doppler shift. The other way is to set the laser frequency to be resonant with fast atoms ($\delta_0 \sim -kv_i$) without any magnetic field and then keep them in resonance while their speed decreases using an increasing magnetic field. This is called σ^- configuration because, in this case, we want the Zeeman shift to have the same sign of the Doppler shift.

The latter configuration employs light far detuned from the atomic transition, thus it is not resonant with the slow atoms out of the slower, where there is no magnetic field. This is very useful in a setup were atoms are trapped in a MOT at the end of the slower because the slower light do not affect them.

It is also possible to realize a spin flip Zeeman slower schematized in figure 2.2 that consists in a first part in the σ^+ configuration and a second part in the σ^- configuration, i.e. with the magnetic field oriented in the opposite direction. For the reasons described in section 3.2.5 we decided to realize this kind of Zeeman slower.



Figure 2.2: Simple energy scheme of a spin flip Zeeman slower working on a transition $J = 0 \rightarrow J = 1$. The magnetic field is 0 on the y axis and increases in opposite directions leaving the axis. The laser beam comes from the right. In figure it is reported only the Zeeman shift and not the shift due to the Doppler effect. (The Zeeman shift and the detuning is not to scale with the atomic transition). Adapted from [20]

An important parameter in the design of the Zeeman slower is its length. In fact it exists a maximum force that can be impressed to the atoms, as expressed in equation 2.3, and therefore a minimum length for the Zeeman slower. Solving the equations of motion we obtain the minimum length necessary to slow atoms with initial velocity v_i

$$L_{\min} = \frac{v_i^2}{2} \frac{m}{F_{\max}} = v_i^2 \frac{1}{\Gamma} \frac{m}{\hbar k}$$
(2.28)

Usually this length is of the order of less than 50 cm, as an example, for Dy starting at 450 m/s the minimum length is 14 cm. However it is practically impossible to guarantee the maximum deceleration over the whole slower, and so a Zeeman slower is usually designed to achieve an acceleration $a = \eta a_{\text{max}}$ with $\eta < 1$, and thus $L = L_{\text{min}}/\eta$.

It is important to remember that also Zeeman slower application uses the scattering force and therefore it has a collateral heating effect as painted out in the case of the molasses. In this case the random direction of spontaneous emission is not critical for the final velocity that can be achieved but it produces an increasing of the mean square transverse velocity that can be interpreted as a divergence of the atomic beam. To get a quantitative idea, let us estimate the number of scattered photons between t = 0 and t as

$$N(t) = \frac{v_i - v(t)}{v_{\text{rec}}} \tag{2.29}$$

where $v_{\rm rec} = \hbar k/m$ is the recoil velocity. The root mean square values of transverse velocity components is

$$\left\langle v_{x,y}^2(t) \right\rangle = \frac{v_{\text{rec}}^2}{3} N(t) \tag{2.30}$$

In our case, assuming the initial velocity of 450 m/s and the final velocity null, the number of photons scattered by the atom before stopping is $N \simeq 8 \cdot 10^4$. The corresponding transverse RMS velocity is therefore around 1 m/s. This velocity is negligible while the atom longitudinal velocity is high but becomes important at the end of the Zeeman slower. In fact, considering a final velocity of 10 m/s, the atomic beam divergence θ is

$$\theta = \arctan\left(\frac{v_{x,y}}{v_z}\right) = 6^{\circ}$$
(2.31)

2.3.3 Magneto Optical Trap (MOT)

Optical molasses and Zeeman slowers are efficient techniques to cool an atomic gas, but they have no confining effects. In fact atoms emerging from a Zeeman slower are no more in resonance with the slowing light, so they are substantially free. Similarly atoms accumulate in the optical molasses region because it takes a long time for them to escape the molasses region thanks to their random walk or collisions between themselves, but their density is still low. To add a confinement effect we have to make the scattering force dependent on the position. This can be done by choosing the proper polarization of beams and adding a magnetic field gradient to an optical molasses configuration. This technique takes the name of Magneto Optical Trap (MOT).

Magnetic field gradients can be produced by two identical coaxial coils with opposite current direction (anti-Helmotz configuration). In this way a quadrupole field is created where the field is zero in the center of the coils and its magnitude increases linearly in any direction for small displacements from the central position.

The magnetic fields does not confine the atoms by itself, but produces a variable Zeeman shift of the atomic hyperfine levels. This position depending

shift causes an unbalance in the scattering forces of the counterpropagating beams, thus a confinement force.

Let us take for example a cooling transition between a J = 0 ground state and a J = 1 excited state as schematized in figure 2.3. The field gradient produces a splitting in the Zeeman levels of the excited state (J = $1 \Rightarrow m_J = 0, \pm 1$) that increases as the atoms leave the center of the MOT.



Figure 2.3: Energy scheme of a 1D MOT working on a closed transition between $J = 0 \rightarrow J = 1$. The magnetic field increases, with opposite directions leaving the center of the mot z_0 . The laser beams are red detuned with respect to the transition (The Zeeman shift and the detuning is not to scale with the atomic transition). Adapted from [20]

To produce a restoring force the red detuned light has to be circularly polarized with the right elicity in order to be resonant with the hyperfine level lowered in energy by the Zeeman effect. In this way a still atom displaced from the center of the trap is more resonating with the laser that pushes it back to the center. Near the center of the trap and far away from the coils the field gradient is constant and the Zeeman effect is linear in the displacement, thus the detuning between laser and atomic transition can be expressed, similarly to equation 2.16, as

$$\delta = \omega - \omega_0 \mp kv \mp \beta z$$

$$= \delta_0 \mp kv \mp \beta z$$
(2.32)

with

$$\beta = \frac{g\mu_B}{\hbar} \frac{\partial B}{\partial z}$$

We can replace δ in equation 2.18 to obtain whit analougus approximations:

$$F = F_{\text{scatt}}^{\sigma^+}(\delta_0 - kv - \beta z) - F_{\text{scatt}}^{\sigma^-}(\delta_0 + kv + \beta z)$$
$$\simeq -2\frac{\partial F}{\partial \delta_0}(kv - \beta z) = -\alpha v - \frac{\alpha \beta}{k}z$$
(2.33)

An atom that enters the MOT region will be slowed by the molasses and the cold atoms will be pushed to the center of the MOT by the position dependent force. The lowest temperature reachable in a MOT is in principle the same as in the molasses. However the higher densities reached in the MOT allow a heating mechanism neglected since this point: collisions between atoms. For this reason a MOT is usually hotter than a molasses.

What is the fastest atom that can be captured and cooled by a MOT? The MOT capture velocity is a critical parameter in the design of a cold atom experiment. Let us take the best condition we can: a molasses of diameter D always applying the maximum force $F_{\text{max}} = \hbar k \Gamma/2$ to the atoms. To effectively stop the atoms before they pass through the molasses their velocity have to be below

$$v_c \simeq \sqrt{\frac{\hbar k \Gamma D}{m}} \tag{2.34}$$

This is clearly an overestimation of the real value of v_c . For the 626 nm transition and a beam diameter of $4 \text{ cm} v_c = 11.6 \text{ m/s}$ while for the 421 nm transition $v_c = 217 \text{ m/s}$.

2.3.4 Optical dipole trap

Besides MOT there is another commonly used trapping mechanism: the dipole trap. It uses the dipole force instead of the scattering force. Looking at the approximate expression for the dipole force in equation 2.5 we can express this force as generated by a potential $\vec{F}_{\rm dip} = -\vec{\nabla}U_{\rm dip}$, thus:

$$U_{\rm dip} = \frac{\hbar \Gamma^2}{8\delta} \frac{I}{I_{\rm sat}}$$
(2.35)

For a red detuned laser beam ($\delta < 0$) the potential is attractive toward the direction of high intensity.

A focused laser beam has the following Gaussian intensity profile:

$$I(r,z) = I_0 \frac{1}{1 + (z/z_R)^2} \exp\left(\frac{-2r^2}{w^2(z)}\right)$$
(2.36)

where the origin of z axis if fixed where the beam waist w is minimum and equal w_0 to $z_R = \pi w_0^2 / \lambda$; $w(z) = w_0 \sqrt{1 + (z/z_R)^2}$. The Gaussian beam profile has increasing intensity towards the focal point in both radial direction ad axial direction. To realize a dipole trap is therefore enough to focus a red detuned laser beam. Let us note that the dipole force is conservative, thus a dipole trap does not cool atoms like the optical molasses technique, but simply keep them in the center of the trap if their kinetic energy is low enough not to escape the potential.

This characteristic may look like a handicap of this method, instead it can be used to perform evaporative cooling. The basis of this cooling technique is not to slow every single atom of the sample, but to let hotter atoms leave the sample in order to decrease the average temperature. There are various ways to apply this concept but one of the simplest is to trap a cloud of atoms in a dipole trap and then slowly lower the laser intensity.

In a dipole trap the heating due to the spontaneous emission that was the limit to the MOT temperature is not present. The laser light used to realize dipole traps needs to be far detuned from any atomic transition in order to reduce scattering events. The scattering force in fact depends on the detuning as $1/\delta^2$ while the dipole force has a dependence on $1/\delta$. It is important to notice that even if a single scattering event occurs, since the recoil temperature is far beyond the trap temperature, the atom leaves the trap reducing the number of atoms trapped but not increasing their temperature.

Dipole trap are in principle very simple to realize. The only need one high power laser beam far detuned from any transition but not necessarily monochromatic or locked to a frequency reference. With reasonable light power, traps with depths of the order of few μ K can be realized. Optical trap can also be used to move a cold atomic cloud within a vacuum setup: the movement can be realized simply by moving the focal point of the laser beam.

2.4 Second Harmonic Generation

As already told, in the experiment we exploit the 421 nm atomic transition in the Zeeman slower and the transverse cooling stages. To obtain light at this frequency we double the frequency of a 842 nm commercial laser using a non linear crystal in a ring cavity. The details of how we implement this technique in the experiment will be explained in section 3.1 Here I am going to introduce some basic concepts of non linear optics and in particular of second harmonic generation (SHG).

Non linear optics is the branch of optics that describes all the phenomena in which the response of a material is not linear in the light field. The polarization \vec{P} must then be described using not only a linear term in the electric field \vec{E} but also terms of higher orders:

$$\vec{P}(\vec{x},t) = \epsilon_0 \left[\chi^{(1)} \vec{E}(\vec{x},t) + \chi^{(2)} \vec{E}^2(\vec{x},t) + \dots \right]$$

where χ^n is the susceptivity at the *n*-th order and ϵ_0 is the vacuum per-

mittivity. Second harmonic generation is one of the practical application of nonlinear optics. In a classical picture a nonlinear crystal is made of a great number of dipoles, generating a global polarization as stated above. Expanding to the second order, it is trivial to see that a monochromatic field, $E(t) = A \cos(\omega t)$, will produce a component of the polarization oscillating at twice the original frequency:

$$P(t) = \epsilon_0 \left[\chi^{(1)} E(t) + \chi^{(2)} E^2(t) + \dots \right] =$$

= $\epsilon_0 \left[\chi^{(1)} A \cos(\omega t) + \chi^{(2)} (A \cos(\omega t))^2 + \dots \right] =$
= $\epsilon_0 \left[\chi^{(1)} A \cos(\omega t) + \chi^{(2)} \frac{A}{2} (1 + \cos(2\omega t)) + \dots \right]$

In an extended crystal every dipole oscillates following the local phase of the incident field. If the dipoles interfere constructively they generate a second harmonic light that propagates in the crystal. For a nonlinear isotropic and non magnetic material, assuming no free charges or currents, the wave equation is [21]

$$\nabla \times \nabla \times \vec{\mathbf{E}} + \frac{1}{\epsilon_0 c^2} \frac{\partial^2}{\partial t^2} \vec{\mathbf{D}} = 0$$
 (2.37)

Where ϵ_0 is the vacuum permittivity and \vec{D} is the electric displacement field define as

$$\vec{\mathbf{D}} = \epsilon_0 \vec{\mathbf{E}} + \vec{\mathbf{P}}$$

In most practical cases the the amplitude of the involved fields is slowly varying over the wavelength, thus equation 2.37 can be simplified in

$$\nabla^2 \vec{\mathbf{E}} - \frac{1}{\epsilon_0 c^2} \frac{\partial^2}{\partial t^2} \vec{\mathbf{D}} = 0$$
 (2.38)

We can split \vec{P} and \vec{D} in a component linear with the electric field and a non linear component

$$\vec{\mathbf{P}} = \vec{\mathbf{P}}^L + \vec{\mathbf{P}}^{NL}$$
$$\vec{\mathbf{D}} = \vec{\mathbf{D}}^L + \vec{\mathbf{P}}^{NL}$$

where the linear component of $\vec{\mathbf{D}}$ is defined as

$$\vec{\mathbf{D}}^L \equiv \epsilon_0 \vec{\mathbf{E}} + \vec{\mathbf{P}}^L \equiv \epsilon_0 \epsilon^L \vec{\mathbf{E}}$$

In this way is possible to write equation 2.38 as

$$\nabla^2 \vec{\mathbf{E}} - \frac{\epsilon^L}{c^2} \frac{\partial^2}{\partial t^2} \vec{\mathbf{E}} = \frac{1}{\epsilon_0 c^2} \frac{\partial^2}{\partial t^2} \vec{\mathbf{P}}^{NL}$$
(2.39)

Let us simplify the electic field dropping the polarization and assuming that it is composed by monochromatic plane waves propagating along z, the total electric field can therefore be expressed as

$$E(z,t) = \sum_{n} E_n(z,t) = \sum_{n} A_n(z)e^{-i(kz-\omega t)} + c.c.$$
(2.40)

The same can be done for \vec{D} and \vec{P} . Equation 2.38 now becomes:

$$\frac{\partial^2}{\partial z^2} E_n(z) + \frac{\epsilon^L(\omega_n)}{c^2} E_n(z) = -\frac{1}{\epsilon_0 c^2} P_n^{NL}(z)$$
(2.41)

Let us now evaluate equation 2.41 for the case of second harmonic generation. There are only two frequencies involved ω_1 and $\omega_2 = 2\omega_1$, the relative electric fields can be expressed as

$$E_{1,2} = A_{1,2}e^{-i(k_{1,2}z - \omega_{1,2}t)} + c.c.$$

The non linear polarizations can be written as:

$$P_1(z) = 4\epsilon_0 d_{\text{eff}} E_1(z) E_2^*(z)$$
(2.42)

$$P_2(z) = 2\epsilon_0 d_{\text{eff}} E_1^2(z) \tag{2.43}$$

where d_{eff} is the effective nonlinear optical coefficient[21, p. 1.5]. Since $\chi^{(2)}$ depends on the orientation of the crystal, the polarization of light whit respect to the crystal and on the frequency involved in the non linear process, it is more convenient to tabulate the values of $d_{\text{eff}} = \chi_{\text{eff}}/2$. The first term in equation 2.41 can be neglected since the variation of amplitude in a wavelength is very small. In a lossless medium, we obtain:

$$\begin{cases} \frac{dA_1}{dz} = \frac{2id_{\text{eff}}\omega_1^2}{k_1c^2} A_1^* A_2 e^{-i\Delta kz} \\ \frac{dA_2}{dz} = \frac{id_{\text{eff}}\omega_2^2}{k_2c^2} A_1^2 e^{i\Delta kz} \end{cases}$$
(2.44)

If we assume A_1 constant and $\Delta k = 2k_1 - k_2 = 0$, which is the ideal case, it can be seen that the intensity of the second harmonic light $I_j = 2n_j\epsilon_0 c |A_j|^2$ grows quadratically with z as expected for a second order phenomena. If the pump field and the second harmonic field are out of phase i.e $\Delta k \neq 0$, we can integrate equation 2.44 over the direction of propagation to obtain the amplitude of the second harmonic field in function of the crystal length and the phase mismatch. It is possible to show that, in this case, the power rapidly oscillates between the pump beam and the second harmonic beam along the crystal. It is therefore strictly necessary to have a good control on Δk in order to have an efficient doubling conversion. We will treat how to control Δk in the next section. For practical purposes A_1 can not be taken as constant over the crystal because the aim is to convert as much radiation as possible and therefore A_1 decreases significantly along the crystal. The solution of equation 2.44 for $\Delta k = 0$ is

$$\begin{cases}
A_1 = \operatorname{sech}(z/l) \\
A_2 = \tanh(z/l)
\end{cases} (2.45)$$

where l is the length scale of the conversion process defined as

$$l = \frac{c\sqrt{n_1 n_2}}{2\omega_1 d_{\text{eff}} |A_1(z=0)|}$$
(2.46)

Equation 2.45 is plotted in figure 2.4 where we can note that the conversion is almost completed after a couple of characteristic length. In the general case of $\Delta k \neq 0$ there is a oscillatory behavior in z. As plotted in figure 2.5.

In this case it's crucial to choose a crystal long enough to convert as much light as possible but not too long to avoid the reconversion back to the pump frequency.



Figure 2.4: Spatial variation of second harmonic intensity (blu) and pump intensity (red) in units of characteristic distance plotted for $\delta k = 0$.

2.4.1 Phase matching

We have seen in the previous section how critical it is to satisfy the phase matching condition in order to have an efficient and stable conversion of the pump light into the second harmonic beam. In this paragraph we see in detail how we can get the phase match. Generally speaking in second harmonic generation we require

$$\Delta k = 0 \implies 2k_1 = k_2 \implies \frac{2n_1\omega_1}{c} = \frac{n_2\omega_2}{c} \implies n(\omega_1) = n(\omega_2) \quad (2.47)$$



Figure 2.5: Spatial variation of second harmonic intensity (normalized) in units of characteristic length l plotted for different values of $\Delta s = \Delta kl$ [21, fig 2.7.4].

This condition cannot be achieved in a normal dispersion material since the refraction index increases with the frequency. The typical approach to solve this problem is to take advantage of the birefringence, i.e. the characteristic of a materiel to show different refractive indexes in the different oscillation directions of the electric field.

For example, in a positive uniaxial crystal along one direction, called extraordinary, the refractive index n^e is higher than along the other two directions, called ordinary.

Equation 2.47 can be satisfied using a pump beam polarized along the extraordinary direction.

$$n^e(\omega_1) = n^o(\omega_2)$$

In this way second harmonic light $\omega_2 = 2\omega_1$ is generated along the ordinary axes which is the one that matches the refractive index. Certainly the two refractive indexes wont be exactly of the right values that satisfy the phase matching conditions. To solve this problem the crystal can be cut with a angle θ between the extraordinary axis and the direction of propagation so that the phase matching can be reached

$$n^{\prime e}(\omega_1) = n^o(\omega_2)$$
$$n^{\prime e}(\omega_1) = n^e(\omega_1)\cos\theta + n^o(\omega_1)\sin\theta$$

All this consideration are done for an uniaxial material but, to be complete, there are also biaxial materials. They exhibits three different n in the three directions. All the dissertation done since this point is still valid. The only difference in this case is that we can play with two angles in the orientation of the crystal with respect to the direction of propragation. In particular
we can choose one angle to be 0 and reach the phase match using the other one.

It is also possible to reach the phase matching changing the temperature, in fact the lattice spacing of the birefringent crystal changes with the temperature and thus also the refractive index changes. In some crystals the ordinary and extraordinary refractive indexes follows different functions of the temperature allowing to reach the phase matching with the temperature. Anyhow the temperature of the crystal needs to be efficiently stabilized in order not to affect the conversion efficiency.

In order to have a realistic estimation of the expected conversion efficiency we have to reconsider some of the approximation done since this point. First of all there will always be a little wavevector mismatch that limits the maximum conversion efficiency and makes the SH power fluctuating over the length of the crystal. Furthermore we should consider that it is convenient to use a focused beam in order to achieve the highest intensity possible for the second order process. The drawback of a narrow focusing is that we have a small confocal parameter $b = 2z_R$, so the SHG will take place only in a small length of the crystal. Also an highly focused beam has a small beam radius, thus there will be various \vec{k} in the incident beam and not all of them will be phase-matched with the crystal. It is now clear that there will be an optimal length for the crystal and an optimal focusing of the beam in order to achieve the maximum conversion efficiency. The calculations that lead to the optimal parameters are rather complex [21, 22, 23]; here I will simply report the main results.

Is possible to show that the best length L is given by

$$\xi \equiv \frac{L}{b} = \frac{\lambda L}{2\pi n w_0^2} = 2.84 \tag{2.48}$$

The expected efficiency is

$$\eta_{\rm SHG} = \frac{P_{2\omega}}{P_{\omega}} = h \frac{2\omega_1^3 d_{eff}^2}{\pi c^4 n^2 \epsilon_0} \mathcal{P}_{\omega} L \tag{2.49}$$

Where $P_{2\omega}$ and P_{ω} are the power of second harmonic light and pump light respectively, whereas h depends on the optical parameters of both the crystal and the beam in a non trivial way. For the optimized parameters and in the reasonable approximation of no absorption by the crystal and focal point placed in the center of the crystal h is of the order of the unit. The expected efficiency for available crystals and for the typically available pump power is far below the 1 %. For this reason SHG is usually performed in a cavity so that the effective light intensity on the crystal is enhanced of a typical factor of 100.

2.5 Locking techniques

In all the dissertation above we have assumed to be able to select the laser frequency at wish. The reality is quite more complicated: any laser source shows slow frequency drifts but also fast phase noise. As a consequence of this behavior there is the necessity to stabilize the frequency, this procedure is called locking. A laser can be locked to various frequency references. One of the most common is the spectroscopy signal of an atomic transition but it is also possible to lock a laser on the resonance of an optical cavity or to another laser. In this section I analyze the two locking techniques used in this experiment. The first one is called Hansch-Couillaud and it will be used to lock the length of the SHG cavity to the pump laser frequency. The second one, which also takes its name from the inventors Pound-Drever-Hall , is used in our experiment to lock the laser frequency to the atomic transition.

2.5.1 Hansch-Couillaud locking

Before entering in the details of the cavity lock, let me first say few words on the effects that can modify the length of a cavity. The cavity length must be a integer multiple of the wavelength in order to be resonating with the injected light: $L = n \cdot \frac{c}{\nu}$. In fact the length must be stabilized within a fraction of the wavelength. The higher is the cavity finesse, tha smaller is the allowed cavity length variation:

$$\delta L \ll \frac{\lambda}{F}$$

One can understand this considering that every round trip the wave sums on itself with a small phase mismatch. If the cavity has low finesse, waves will do few round trip before being lost due to mirror losses or absorption thus even a big phase mismatch will not propagate enough to reach the opposite phase and cancel the resulting field inside the cavity. On the contrary in a cavity with high finesse the light will do lots of round trips allowing small phase mismatches to sum over and cancel the electric field in the cavity.

This length is in the order of the nm thus is very difficult to build a cavity with thermal expansion lower than a nm. There are also sound and mechanical vibrations to keep into account and also air flow that can change locally the refraction index and the distance between elements. Also the laser frequency may vary and in order to keep the cavity resonating this variation must be less than

$$\delta \nu \ll \frac{FSR}{F} \simeq 1 \,\mathrm{MHz}$$

As we will see later all this sources can be reduced but the most practical way to keep the cavity length correct is to provide an active feedback to one of the mirrors, which can be mounted on a piezo driver. The cavity length can be adjusted applying a voltage to the piezo.

To provide the feedback signal we use a Hansch-Couillaud locking scheme [24]. This method is based on the fact that the cavity has a polarizing element within. if Θ is the angle between the direction of transmission of the polarizing element in the cavity and the polarization of the input beam, it can be expressed as:

$$E^{i} = E^{i}_{\perp} + E^{i}_{\parallel}$$
$$E^{i}_{\parallel} = E^{i} \cdot \cos(\Theta) \qquad E^{i}_{\perp} = E^{i} \cdot \sin(\Theta)$$

the reflected field is

$$E_{\parallel}^{r} = E_{\parallel}^{i} \left[\sqrt{R_{1}} - \frac{T_{1}}{\sqrt{R_{1}}} \frac{Re^{i\delta}}{1 - Re^{i\delta}} \right]$$

$$E_{\perp}^{r} = E_{\perp}^{i} \sqrt{R_{1}}$$

$$(2.50)$$

Where δ is the phase difference after a round trip, R_1 and T_1 are the reflectivity e trasmissivity of the input coupler and R < 1 represents the losses after a round trip of the cavity. As we can see at resonance $\delta = 2\pi n$ the two components remain in phase i.e. linearly polarized. If the cavity is slightly out of resonance the reflected light will be elliptically polarized. To detect this ellipticity we put a $\lambda/4$ retarder and a polarizing beam splitter. The intensities of the two beams I_a and I_b are analyzed with a differential photodiode. The intensity difference can be approximated as:

$$I_a - I_b = 2I^{(0)} cos(\Theta) sin(\Theta) \frac{T_1 R sin(\delta)}{(1 - R)^2 4 R sin^2(\frac{\delta}{2})}$$
(2.51)

The first thing we notice from equation 2.51 is that injecting the cavity parallel or perpendicular to the axis of the polarization sensitive element in the cavity will produce no error signal. In our case the LBO crystal is not a proper polarizing element but it shows a different phase shift for ordinary and extraordinary polarization and also has losses, due to conversion, only along the extraordinary axis. This complicate the equations but doesn't change the trend of equation 2.51.

In figure 2.6 is shown a simulation of the error signal. A PI lock fed by this signal has a capture range of the whole free spectral range in fact the signal is nonzero in a large part of the free spectral range and has the opposite slope of the the locking point. This is very useful because the system will lock itself without the necessity of searching the line and will re-lock if, for some reason, it jump out of locking point.



Figure 2.6: Calculated error signal for $R = 0.95 T1 = 0.98 \Theta = 0.1$. The divergence in the plot are due to the approximated equation used.

2.5.2 Pound-Drever-Hall locking

The Hansch-Couillaud technique is very powerful but it can to be used when there are no polarizing elements. In this cases it is possible to use the Pound-Drever-Hall technique [25]. The light of the laser is passed through device that introduces a modulation on the phase of the electric field of the light that passes through it.

$$E(t) = Ae^{i\omega t} \to E(t) = Ae^{i(\omega t + \beta \sin \Omega t)}$$
(2.52)

where $\Omega/2\pi$ is the frequency of the modulation provided by a local oscillator and β is a dimensionless parameter that characterize the amplitude of the modulation. An arbitrary modulation gives all the side bands with intensities described by the Bessel functions but for this technique only the first two sideband are usefull so I am going to consider only small modulation amplitude $\beta \ll 2\pi$. Under this approximation it is possible to write a first order taylor expansion for the exponential in order to separate the frequencies in the electric field above obtaining:

$$E(t) \simeq A e^{i\omega t} \left(1 + i\beta \sin(\Omega t)\right)$$

$$= A \left(e^{i\omega t} + \frac{\beta}{2} e^{i(\omega + \Omega)t} - \frac{\beta}{2} e^{i(\omega - \Omega)t}\right)$$
(2.53)

After the modulator the beam is injected to the frequency reference used for the stabilization. It can be a Fabry-Perot cavity or a saturated spettroscopy line. In any case if we analyze the output of the frequency reference it can be expressed as $E_{\text{out}}(\omega) = F(\omega)E_{\text{in}}(\omega)$ where F is a complex number. so if we analyze the power output we find:

$$P_{\text{out}} \simeq P_0 |F(\omega)|^2 + P_0 \frac{\beta^2}{4} \left\{ |F(\omega + \Omega)|^2 + |F(\omega - \Omega)|^2 \right\}$$

$$+ P_0 \beta \left\{ \Re[\chi(\omega)] \cos(\Omega t) + \Im[\chi(\omega)] \sin(\Omega t)] \right\}$$
(2.54)

where $P_0 \propto A^* A$ and

$$\chi = F(\omega)F^*(\omega + \Omega) - F^*(\omega)F(\omega - \Omega)$$
(2.55)

The output is read using a photodiode and it is sent to a mixer together with the local oscillator signal shifted in phase. A mixer is an electronic device which output is the product of the input signals.

$$V_{\rm mix} \propto P_{\rm out} \cos(\Omega t + \phi)$$
 (2.56)

The output of the mixer is then sent to a low-pass filter set at frequency much lower that Ω . In this way the only part of the signal that pass the filter is proportional to that oscillating at frequency Ω and is is the error signal for the locking system:

$$V_{\rm fb}(\omega) \propto \Re[\chi(\omega)] \cos(\phi) + \Im[\chi(\omega)] \sin(\phi)$$
(2.57)

With an appropriate choice of ϕ it is possible to maximize the error signal.

Looking at the feedback signal in figure 3.14 it is clear that a PI system can be locked on the central slope. Chapter 3

Project assembling and measurements



Figure 3.1: A picture of the optical table. On the bottom of the image there is the Ti:S laser with its green pump, while in the center it is possible to see the SHG cavity in the vacuum box (open).



Figure 3.2: A picture of the complete vacuum apparatus. The atoms are moving from right to left.



Figure 3.3: SolidWorks project of the vacuum apparatus and its support. The letter indicates the following parts: a)Effusive oven; b)Transverse cooling cell; c)differential pumping stage; d)Gate valve to isolate the two parts of the setup; e)Zeeman slower; f)6-ways cross for test measurements; g)Optical breadboard; h)Vacuum pumps; i)Valves to connect turbomolecular pumps.

In this chapter I will speak about the assembling of the apparatus and of the first measurements we performed to characterize the system. After testing the proper functioning of the effusive cell and of the Zeeman slower the apparatus has been moved to the CNR research area in Pisa in a equipped laboratory. The experiment is built on two separated optical benches. On the first one there are the laser sources as shown in figure 3.1. Light is passed from one table to the other through optical fibers. On the other table there are the vacuum setup and the optics necessary to shape the light beams used in the experiment as shown in figure 3.2. In figure 3.3 I reported the SolidWorks representation of the vacuum setup configuration we employed for the first characterization of the experiment reported in this chapter.

3.1 421nm light generation

In this experiment we use 421nm light for different purposes. The choice of the laser source is done considering the expected optical power needed to have efficient cooling. For the Zeeman slower we expect to have a proper slowing with less than 300 mW (see section 3.2.5), and the transverse cooling should work with 200 mW 3.2.4. A smaller part of the blue light is used for imaging purposes and for the locking of the light on the atomic reference. To make a conservative estimation we need 500 mW of blue light on the atoms.

However we have to consider that we use optical fibers to bring the light from the light sources table to the vacuum setup table. Also we need to use a set of independent frequency shifted beams for the various purposes. I'm going to explain how this can be done using AOMs in section 3.1.4. Unfortunately both fibers and AOMs have an efficiency far from 100% so we need at least 1 W of blue light.

To obtain such a high UV power we decided to double the frequency of a 842 nm Ti:S laser as described in section 2.4 using a LBO non linear crystal in a home made bow-tie ring cavity.

3.1.1 The TI:S laser

To generate the 842nm radiation necessary to pump the SGH cavity we chose a commercial Ti:Sapphire laser: SolsTiS by M2. This laser uses a Sapphire (Al_2O_3) crystal doped with Ti ions in a ring cavity with about 700 MHZ free spectral range. Ti:S shows an absorption peak from 400 nm to 600 nm and can emit radiation in the range between 700 and 1000 nm. In principle every mode resonant with the cavity and under the gain curve of Ti:s can oscillate in the cavity. A tunable birefringent filter provide a rough selection of the desired wavelength while an etalon in the cavity is used for fine frequency adjustments and to lock to the selected mode by sending a feedback to one or more cavity mirrors mounted on piezo drives. An external

Name	Sensitivity	Cut ν
RESN FAST	$3.4 \mathrm{MHz/V}$	$30\mathrm{kHz}$
RESN SLOW	$2.5\mathrm{GHz/V}$	$50\mathrm{Hz}$
REF CAV	$3\mathrm{GHz/V}$	$300\mathrm{Hz}$

Table 3.1: Feedback inputs for the SolsTi:S laser and relative characteristics

high finesse reference cavity with a free spectral range of 3 GHz allows to narrow the laser line down to 50 kHz. The laser frequency can be controlled via the ethernet user interface or applying feedback signal directly to the laser. The direct feedback can be applied to different parts of the tuning system as explained in table 3.1. The resonator input can be used when the laser is not locked to the reference cavity while, when the laser is locked, only the reference cavity input is available.

To pump the Ti:S laser we use a green 532 nm 10 W laser in order to obtain a power output of 2.4 W^1 . In figure 3.4 we can see the whole laser system.



Figure 3.4: A picture of the SolsTi:S laser by M2. The 532 nm, 10 W pump laser is on the right, the first metallic block contains only the injection optics, the second box with the M2 logo is the Ti:S laser and the last block on the left is the external reference cavity.

3.1.2 Cavity enhanced second harmonic generation (SHG)

Second Harmonic Generation has already been introduced in section 2.4; let us now focus on the implementation of this technique beginning from the choice of the non linear crystal.

There are various parameters to be considered. First of all the crystal has to be transparent for both the pump and the second harmonic frequencies, also the refractive indexes have to allow the phase matching. Finally,

 $^{^{1}}$ The power dropped to $1.9 \,\mathrm{W}$ during the first year of operations.

for our application, we require high damage threshold and a wide angular acceptance in order to obtain high SH power and to realize a setup easy to align and very robust.

LBO is a suitable choice because it is a good mix between high non linear coefficient, d_{eff} , and high power damage threshold. LBO is the short name for lithium triborate (LiB₃O₅) and it has successfully been used for similar applications. The refractive indexes of LBO are reported in table 3.2.

	n_x	n_y	n_z
$842\mathrm{nm}$	1.5692	1.5946	1.6099
$421\mathrm{nm}$	1.5877	1.6177	1.6321

Table 3.2: refraction index along the three directions of an LBO crystal for the wavelengths of interest

SHG is a process of the second order in the light field amplitude. A high light intensity is therefore required in order to enhance the conversion efficiency. We thus decide to realize an external optical cavity to amplify the optical power on the non linear crystal. Linear cavity is the simplest scheme possible but it can't be used in this case because it supports a standing wave, which has a $\cos^2(z)$ intensity pattern that excites only a part of the crystal. This would reduce the conversion efficiency but also may lead to spatial hole burning in the crystal. For this reason we decided to realize a bow tie ring cavity using the scheme drawn in picture 3.5.



Figure 3.5: Optical scheme of the SHG cavity.

It is made by four mirrors. The plane input coupler has high reflectivity at 842 nm to optimize the conversion efficiency as better explained below. The piezo mirror is highly reflective and it is 5 mm diameter in order to reduce its mass and allow a faster response to the locking system. The two curved mirrors are highly reflective for 842 nm and highly transmissive for 421 nm so that the red light is resonating in cavity while the blue light exits the cavity once generated.

The LBO crystal is fixed on a copper basement mounted on a precision positioning stage that can be both rotated and shifted to reach the phase match using the orientation method. The piezo basement is also thermally coupled to a Peltier cell. The thermal stabilization is a critical component of the SHG process because the LBO crystal is very sensitive to temperature. In fact, as stated in section 2.4.1, it can be also phase matched using temperature however, for the wavelength of interest, we should heat the crystal up to 150 °C to get the phase match. Even though we are using an angle matching scheme, it is crucial to keep the crystal at the temperature for which it has been designed (20° C). Active thermal stabilization is very important to stabilize thermal fluctuations that are critical at the moment of the cavity lock when, in a fraction of a second, the power in cavity pass from 0 to hundreds of W. The crystal starts to heat due to absorption and its length and refractive index changes. This makes the effective cavity length change very fast and the piezo actuated mirror have to move a lot to compensate this effect. A robust thermal stabilization of the crystal reduces this effect and it allows to lock in any point in the piezo stroke in order to perform long frequency scans without loosing the lock. For this reason the basement contain a fast NTC placed very close to the crystal to immediately detect its thermal fluctuations. The measured temperature provides the feedback to a PID controller that drives the Peltier cell. The optimization of the feedback parameter is of crucial importance to reduce the thermal drift that happens when the cavity is locked.

We enclosed the bow tie cavity in an aluminum box in which it is possible to realize a medium vacuum to prevent any interference on the SHG system. The box has two glass windows to let the light in and out. The walls where the viewports are glued are tilted in order to avoid a Fabry-Perot effect between the two glass-to-air interfaces of the window itself. The input window is 2" large to let the beam reflected from the input coupler exit the box and it has a 842 nm Anti Reflection (AR) coating, while the output windows has a 1" diameter and AR coating at 421 nm. The electrical wiring needed to control the piezo, the Peltier cell and to measure the temperature are linked to the outside with a feed-through connector. A picture of the open box with the SHG cavity inside is reported in figure 3.6 Inside the box, four sorbothane chips isolate the aluminum block upon which the cavity is build from the vibration of the optical table. The vacuum inside the box stops the acoustic noise but also prevents air flow, chemical deterioration of the crystal and changes in air refractive index due to heating.

To lock the length of the cavity to the wavelength of the Ti:S laser we decided to use the Hansch-Couillaud locking scheme described in section 2.5.1as schematized in figure 3.7. The mirrors arrangement of a bow-tie cavity provides an easy access to the beam reflected from the input coupler, which is spatially separated from the input one because the light is injected in the cavity with an angle respect to the mirror. The reflected



Figure 3.6: A picture of the SHG cavity inside the vacuum box (open). The 842 nm laser enters from the top of the picture and the 421 nm light exits from the bottom. The LBO crystal is mounted in the copper basement in the center of the cavity where it is possible to see also the wiring for the peltier cell and the NTC temperature sensor.

beam is attenuated and split in two parts: the first goes through a $\lambda/2$ and a $\lambda/4$ retarder before being split by a Polarizing Beam Splitter (PBS). The waveplates are useful to maximize the error signal and to remove any offset component in a handy and reliable way. The intensities of the beams with opposite polarization I_a and I_b are analyzed with a differential photodiode. The detector is a two stages differential amplifiers. The first stage is a low impedance current to voltage converter that works as an analog adder for the currents from the two photodiodes taken with opposite sign. It is possible to add a third input current using a trimmer to balance the offset between the photodiodes. The second stage is an inverting amplifier for which is possible to change the gain with another trimmer. Both stages have a low pass filter in order to remove noise in the frequency range that can't be reached by the feedback system, at the same time the filtering avoids possible oscillations due to high frequency poles in the system. The differential photodiode signal provides the feedback to a PID (figure 3.8). The output of the PID is amplified by a high voltage driver and controls the piezo mirror in the cavity. The amplification stage allows to exploit the whole piezoelectric traslator range (-30 V + 150 V) in order to span about 11 FSR of the duplication cavity. The lock is therefore very robust to thermal fluctuations of the room.



Figure 3.7: Optical scheme of the Hansch-Couillaud locking system for the SHG cavity. $\lambda/2$ waveplates are painted in green, $\lambda/4$ in red.



Figure 3.8: An image of the Hansch-Couillaud locking signal obtained scanning the cavity length with the piezo mirror and acquiring the signals with an oscilloscope. In yellow the signal of the differential photodiode, in green the signal of the coupling efficiency photodiode that shows a coupling of about 80%.

The other part of the beam is sent directly on a photodiode without any polarization analysis to optimize and measure the cavity coupling efficiency defined as $\eta_{\text{coupling}} = 1 - P_{\text{lock}}/P_{\text{unlock}}$ where P_{lock} is the reflected power when the cavity is resonant and P_{unlock} when it is far from resonance.

Let us now consider the design of the optical crystal and of the enhancement cavity. Equation 2.49 suggests that a long crystal enhances the conversion efficiency, however the longer the crystal is, the wider will be the beam at the entrance of the crystal itself. In fact, in order to optimize the conversion efficiency it is better to focus the beam in order to increase the optical intensity. However the smaller the waist, the smaller the Rayleigh range, thus the beam rapidly diverges. In order not to cut the tails of the beam would therefore be necessary to choose a wider crystal but this leads to problems in the thermal dissipation and increases a lot the cost of the crystal. Also, a long crystal amplifies the effects of a wavevector mismatch on the efficiency, making the phase matching difficult. We choose a 15 mm long crystal that is a good trade between the two extremes. To obtain the ideal focusing (eq.2.48) the waist inside the crystal has to be:

$$w_0 = \sqrt{\frac{\lambda_1 L}{2\pi n} \frac{1}{2.84}} = 26.6\,\mu\mathrm{m} \tag{3.1}$$

Considering a pump power of 2 W, equation 2.49 predicts an efficiency for single passage conversion of 0.14 %. Looking at this small value it appears clear why an external cavity is indispensable in order to achieve a reasonable conversion efficiency.

There are two critical parameters in the design of the cavity: the curvature of the curved mirrors to obtain the right focusing of the beam on the crystal and the reflectivity of the input coupler. We decided to use two spherical mirrors with a curvature radius of 100 mm placed in a confocal configuration. In this way a 500 μ m waist collimated input beam is focused with the correct spot size in the middle of the two mirrors. The second mirror collimates back the beam at its original size.

The other important parameter is the choice of the reflectivity of the input coupler (IC) mirror. The enhancement factor \mathcal{K} , defined as the ratio between the optical power circulating in the cavity and the injected power can be expressed, assuming that only the IC has a reflectivity R smaller than 1, as

$$\mathcal{K} = \frac{1-R}{\left(1 - \sqrt{R(1 - \eta_{\text{SHG}})(1 - Loss)}\right)^2}$$
(3.2)

where η_{SHG} is the single passage conversion efficiency and Loss are the losses in the cavity not due to the conversion. The maximum of \mathcal{K} is $\mathcal{K} = 1/((1 - \eta_{\text{SHG}})(1 - Loss))$ and is reached for:

$$R_{\rm opt} = (1 - \text{Loss})(1 - \eta_{\rm SHG}) \tag{3.3}$$

Since it is not possible to exactly estimate the losses and the conversion efficiency we are going to obtain, we decided to test two IC with reflectivity of 95% and 98% respectively. In the end we decided to use 98% one.

3.1.3 SHG characterization

The output power and the conversion efficiency are the most important characterization of the SHG. To perform this measurements the injection power is measured before the input window of the vacuum box and the output power out of the box. Since the shape of the Ti:S beam may change varying the pump power affecting the coupling of the cavity, we vary the 842 nm power deviating some light with a $\lambda/2$ wave plate in front of a PBS. The results of this measurement are plotted in figure 3.9. As expected for a second order process, the efficiency grows as the pump power increases and reaches 60% for the maximum power reachable by the pump laser.

Another important parameter is the stability of the lock. I tested the frequency range of the lock applying a modulation to the fast input of the Ti:S laser resonator and measured the residual modulation on the error signal. As one can see in figure 3.10 the lock shows a cut frequency of about 400 Hz. It is possible to notice an overshoot from 2 to 8 kHz probably due to a resonance of the piezo mirror mount.

3.1.4 Optical scheme for the 421 nm light

The full otical scheme of the 421 nm light is drawn in figure 3.11. Generating 421nm light is just the first step to implement the techniques described in chapter 2. In fact all the techniques based on the scattering force need a quasi resonant laser beam. The long term stability and the accuracy in the wavelength regulation of the laser itself are far from the required precision for this applications. For this reason the laser needs to be locked to a frequency reference.

But this is not enough since different and indipendently tunable frequencies are needed for the different purposes. The Zeeman slower light need to be red shifted by 1 GHz from the atomic transition while for the transverse cooling a detuning of the order of the linewidth is enough. The imaging light instead have to be resonant with the transition. To obtain this set of lights with slightly different frequency we plan to use the device called acousto optic modulator (AOM).

An AOM is composed by a transparent crystal with a piezo electric actuator fixed to one extremity and a vibration absorber on the other side. The oscillation of the piezo produces pressure waves inside the crystal that correspond to a periodic variation of the refractive index that travels inside the crystal. This periodic structure works like a diffraction grating but, since the structure is traveling in the crystal, the light is not only deviated



Figure 3.9: Above: the 421nm power in function of the pump power. Below: the conversion efficiency in function of the pump power.



Figure 3.10: residual noise on the error signal generated by the differential photodiode of the Hansch-Couillaud lock. The peak- to peak voltage is plotted in red and the RMS voltage in black. The horizontal lines represents the unlocked value.

according to the Bragg law $\sin \theta = m\lambda f/c_s$, where c_s is the speed of sound in the crystal, but also gets a frequency shift of mf depending on the order of diffraction m and the frequency of the piezo modulation f. Typically the work frequency for an AOM is not freely tunable but it has to be in a small range of frequency for which the AOM is designed. We choose to use Crystal Technologies AOM (AOMO 3350-125) that can work at (350 ± 50) MHz.

With this choice to obtain the frequency shift of 1 GHz we need to pass 3 times in an AOM. To do this we lock the laser light to the atomic line after a double passage in an AOM so that the light that exits the SHG cavity is red detuned of 350×2 MHz with respect to the atomic transition. The Zeeman slower light is then obtained shifting again with a single passage in another AOM to obtain a total shift of 1050 MHz tunable in a range of ± 150 MHz. The light for the transverse cooling and the imaging is shifted back with a double pass AOM to be resonant with the atomic transition. This shifting setup may result complicated but allows to have the transverse cooling light independent of the imaging light and the transition and both to be tunable in a range of 100 MHz. To provide the modulation signal to the AOMs we amplify 4 drivers (one for each AOM) designed by the LENS electronics workshop and specifically modified for the purpose.

The drivers provides to inputs to modulate both the amplitude and the frequency of the generated radiofrequency. that can be used to tune the



Figure 3.11: Optical scheme of the 421 nm laser. Only the main optical elements are reported. The frequency indicated beside the AOM is the total frequency shift due to the single or double pass in the AOM. $\lambda/2$ waveplates are indicated in green while $\lambda/4$ are in red.

frequency shift of the laser and its power.

3.1.5 Atomic line lock

The lock of the laser frequency on the atomic transition uses the Pound-Drever-Hall technique described in section 2.5.2. The frequency reference is provided by modulation transfer spectroscopy of the atomic transition [26]. The optical scheme for the locking is outlined in figure 3.12: let me describe this setup in detail. We performed saturated absorption spectroscopy, using



Figure 3.12: Optical scheme of the Pound-Drever-Hall locking scheme. The pump beam is in blu and the probe beam is in green.

a Madatech see through hollow-cathode lamp as a source of Dysprosium vapor shown in figure 3.13. A hollow-cathode lamp consists of a glass tube



Figure 3.13: A picture of the hollow cathode lamp.

filled with a buffer gas containing an anode shaped as a disc with an aperture and a hollow cylindrical cathode in order to let the spectroscopy laser beam pass through them. The cathode has a deposition of Dysprosium on its inner surface. When an high voltage is applied to the electrodes the buffer gas ionizes and the ions are attracted to the cathode. When the ions impact the cathode they sputter Dysprosium atoms producing a Dysprosium gas in the region inside the cathode. We choose to use a source of Dysprosium independent of the effusive cell of the experiment because, in this way, we are sure to not interfere with the atomic beam and we can reduce the number of optics on the vacuum setup table.

The probe beam is the s-polarized component of the light inciding on the PBS. It simply goes through the hollow cathode lamp and is read by a fast photodiode. The pump beam is instead the p-polarized component that passes through the PBSs. This beam passes through an Electro Optical Modulator (EOM) that introduces a phase modulation before being sent on the atoms. The relative balance between pump and probe can be adjusted by rotating the $\lambda/2$ wave plate in front of the PBS.

An EOM is a device used to modulate the phase or the polarization of a laser beam. It is basically made of a crystal that changes its refraction index along an optical axes when an electric field is applied. If the light is polarized along the optical axes its phase changes when an electric field is applied because the refraction index changes. If the polarization of light is tilted with respect to the optical axes the EOM acts like a variable wave plate changing the polarization from linear to elliptical or rotating it. Usually to achieve a π phase shift very high voltages are needed but if the applied electric field is sinusoidal it is possible to realize resonant EOMs that require much lower amplitude of the fields to perform a consistent phase shifting.

We choose a 12.5 MHz resonating EOM because it is one of the higher frequencies achievable with a cost effective device, of course the higher is the frequency, the higher is the band of the feedback system. We oriented the EOM to impress only a phase modulation without rotating the polarization of the beam. We use a setup called modulation transfer spectroscopy because the modulation that creates the side-bands on the pump beam is "transferred" on the probe beam [26], in fact the EOM modulates only the pump beam. This has the great advantage of preventing any spurious amplitude oscillation on the probe beams due to the non ideal orientation of the EOM, since we are extracting the signal right from the amplitude oscillation of the probe beam. An amplitude modulation on the pump beam instead is less critical because it is around the saturation intensity so it doesn't change the shape of the dip for small fluctuations. The phase modulation introduced by the EOM creates sidebands of the main frequency and can be used for the Pound-Drever-Hall locking.

Let us now calculate the response function $F(\omega)$ of this spectroscopy signal to put it in equation 2.57. We can suppose that saturated spectroscopy [20] only sees the natural linewidth since the probe beam is well below I_{sat} . In this case we can assume

$$F(\omega) \propto I \frac{\Gamma/2\pi}{(\omega - \omega_0)^2 + (\Gamma/2)^2}$$
(3.4)



Figure 3.14: Simulation of the Pound-Drever-Hall error signal for a single transition of width $\Gamma = 32.2 \times 2\pi$ MHz. In red it is plotted the natural linewidth and in blue the expected error signal with a modulation frequency of $\Omega = 12.5 \times 2\pi$ MHz.

Using this response function in equation 2.57 we obtain the error signal plotted in figure 3.14

I assembled the spectroscopy photodiode because there was no cost effective detector for our purposes. The detection system consists in an amplified photodiode with two outputs: a low frequency one and a high frequency one. I used a high speed photodiode with high responsivity in the blue region. Whose reverse bias current is filtered by a crossover filter at the frequency of 8 MHz. The crossover is a two branches filter that splits the input signal in an high frequency part and a low frequency part. The high frequency part is then amplified by a high frequency current to voltage converter configuration. This output is used for the Pound-Drever-Hall circuit it is non zero only when the pump and the probe beam are spatially overlapped and the phase modulation is on. The low frequency is also amplified and it is used for alignment purposes since it is possible to see both the Doppler profile of the line when the pump and probe beams are overlapped. In figure 3.15 it is possible to see the measured PDH signal. It is important to notice that this feedback uses the 421 nm light to generate the error signal but the feedback signal is sent to the reference cavity of the Ti:S so that this locking system is not independent on the SHG cavity lock. The SHG lock keeps the cavity in resonance with the Ti:S light, which is referenced to the atomic transition.

To test the efficiency of this lock we can modulate the light frequency through the AOM that shift the spectroscopy light, and measure the residual fluctuation on the error signal. The results are shown in figure 3.16.

It is possible to find an upper limit for the residual linewidth after locking measuring the amplitude of the error signal and converting it in frequency. We calibrate the voltage-frequency conversion measuring the distance be-



Figure 3.15: Error signal of Pound-Drever-Hall locking scheme measured scanning the laser frequency with the external cavity input of the Ti:s laser.



Figure 3.16: Residual noise on the Pound-Drever-Hall locking system in function of the frequency. The residual noise with the system locked is plotted in black while in red is plotted the modulation with the system unlocked.

tween the spectroscopy signal of the two bosons. Then we measured the residual electrical noise on the Pound-Drever-Hall error signal that is composed of two components: the residual frequency fluctuation and the electronics noise. By doing so we find that the residual laser linewidth is below 10 MHz.

3.2 Vacuum setup assembling

To guarantee easy optical access to the vacuum setup we decided to build a structure that suspends the system center 40 cm above the optical table. The structure is built using 45 mm section aluminum bars to keep the system stable without encumbering. The structure we built is totally modular to allow any future modification to the vacuum setup, and has only 6 legs in order to leave as much free space as possible on the optical table. Also we realized a raised breadboard around the final cell to provide a stable support for the optics and to work on two levels.

The realized setup is very close to the one designed in figure 1.7 except for the final part. Instead of the final octagonal chamber we mounted a 6-ways cross mainly to avoid delay in the experiment. In fact we were ready to assemble the vacuum setup except for the final cell so we decided to start mounting a test setup for first testing all the apparatus. This approach will force us to open the vacuum system and to perform a second baking for the first configuration, but we have made some useful tests and measurements that will help us to improve the experiment. The design of the test apparatus and its picture can be seen in figure 3.3 and figure 3.2.

3.2.1 Baking

In a vacuum setup there are many sources of gases and vapors that make difficult to reach ultra high vacuum (UHV) pressures. Let us take a brief overview on the mechanisms that limit the ultimate pressure reachable in a vacuum apparatus.

Every element or compound has a different vapor pressure in function of the temperature. It is clear that if an compound with a given vapor pressure is present in the vacuum system it is impossible to go far below that pressure. Also the walls of the apparatus have been exposed to air and to manufacturing materials before being enclosed in the apparatus. Some of this materials will be absorbed in the walls and slowly released when the pressure drops. An effective way to prevent these phenomena is to accurately clean all the vacuum components before assembling them but the final solution is to bake the apparatus. The baking procedure simply consist in heating the apparatus while pumping the vacuum. In this way the vapor pressure of any material is increased so the evaporation accelerates drastically when the vapor pressure is above the system pressure that is kept low by the vacuum pumps. For absorbed material the situation is cumbersome but from a qualitative point of view it is possible to say that desorption accelerates when the temperature is high enough to overcome the adhesion force. The two critical materials to evaporate are hydrogen and water. Hydrogen is the dominant gas that is desorbed for stainless steel [27] and it is easily desorbed for temperature above 200-250 °C; water is the most common vapor in air that can not be desorbed at room temperature. It is important to notice that if a small surface inside the vacuum setup remains unbaked its outgassing will be dominant compared to all the rest of the walls.

The other sources of gas that may limit the minimum pressure reachable in a vacuum setup are leaks from outside the apparatus or from closed regions inside the apparatus and the backstreaming from the vacuum pumps. Closed regions inside the setup have been carefully avoided cutting a side of all threads and drilling an exhaust hole in every housing mounted inside the vacuum system. To be sure that our setup has no leak from outside we performed a leak test using a RGA (residual gas analyzer) and helium. A RGA is substantially a mass spectrometer that analyzes the gas inside the vacuum and allows to detect also oil and solvents contamination. The leak test consists in spraying helium in various parts outside the apparatus and check with the RGA if the relative helium pressure rises. Helium is used because it is the most permeable gas thus reveals even the smallest leak.

We use two Varian SH-110 scroll pumps $(100 \,\ell/\text{min}; \text{ultimate pressure of } 5 \cdot 10^{-2} \text{ Torr})$ as forepumps for two Varian Turbo-V 81 M turbomolecular pumps $(80 \,\ell/\text{s}; \text{maximum inlet pressure } 13 \text{ Torr}; \text{ultimate pressure of } 3.8 \cdot 10^{-10} \text{ Torr}).$

For the final vacuum we use 2 Varian VacIon Plus 20 ion pumps $(20 \,\ell/s;$ maximum starting pressure $5 \cdot 10^{-2}$ Torr; ultimate pressure below 10^{-11} Torr) and a NEXTorr D 100-5 hybrid pump ($100 \,\ell/s$ getter + $6 \,\ell/s$ ion pump).

To reach the UHV we first activated the Scroll pumps to reach a pressure of 10^{-1} Torr then we activated the turbomolecular pumps. So the pressure dropped down to 10^{-5} Torr in 5 minutes; 10^{-6} Torr in an hour and 10^{-7} Torr in 12 hours and $5 \cdot 10^{-8}$ Torr in 3 days. While the pressure was dropping we prepared the apparatus for the baking winding all the parts with heating tapes and then covering up with various layers of aluminum foil to prevent heat dissipation. in figure 3.17 it is possible to see a picture taken during the baking. We baked the apparatus for 8 days at the temperatures reported in table 3.3. We decided to keep the Zeeman slower far under the maximum temperature reachable because the different thermal expansion of the glue, the copper of the coils and the steel may cause damages to the coils. Also we raised the effusive cell temperature to $800 \,^{\circ}\text{C}$ to let all the impurities and manufacturing residuum evaporate when the vacuum pumps are still on. During the baking the pressure raised from $5 \cdot 10^{-8}$ Torr to 10^{-5} Torr at the beginning of baking and went back to $5 \cdot 10^{-8}$ Torr at the end of baking.



Figure 3.17: The vacuum apparatus enveloped in aluminum foil during the baking. On the right it is possible to see one of the pumping benches composed by a turbomolecular pump, a scroll pump and a RGA.

We performed the activation procedures of the final pumps while the apparatus is still hot and the turbo pumps are attached to the setup. In facts ion pumps release gases upon starting if they have been exposed to air so we cooled down the ion pumps to $100 \,^{\circ}$ C, then we mounted the magnets and put them on and off until the pressure stopped rising at the activation. Also the getter part of the NEXTorr pump needs to be activated heating the getter element to very high temperatures. During this process the getter release a lot of vapor thus the turbo pumps need to be still connected to draw all the residual material evaporated.

Once all these operations are concluded we closed the values to the turbo pumps, activated the ion pumps and cooled down the apparatus to room temperature. The pressure drops very fast in the whole apparatus to the minimum pressure readable by the ion pumps corresponding to 10^{-10} Torr.

After the baking we notice that all the view ports of the transverse cooling cell where coated everywhere except than on the oven side. This coating is likely caused by Dysprosium effused by the cell because we underestimated the effusion temperature and the divergence of the tails of the atomic beam effused by the cell. A small number of dysprosium atoms are enough to produce a coating with a very high optical density. This coating prevented us the possibility to test the transverse cooling. In the final configuration of the apparatus we will add a cold skimmer right in front of the oven to prevent any atom to reach the viewport surface.

3.2.2 Differential pumping

We have seen in the previous section that it is possible to reach UHV pressures in our apparatus but, when the oven reaches its operative temperature of 1250 °C, we expect the pressure in the first part of the apparatus to rise.

Part	$T_{\rm bake}$	$T_{\rm max}$
Vacuum tubes	450 °C	$280^{\circ}\mathrm{C}$
Bellow	$80^{\circ}\mathrm{C*}$	80°C
NEXTorr	$250^{\circ}\mathrm{C}^{\dagger}$	200 °C
Ion pumps	$350^{\circ}\mathrm{C}^{\dagger}$	200 °C
Valves	$300^{\circ}\mathrm{C}$	$170^{\circ}\mathrm{C}$
View ports	$300^{\circ}\mathrm{C}^{\star}$	190 °C
Zeeman slower glue	$260^{\circ}\mathrm{C}$	160 °C

Table 3.3: Baking temperatures for the various components of the experiment. T_{bake} is the maximum baking temperature indicated by the manufacturer, T_{max} is the maximum temperature we reached during the baking. * Higher temperatures are possible but this reduces the flexibility. \dagger Without magnets, with magnets the maximum temperature is 150° C.* with a maximimum temperature variation of 2° C/min

To prevent this higher pressure to affect the final part of the apparatus we design a differential pumping setup. Before looking at the details of the design let me introduce the basis of the study of gas flow [27, chap. 3]. In an UHV apparatus we can study the gas flow in the molecular approximation. This approximation neglects the collisions between particles and consider only the interaction with the walls. A particle that hits a wall is then reemitted with random direction and a speed distribution corresponding to the temperature of the wall. This approach is correct when the Knudsen number $K_n \equiv \lambda_{mfp}/d > 1$, where λ_{mfp} is the mean free path of a gas particle and d is the typical dimension of the vacuum apparatus. In UHV conditions λ_{mfp} may be greater than 1 m while the tube diameter of our apparatus is 4 cm so this condition is easily satisfied. Let us now define some useful quantities: the throughput is the quantity of gas that passes a plane in in an unit of time

$$Q \equiv \frac{d}{dt}(PV) \tag{3.5}$$

The throughput of a tube is linearly dependent on the pressure difference at the ends so it is possible to define the conductance C as

$$Q = C(P_2 - P_1) (3.6)$$

This equation can be compared to the Ohm law for electricity I = CV where the throughput have the role of the current and the pressure is analogous to the voltage. The last quantity we need is the pumping speed defined as $S \equiv Q/P$.

Let us now schematize our vacuum apparatus as in figure 3.18

The conductance of the channels separating the chambers can be calculated as the conductance of a long tube in the molecular flow approximation:



Figure 3.18: Differential pumping scheme

$$C_{\rm tube} = \frac{\pi}{12} v \frac{d^3}{l} \tag{3.7}$$

where v is the thermal average velocity of the gas, d is the tube diameter and $l \gg d$ is the tube length. In our case the each tube is 5 mm in diameter and 81 mm in leght. The conductance for H₂ at 25 °C is 0.71/s.

For every chamber at the stationary state it is possible to write an equation setting the incoming throughput equal to the outgoing throughput, We obtain a system of 3 equation that can be solved fixing some parameters. The pumping speed is known, and we can neglect the desorption in the differential pumping stage since its surface is very small. Assuming that desorption in the final cell is $Q_{\text{des}} = 101/\text{s} \cdot 10^{-11}$ Torr, to keep the pressure of $P_3 = 10^{-11}$ Torr in the final cell we find the maximum allowed pressures in the other sections of the setup:

$$P_2 = 1.3 \cdot 10^{-9} \,\mathrm{Torr} \tag{3.8}$$

$$P_1 = 4 \cdot 10^{-8} \,\text{Torr} \qquad L_1 = 8 \cdot 10^{-7} \,\text{Torr} \cdot 1/\text{s}$$
 (3.9)

Where L_1 is the maximum desorption rate in the oven cell. The relative high pressure in the transverse cooling cell could be in principle measured from the ion pump current. The pressure measurement we read is compatible with the allowed one, even when the effusive cell is set to it maximum temperature, even if the reading of the ion pump is not completely trustable at this low pressures.

Inside the differential pumping section we include a rotary shutter to block the atomic beam when it is not not necessary in the final cell. In fact we plan to keep the Zeeman slower on and the beam shutter open until the MOT is loaded, then the next cooling processes will be performed turning off the magnetic field of the Zeeman slower , the Zeeman slower light and closing the beam shutter because fields light and atoms would be disturbing in these phases of the experiment.

We designed the shutter as a wobble stick moved by a pneumatic rotary drive. The shutter stop position can be regulated mechanically while the rotation speed can be regulated by flow control valves. I designed the remote control for the shutter using a pneumatic valve 5 ways, 3 positions. I designed an adapter to drive the high power signal 24 V 2 A that control the valve with a standard TTL output. For test purposes I programmed an Arduino Uno board to work as a temporary control unit for the shutter.

3.2.3 Effusive cell design and atomic beam measurements

In this paragraph I will describe the design of the effusive cell but first let me outline some characteristics of the atomic beam effused by a cell. We consider a hot chamber containing an ideal gas connected through a channel of diameter d and length L to a cool chamber that we can consider empty. Considering a density low enough to use the molecular flow approximation $(K_n = \lambda_{mfp}/d \ll 1)$. The flow intensity to the empty chamber in a given direction θ can be calculated summing the intensity that comes directly from the back part of the hot chamber with the intensity that arises from the collisions with the walls of the channel. [28, chap 2.4]. The result is:

$$I(\theta) = \nu_0 r^2 j(\theta) \tag{3.10}$$

Where ν_0 is the rate per unit area at which particles enter the channel entrance, r is the channel radius and $j(\theta)$ is the angular distribution. When $\tan \theta \ge d/L$ only particles colliding with the channel walls contribute to the flux and $j(\theta)$ can be expressed as:

$$j(\theta) = \cos(\theta) \left[\alpha + \frac{4(1-2\alpha)}{3\pi q} \right]$$
(3.11)

When $\tan \theta < d/L$ instead there are contribution from both the channel walls and the chamber so:

$$j(\theta) = \alpha \cos(\theta) + \frac{2}{\pi} \cos(\theta) \left[(1 - \alpha)R(q) + \frac{2(1 - 2\alpha)}{3q} \left(1 - \sqrt{(1 - q^2)^3} \right) \right].$$
(3.12)

In this expression we define:

$$q = \tan \theta / \beta \tag{3.13}$$

$$\beta = d/L \tag{3.14}$$

$$R(q) = \arccos(q) - q\sqrt{1 - q^2}$$
(3.15)

$$\alpha(\beta) = \frac{1}{2} - \frac{1}{3\beta^2} \left[\frac{1 - 2\beta^3 + (2\beta^2 - 1)\sqrt{1 + \beta^2}}{\sqrt{1 + \beta^2} - \beta^2 \left\{ \ln\left(1 + \sqrt{1 + \beta^2}\right) - \ln\beta \right\}} \right].$$
 (3.16)



Figure 3.19: Polar plot of the expected flux density. In black the flux density of a 4 mm aperture and in red for a channel 4 mm in diameter and 30 mm long. The two plots are normalized to the same value.

Despite the equations above may result quite cumbersome it is possible to notice that the presence of the tube does not change the front directed intensity, in fact for $\theta = 0$ there is no dependence on the value of q. Only the crooked flux is heavily reduced as shown in figure 3.19. This increases the duration of the material in the crucible and also avoid the effused vapor to go in unwanted directions. We decided to use a crucible shown in figure 3.20



Figure 3.20: A picture of the tantalium crucible unmounted. From the left: the crucible chamber, the collimating tube and the crucible protection.

with a diameter of 4 mm and a length of 30 mm corresponding to a $\beta = 0.1\overline{3}$. It is also possible to calculate the total flux reduction W integrating the expression for $I(\omega)$ over the solid angle.

$$W = 1 + \frac{1 - \sqrt{1 + \beta^2}}{\beta^2} + \frac{2\left[1 - 2\beta^3 + (2\beta^2 - 1)\sqrt{1 + \beta^2}\right]^2}{9\beta^4 \left[\sqrt{1 + \beta^2} - \beta^2 \left\{\ln\left(1 + \sqrt{1 + \beta^2}\right) - \ln\beta\right\}\right]}$$
(3.17)

According to equation 3.17 the total flux is reduced by a factor of 7 with respect to the case of a simple aperture at the given oven temperature.

The effusive cell is a commercial double filament cell by SVT Associates Inc. The effusive cell have to be heated slowly to avoid temperature differences that can crack the crucible due to different thermal expansion in different parts. The temperature of the oven is controlled by two digital PIDs, one for the tip filament and one for the main filament. Two type K termocouple provides the feedback signals for each PID that controls a current generator to remain at the given temperature or to ramp without exceeding the heating or cooling rate of 20 °C/s. There are two filaments to keep the two parts of the oven at different temperatures, in particular we want to keep the tip slightly hotter than the dysprosium chamber to avoid clogging in the collimation tube.

We performed two kinds of measurements to characterize the divergence of the atoms emitted by the cell and the flux.

We estimate the angular distribution of the atomic beam by shining a probe laser beam from the 45° viewports of the transverse cooling cell. The beam has a 1 mm diameter and it is 23 ± 2 mm far from the end of the crucible. On the upper part of the apparatus an imaging system with a reduction factor of 10 focuses the image of the fluorescence on a CCD camera. To calibrate the pixel dimension we count the pixels occupied by the image of the movable beamsplitter that we assembled in the cavity to have control the oven status and to align the Zeeman slower beam (whose dimensions are known).

It is important to remember that all the viewports of the transverse cooling chamber have been coated by Dysprosium during the baking, however a little slice on the oven side has remained uncoated due to the internal geometry of the chamber that have protected them. We used this area to perform measurements even if the optical access is very small.

The frequency of this beam can be tuned using the reference cavity controls of the SolsTi:S since the laser is not locked to the atomic line. We set the same frequency on the spectroscopy AOM and the probe beam AOM in order to use the spectroscopy reference to measure the detuning of the beam from the transition. The frequency of the AOM drivers is stable below 0.1 Mhz and drift less than 0.5 Mhz/h so that it can be considered constant over the measurement. To calibrate the frequency axes we use the isotopic shift between the two bosons 164 Dy and 162 Dy for the 421 nm transition that is known from literature [19] to be 913.2 \pm 0.8 MHz.

Changing the frequencies we measure different spatial profiles of fluorescence, the CCD images are reported in figure 3.21. We record the peak position and intensity of the of the strongest ¹⁶⁴Dy isotope signal. Knowing the distance between the end of the collimation channel and the probe beam we can convert the position in angle considering also that the probe beam is inclinated with respect to the imaging direction of 45 °C. In this way a first measurement of the longitudinal atoms velocity v_z is possible. Assuming that the speed distribution is independent on the angle we can convert the frequency Doppler shift to the corresponding velocity. Data are plotted in figure 3.22. We expect to find a linear dependence: $v_{transv} = v_z \cdot \sin \theta$. We fit the data for a crucible temperature $T_{\rm tip} = 1050\,^{\circ}{\rm C}$ with a line and we extract a beam velocity of $v_z = (364 \pm 11) \text{ m/s}$ that has to be compared with the expected $v_{\rm theo} = 450 \,\mathrm{m/s}$. The difference in the values clearly suggests that the assumption of independence of the speed distribution on the angle is not valid. This is probably mainly due to fact that the crucible aperture size is not negligible with respect to the distance at which the measurement is taken.

Another important information we can extract for this measurement is the angular divergence of the atomic beam, in fact plotting the measured fluorescence intensity as a function of the angle we find the results in figure 3.23. It is possible to appreciate the good agreement of the measured data with the theoretical calculation, however for higher angles the measured data seem to be under the theoretical values, however this point in the graph corresponds to angles at the borders of the viewport and away from the axis of the optical system thus it is possible that light losses make those points non reliable.

We can define the beam divergence as the HWHM of the intensity profile. In this case the divergence from an aperture is 60° while our tube should provide a divergence of 6.3° . As we can see from the measured data in figure 3.23 the collimation channel works as expected.

Another interesting quantity to experimentally estimate is the atomic flux out of the oven. To the aim we perform an absorption measurement [29]. The setup for this measurements is here described. We split a beam with a non polarizing beam splitter, one part is directly revealed by a photodiode, the other is enlarged to a waist of 3.5 mm, sent through the CF16 viewports placed after the transverse cooling cell and finally is detected by another photodiode. The two photodiodes are read by a differential amplifier. We measure the absorption for different oven temperatures from 950 °C, which is the minimum temperature that gives an appreciable signal, to 1300 °C. The results are shown in figure 3.24

It is possible to derive the flux from this measurement even if we need to make some rough approximations. We assume an uniform and collimated



Figure 3.21: Fluorescence images taken with the CCD camera for frequency detuning from -225 MHz to +225 MHz in steps of 45 MHz. Note that the probe beam is oriented with an angle of 45° with respect to the observation direction. In this picture the atomic beam is traveling from the top to the bottom and the probe beam from the right to the left.



Figure 3.22: Transverse velocity as a function of the angle of detection with respect to the atomic beam axis. The fit error is within the data points.



Figure 3.23: Normalized flux intensity of the atomic beam effused by the oven in function of the angle. In black the profile for a 4 mm hole, in red the profile for a tube with 4 mm diameter and 3 cm legth. In blue the measurement in our setup. The fit error is within the data points.



Figure 3.24: Measured absorption for various temperatures of the effusive cell

atomic beam and consider the atom as a two levels system. If an atom absorbs a photon, after a lifetime τ the photon will be re emitted in a random direction, thus we assume that every absorbed photon is lost from the laser beam. Thus the absorbed light is proportional to the atoms number in the beam volume. For laser beam under the saturation intensity the Beer-Lambert law describes the transmission of light through the atomic beam:

$$I = I_0 e^{-D} (3.18)$$

Where D is the optical density of the sample and can be expressed in function of the atomic density, the absorption cross section σ_0 and the length of interaction that in our case is the atomic beam diameter 2r.

$$D = n\sigma_0 2r \tag{3.19}$$

We take the following expression for the cross section:

$$\sigma_0 = \frac{3\lambda^2}{2\pi} \tag{3.20}$$

Under these approximations we can estimate the atomic density and then calculate the flux with the following expression:

$$\phi = \pi r^2 u n \tag{3.21}$$

where we are assuming the all the atom are travelling at the most probable speed in the beam $u = \sqrt{3k_BT/m}$. The results are plotted in figure 3.25.



Figure 3.25: Atomic beam flux: in blue the total flux calculated for our setup, in red the flux estimated from the absorption measurements.

The values are more the one order of magnitude smaller than the theoretical calculation above. However we are measuring only the flux of a single isotope, also we are neglecting all the atoms that does not arrive in the laser region.



Figure 3.26: Spectroscopy of the 421 nm transition. In blue the measured data averaged over 32 scan, in red the fit with a series of Lorentzians of fixed FWHM=37.9 MHz. The two highest peaks correspond to the two bosons 164 Dy and 162 Dy, while the smaller ones come from the iperfine splitting of the fermionic isotopes 163 Dy and 161 Dy.

To conclude this section I show a fluorescence transverse measurement of the 421 nm transition we performed at the end of the Zeeman slower tube. The atoms that can reach this point of the apparatus have a very low transversal velocity. This allow us to obtain a nice spectroscopic signal, better than the one reported in figure 1.6 by [19]. From this measurement we can have precise calibration of the laser frequency and, for example, estimate the isotopic abundance of our sample.

A laser beam crosses the atomic beam perpendicularly, while a photomultiplier collects the fluorescence light from above the 6-ways cross. We scan the laser across the transition averaging over 32 measurements to enhance the Signal to Noise Ratio (SNR) keeping the slow drifts of the laser and AOM frequencies negligible during the measurement.

To fit the data, shown in figure 3.26, we used a comb of Lorentzian functions instead of Voigt functions because we expect the Doppler Broadening to be negligible for transverse spectroscopy compared to the natural linewidth. Fitting the highest peaks corresponding to the bosonic ¹⁶⁴Dy, we measure a FWHM linewidth of $2\pi \cdot 37.9 \pm 0.2$ MHz, which is slightly larger than the natural linewidth. This can be due to the laser intensity which is close to the saturation intensity, to the residual transverse Doppler broadening and to an imperfect orthogonality between the laser beam and the probe beam. We then fit the whole spectrum using a series of Lorentzians fixing the FWHM to the found value in order to reduce the number of fitting parameters.

We took another measurement with the same setup discussed above,



Figure 3.27: Spectroscopy of the 421 nm transition performed under an angle of $(13 \pm 2)^{\circ}$. In blue the measured data averaged over 32 scans, in red the convolution between the series of lorentians found in the fit of figure 3.26 and the thermal distribution of an atomic beam at $T = 950^{\circ}$ C

but tilting the probe beam by an angle α of $(13 \pm 2)^{\circ}$ with respect to the direction orthogonal to the atomic beam. In this way the atomic longitudinal velocity has a component along the direction of the laser beam equal to $v_z \sin(\alpha)$. In this configuration we expect the transition to show a blue shift due to the Doppler effect and we expect every line to show a frequency profile corresponding to the velocity distribution of the atomic beam. In fact to reproduce the results of this measurement shown in figure 3.27 for $T_{\rm oven} = 1000 \,^{\circ}{\rm C}$ we have to perform the convolution between the comb of Lorentzians found in the orthogonal case with the thermal frequency distribution of the atomic beam. This convolution is too complex to be used as a fitting function by Mathematica so we impose by hand different temperatures for the distribution finding that the best one to reproduce the data is $T = 950 \,^{\circ}{\rm C}$, however the significant error in the angle measurements has to be taken into account in this measurements.

3.2.4 Transverse cooling

The transverse cooling (TC), even if not essential for the operation of the experiment, allows the atomic beam effused from the oven to be collimated. This allows to reducing atoms losses and working at lowen oven temperatures mantaining an adequate flux to the MOT cell. In this way we can extend the duration of the atomic sample in the oven and postpone its reloading.

The transverse cooling is a 2D molasses, realized with to retroreflected beams entering in the TC cell from the two inclined windows on the bottom. As explained in section 2.3.1 we would like to employ beams red detuned
with respect to the atomic transition and with intensity above $I_{\rm sat}$ in order to apply the maximum viscous force possible. We use elliptical beams to enhance the intensity on the atoms. In fact all the atoms that are already outside the 5 mm diameter of the differential pumping tube can not be pushed back in position so it is useless to use a width of the beam larger than 5 mm. while it is convenient to use all the length of the optical access that is 40 mm.

If we define the average intensity as:

$$I_{\text{avg}} \equiv \frac{P}{\pi w_x w_y} \tag{3.22}$$

The power required to reach an average intensity equal to I_{sat} is about $90 \,\mathrm{mW}$ per beam.

To estimate the range of transverse velocities that are collimated we can suppose to apply the maximum force to the atoms for all the time t they stay in the transverse cooling beams. Assuming the most probable speed of the beam to be 490 m/s we find $t = 82 \,\mu\text{s}$. During this interval the atoms undergoes an acceleration that can be calculated using the equation 2.3.

$$a_{\rm max} = \frac{F_{\rm max}}{m} \simeq 9.3 \cdot 10^4 \,{\rm m/s^2}$$
 (3.23)

Thus the maximum transverse velocity that can be collimated by the TC is $v_{\text{max}} = a_{\text{max}}t = 7.6 \text{ m/s}$ corresponding to a divergence of approximately 1°.

3.2.5 Zeeman slower

In section 2.3.2 I have described the working principle of the Zeeman slower. Here I'm going to describe how we realized our Zeeman slower. We produce the proper magnetic fields by a collection of coils.

To calculate the magnetic field let us start from the Biot-Savart law:

$$d\mathbf{B} = \frac{\mu_0 I}{4\pi} \frac{d\mathbf{l} \times \hat{r}}{r^2} \tag{3.24}$$

Where I is the electrical current. Integrating this equation over a single turn supposing to be in the center of the turn we find:

$$B_z = \frac{\mu_0 I}{2} \frac{a^2}{(a^2 + z^2)^{3/2}} \tag{3.25}$$

The subscript z indicates the in the center of th coil the field is orthogonal to the coil plane. It is important to notice that in the center of the coil the dependence of the magnetic field on the radius a is $\sim 1/a$, so the smaller the coil is, the higher is the field produced with the same current. Also there is another advantage in using smaller coils, in fact the resistance of a coil

scales as $R = \rho l = \rho 2\pi a^2$, where ρ is the resistivity and l the length of the coil, so the dissipated power $P = RI^2$ scales as a^2 . To obtain the total field we simply have to sum over the collection of coils we employ.

Let us now focus on the shape of field needed to have an efficient slowing of the atoms. The force applied by the Zeeman slower is the scattering force expressed in eq. 2.2 where in this case the effective detuning from the transition is given by equation 2.27. First we have to calculate the magnetic moment in equation 2.27 using the known Landè factors [30]. Dysprosium has a ground state with J = 8 and the cooling transition links to a state with J = 9 so the Zeeman effect splits all these levels in a more complex structure background with respect to the simple example I made in section 2.3.2. However is easy to be convinced that the slower light will cycle only between the two higher (or lower) m_J states. In fact a σ^+ photon produce a transition with $\Delta m_J = +1$ while the atom will decade with the same probability with $\Delta m_J = -1, 0, +1$. After a few scattered photons the atomic beam will be totally polarized and will be cycling between the two stretched states. This is a fundamental point because the Landé factor is slightly different for each transition and therefore the slower light will be optimized only for the stretched transition. In the spin flip part the atoms have to change their spin from $m_I = +J$ to $m_I = -J$. In that point of the slower the magnetic field cross 0 so the light is nearly resonant with all the Zeeman levels, also the different Landé factor will be negligible. Considering that the emission process has random change in the m_J while absorbtion changes m_{I} by an unite in the desired direction, on average an atom needs less than 20 cycles to repolarize in the correct Zeeman level. Assuming that there is an absorption every 2τ it take 200 ns to repolarize the atoms. This process is almost instantaneous since in this interval an atom travels less than 0.1 mm. We can therefore consider the atoms in the Zeeman slower to be always cycling between the two stretched levels that support the cooling mechanism.

From the first raw calculation we realized that a pure σ^- slower would need an extremely high field at the end of the slower causing a series of technical issues. The main problem in creating such an high field arises from the power dissipated in the last coil that would be very high. Also the residual field on the final cell, where we want a accurate control on magnetic fields, will be not negligible. To get around this issue we decided to realize a so called spin flip Zeeman slower. It is composed by a first part of σ^+ slower attached to a σ^- part. This configuration produces a field profile similar to a classic σ^- slower but with an offset to negative magnetic fields. This allow us to use a magnetic field ranging from -150 G to 625 G instead of from 0 G to 775 G as would be for a pure σ^- slower. Even if this allow us to use smaller coils, a heat dissipation system is still needed.

For the purpose we realized the slower by inserting a CF16 tube in a slightly larger tube. The atoms passes in the smaller tube and the coils are winded on the external one. Water circulates in the space between the two tubes providing an efficient cooling of the coils.

To obtain the maximum deceleration possible we would like to use an optical intensity exceeding the saturation intensity. The power necessary to have the average intensity equal to the saturation one on a beam with the same diameter of the internal tube, 16 mm, is $P = I_{\text{sat}} \cdot \pi w_0^2 = 112.6 \text{ mW}$. This value is just an indication, in fact, in principle, it is better to use more than the intensity saturation. The atomic beam does not occupy the entire volume inside the slower since it is collimated by the 5 mm wide differential pumping tube, nevertheless we expect the beam to diverge inside the Zeeman slower beacuse of the transverse heating mechanism and the residual divergence. For these reasons we shape the laser beam to match the atomic beam profile. To the aim we use a focused laser beam that increases its intensity moving towards the beginning of the Zeeman slower .

To calculate the ideal field profile we set a constant slowing force $F_{\text{scatt}} = 0.4F_{\text{max}}$ in equation 2.2 and we calculate the required *B* in function of the decreasing velocity B(v) and then we integrate it over time to find B(z). We choose an acceleration lower than the maximum value to assure that all the atoms below a certain velocity will be slowed. In fact if for some reason the slowing process is not effective as expected at a certain velocity, then all the atom that have that velocity or above, are not slowed because the light will never be resonant with their velocity again. A smoother magnetic field gradient assures that all the atoms are cooled even in the presence of small imperfections on the magnetic field.

In our calculation we have to decide an initial velocity at which the slower starts to be effective. The speed distribution of an atomic beam differs from the ideal gas, in fact contains only positive velocities and the speed distribution can be expressed as [20]

$$f(v_z) = v^3 e^{\frac{-v^2}{u^2}} \tag{3.26}$$

where $u = \sqrt{2k_bT/m}$. This distribution is not limited so we have to choose a cut off velocity, all the atoms faster than that velocity will not be cooled. Of course choosing an higher initial velocity allows to slow a larger part of the speed distribution, however this would require an higher magnetic field. As discussed before there are design limits to the higher field we can realize using electric coils. In fact making a bigger coil is not so effective because increasing the number of loops makes the external loop far from the coil center. The effect of the current in the external turn is minimum, also they are longer so their resistance increases and so does the heat generation due to Joule effect. As a compromise we decided to realize the Zeeman slower for an initial speed of $1.1v_{\rm mp}$ calculated at the temperature of 1250 °C, where the most probable velocity in a beam is $v_{\rm mp} = \sqrt{3k_BT/m} = 489$ m/s [20]. This allows to slow more than 50% of the speed distribution (figure 3.28) in every reasonable working condition using affordable coils in the Zeeman slower .



Figure 3.28: The beam speed distribution for $T = 1250^{\circ}C$, the vertical line marks the speed $1.1v_{mp}$.

The result of our calculations is plotted in figure 3.29. Equation 2.2 also gives the value of the laser detuning from the atomic transition $\delta_0 = -32\Gamma = -1$ GHz: this is the value we used in section 3.1 to design the frequency shift scheme.



Figure 3.29: The ideal Zeeman slower field calculated for $a = 0.4a_{max}$, B(x = 0) = -150G, $v_{ini} = 530 \text{ m/s}$, $v_{fin} = 20 \text{ m/s}$

Now that we have all the ingredients we use a Mathematica notebook to design the Zeeman slower coils. With real coils it hard to obtain the large field gradient of the last part of the ideal field profile. The real field will therefore reach the same maximum value with smoother gradient; this leads to an increased length of the slower. In principle it is possible to use progressive coils with a continuously varying number of turn in function of z, but in practice it is very difficult to build this setup exactly as designed. This may cause imperfections in the field with the risk of a not working slower.



Figure 3.30: Scheme of the Zeeman slower project. The coils are painted in black or orange depending on the circulation direction of the current. In blue it is plotted the calculated field profile. Atoms travel from the left to the right, the laser is shined from the right.

For this reason we decided to build our Zeeman slower using a collection of separated coils placed one after the other. As shown in figure 3.30 the slower is composed by 9 coils: 2 in σ^+ configuration at the beginning and 7 in σ^- configuration at the end. The coils are usually simply connected in series, however it is possible to add a resistor in parallel to each coil in order to adjust the current between and thus the shape of the magnetic field at will. This is especially important for the last coil because tuning its field we can tune the velocity of the atoms at the end of the Zeeman slower .

There is also a compensation coil at the very beginning to avoid a residual magnetic field in the first part of the experiment where we think to perform the 626 nm light spectroscopy to get a locking signal. In the final setup this coil is probably useless because the viewports for the 626 nm spectroscopy are placed very far away from the slower where the field is almost zero, however to be conservative we decided to keep the coil in the design.

We also calculate electrical parameters of the slower as reported in table 3.4. Since the slower dissipates 130 W due to the Joule effect it is necessary to keep it cool it.

It is possible to calculate the velocity across the Zeeman slower solving the equation of motion for the calculated field, as shown in figure 3.31. All the atoms with a velocity below $1.1v_{mp}$ are almost unaffected by the slower light until they reach the position where the magnetic field makes them resonant with it. From there they stay resonant with the light and are slowed until they exit the Zeeman slower. The atoms the are faster than this

Parameter	Value
Current	$1.9\mathrm{A}$
Resistence	36Ω
Total dissipated power	$130\mathrm{W}$
Last 2 coils dissipated power	$72\mathrm{W}$
Total wire lenght	$915\mathrm{m}$

Table 3.4: Zeeman slower electrical characteristics



Figure 3.31: Trajectories in the velocity space in function of the position plotted for different initial velocities.

velocity pass through the slower substantially undisturbed. It is interesting to notice that the velocity class between 400 and 500 m/s are already slowed before they enter in the Zeeman slower field because they are resonant with the light at 0 field. Furthermore looking at the small velocity component of the distribution it is possible to see that those atoms are pushed back by the slower light because, even if far detuned from the transition, the scattering force is still present. The results of this simulation can also been presented plotting the final velocity for each initial velocity class weighted with the atomic beam speed distribution as shown in figure 3.32.



Figure 3.32: Speed distribution before and after the Zeeman slower. In blue the atomic speed distribution of a beam at $T = 1250^{\circ}C$, in pink the expected speed distribution at the end of the slower. Both distributions are normalized. The peak at $v \simeq 15$ m/s is due to the slowed atoms and it is out of the plot range.

Before installing the Zeeman slower we measured the magnetic field along the slower axis using a Hall probe to verify the conformity of the effective field to the one calculated for our design. It is possible to see in figure 3.33 that the difference between the two is always below 10 G. The residual discrepancies can be corrected changing the total current in the slower and placing suitable resistors in parallel to the coils with higher fields. However this fine tuning of the field can be done properly only looking at the MOT loading.

To zero the Zeeman slower field correspondence of the MOT we plan to use another coil on the final cell, but we are going to install it only on the final vacuum setup.

It is difficult an maybe useless to optimize the Zeeman slower parameters only on the ground of the spectroscopy we perform at the end of the slower. The proper optimization of the slowing system will be performed



Figure 3.33: Above: in red the measured magnetic field profile on the Zeeman slower axis in compared to the calculated profile in black. Below: The difference between the two. The sensitivity of the Hall probe used is 100 mG, the scattering of the measured points is probably due to radial positioning uncertainties.

when the MOT is assembled and the two parts can be mutually adapted one to the other to maximize the MOT loading. However we performed some measurements of the atomic velocity at the end of the slower to verify that the whole system is working properly.

To verify if the Zeeman slower is working we perform a fluorescence spectroscopy at the end of the Zeeman slower with the probe beam tilted with respect to the direction orthogonal to the atomic beam. In this way, in principle we can measure the longitudinal velocity of the beam calculating the deconvolution of the measured profile with the natural line shape. In fact with this calculation it would be possible to obtain the Doppler profile and thus the velocity distribution.

In principle a laser independent from the Zeeman slower light and resonant with a narrow line transition would be necessary for this kind of measurements. Unfortunately we did not have the 626 nm laser available at the time of the measurements so we tried to use the 421 nm light. We locked the laser to the atomic transition with the proper shifting for the Zeeman slower, and we use the AOM to scan the frequency of the probe beam. In principle the double pass AOM should allow a scan of 200 MHz, however changing the AOM frequency changes also its pointing affecting the optical fiber coupling. For this reason the scan is limited to 90 MHz. Even for this small scan we are forced to normalize the signal in function of the varying power output of the probe fiber during the scan. Another problem arising from the use of the 421 nm transition is the background light from the fluorescence induced by the Zeeman slower beam. This background is also dependent on the residual magnetic field in the detection zone. Also the small range of scanning, 90 MHz compared to the Doppler width of the velocity distribution $\sim 1 \,\mathrm{GHZ}$, forces us to observe the atoms almost perpendicularly to compress the velocity distribution in a small range of frequencies. In this configuration we observe the Doppler shift caused by the projection of the atomic beam velocity over the probe laser direction $v_z \sin \alpha$ where $\alpha = (4.7 \pm 1.0)^\circ$ is the angle between the probe direction and the direction perpendicular to the atomic beam. Unfortunately the natural linewidth remains the same in this configuration thus it masks the details of the resulting speed distribution.

For all these reasons this measurement has to be taken as a qualitative demonstration of the correct working of the Zeeman slower value of the parameters close to the expected ones. We performed different measurements in function of the temperature of the oven, the current in the Zeeman slower and the detuning and focusing of the laser light. As expected a variation in the laser detuning can be mostly compensated by a change in the current. Increasing the power over 200 mW is apparently unaffecting the cooling efficiency while for power lower than 100 mW we observe a clear reduction in the slowing efficiency. As an example in figure 3.34 it is possible to observe that increasing the current in the coils the fluorescence peak moves towards



Figure 3.34: Spectroscopy of the 421 nm transition performed under an angle of $(4.7 \pm 1.0)^{\circ}$ for different currents in the Zeeman slower coils: black I = 0; yellow I = 1.08A; orange I = 1.45A; red I = 1.73A; blue I = 1.94A. $T_{oven} = 1250^{\circ}C$; $P_{zs} = 110 \text{ mW}$; $I_{probe} = 30 \text{ mW/cm}^2$; $\delta = 31\Gamma$

smaller velocities leaving a small tail at high velocities. For too high currents it is possible to see that the signal drops, because in this case the atoms are pushed back into the slower as expected from calculations.

3.3 Next steps towards quantum degeneracy

3.3.1 MOT project

The Magneto optical trap (MOT) whose working mechanism is described in section 2.3.3 is currently under development in Pisa.

The 421 nm transition is not suitable to cool atoms down to the typical capture temperature of an ODT (see section 3.3.2). We therefore decided to use the 626 nm transition because it allows to reach lower temperatures due to the longer lifetime of the excited state.

To realize the 626 nm light we use a commercial Toptica laser TA-SHG Pro. It is composed by a external cavity diode laser (ECDL) with wavelength of 1252 nm amplified in a tapered amplifier (TA) [31] and then frequency doubled in a bow tie cavity locked to the laser frequency using the Hansch-Couillaud method. The SHG part is analogous to the one we realized to generate the 421 nm light described in section 3.1. The optical scheme of the laser is sketched in figure 3.35

This laser has a power of 800 mW and the declared linewidth is below 20 kHz thus narrower than the transition linewidth of $2\pi \cdot 136$ kHz. Therefore we have to reduce this linewidth in addition to lock the laser frequency to



Figure 3.35: Optical scheme of the Toptica 626 nm laser.

the atomic transition. To lock the laser to the atomic transition we will use a lock-in technique on a modulated saturation fluorescence spectroscopy signal Spectroscopy is performed on the same atomic beam we use for the experiment, that is the reason why we designed a custom flange between the transverse cooling cell and the differential pumping stage. On this flange we provide a 3-ways CF16 access to perform spectroscopy. The spectroscopy is placed 15 cm away from the beginning of the Zeeman slower in order to let the atoms return to the ground state before entering the slower. In fact the typical decay length is $l = \tau v \simeq 0.6$ mm, where τ is the lifetime of the transition and v is the atomic mean velocity.

We had not yet tested the cut frequency of the locking system, however we expect to obtain a higher cut frequency with respect to the one measured for the 421 nm lock. In fact in this setup we control the frequency of the laser modulating the current of the laser diode (ad the grating orientation for a slower but greater tunability), which is clearly faster than the piezo-actuated mirror inside the Ti:S laser.

Using the method described in section 2.3.3 we calculate that the capture velocity for a MOT with beam of 4 cm diameter is:

$$v_c = \sqrt{\frac{\hbar k \Gamma D}{m}} = 11.6 \,\mathrm{m/s} \tag{3.27}$$

A rough estimation of the power needed can be extracted considering 3 retroreflected beams with a diameter of 4 cm and with average intensity $(I_{\text{avg}} = P/\pi w^2)$ equal to the saturation intensity, this argumentation leads to a total power of P = 2.7 mW. Of course to let the MOT work properly we have to stay far above the saturation intensity, let us say that $100I_{\text{sat}}$ would be enough in any situation, so the power required on the MOT is approximately 300 mW. As usual we consider 50 % losses in the optical system and fibers so we require 600 mW out of the laser.

However the capture speed is barely higher than the final velocity of the Zeeman slower so the MOT may not capture all the atoms exiting from the slower. Lowering the Zeeman slower output speed further will increase the atomic beam divergence at the end of the slower so the only solution is to increase the capture velocity of the MOT. To do so it is possible to artificially broaden the transition linewidth using an AOM and modulating the drive frequency fast enough to let the atoms average over the modulation.

Clearly the MOT parameters can not be optimized at the same time to have high capture speed together with low temperature and high densities so we plan to change the MOT parameters during the experiment cycle. Larger intensities and broadening of the laser light together with a high magnetic field gradient enhance the capture phase of the MOT. This can understood considering that the capture velocity is proportional to $\sqrt{\Gamma}$ for $I \gg I_{\text{sat}}$.

However as the density increases, also does the probability that a spontaneously emitted photon is reabsorbed by another atom in the MOT resulting in an effective repulsive force between atoms that prevents further increments of the density. When the density approaches this limit, it is useful to perform an optimized cooling stage, in which the the Zeeman slower light is turned off, the shutter of the atomic beam is closed, the intensity, detuning and magnetic field gradient of the MOT are reduced and the line broadening is turned off. to achieve low temperatures a narrow linewidth is, in fact, needed since $T_{\text{Dopp}} \propto \Gamma$ and to get higher densities is better to have low intensities. Under this conditions we can hopefully approach the Doppler temperature of $3.26 \,\mu\text{K}$. After this cooling stage a recompression of the MOT will be performed to reduce its size and increase its density to optimize the loading in the ODT.

3.3.2 The final system

The final part of the experiment will be modified as in figure 1.7. We are going to replace the 6 way cross with an octagonal cell to get the optical access needed for the MOT and the ODT. The vacuum pump will be placed on a side of the cell using a tee in order to keep it near the cell to increase its pumping efficiency leaving one of the optical side access free. Also we are going to add a tee in front of the slower to install a metallic mirror. The Zeeman slower light will enter the vacuum apparatus from below and the mirror will reflect it into the slower. In this way the Dy atoms that are not captured by the MOT will deposit on a reflective surface instead than on a transparent surface. We in fact already could note a coating of the final viewport in the test configuration that we want to avoid in the final one.

Finally we are going to add a glass cell where we will perform the final evaporative cooling and have the degenerate sample ready to be manipulated and to perform measurements.

3.3.3 Optical trap transport

Using a glass cell displaced from the MOT region surely has many advantages, however this setup requires to transport the ultracold atomic cloud from one cell to the other. Magnetic transport can be performed by trapping the atoms in a magnetic trap and then moving the trap coils from one cell to the other. This method in our case has two overwhelming drawbacks: in first place we decided to use the glass cell to obtain a wide optical access but the mechanism to move two electric coil is very bulky. Also we would like to perform a fine magnetic field control so the less magnetic fields and magnetizable objects are present the better it is for our experiment. For these reasons we decided to use a movable ODT. We use an innovative focus-tunable lens that allows to realize this setup in a very compact way [32]. In fact using the optical scheme in figure 3.36 (a) where the first fixed lens has a focal length $f_1 \simeq f_{\min}$ where f_{\min} is the minimum focal length of the tunable lens. The tunable lens is placed at a distance $2f_1$ from the first lens and a third lens of focal f_3 is placed at a distance f_3 from the tunable lens. In this way the focal point of the beam can be moved from a distance f_3 from the last lens (when the tunable lens is set to its minimum focal length), ahead.

We tested this setup using an Optotune EL-10-30 tunable lens, which is basically composed by a bag or highly refractive liquid kept in position by an electrically actuated ring. When the ring is pressed against the bag the center of the bag bulges, its curvature increases thus the focal length decreases. This lens can tune its focal length between 40 mm and 120 mm, using the appropriate optical system we verified that it is possible to move a focal point with 50 μm waist over a range of more than 30 cm without any appreciable trembling in any position.

Furthermore if control on the waist is required it is possible to replace the first fixed lens with another tunable lens as illustrated in figure 3.36(b).



Figure 3.36: Optical scheme of the transport system. (a) With one tunable lens it is possible to control only the focus position while (b) using two tunable lenses it is possible to control both the position and the waist of the ODT laser beam.[32]

Chapter 4

Conclusions

The goal of my thesis work was to contribute to the design and development of a dipolar quantum gas experiment. During the two years I dedicated to this experiment I worked in a research team taking an active part in the laboratory life and in the studies and calculations that lie behind the construction of such an experiment.

I started my thesis work in June 2013 with an empty optical bench in the laboratory, since then I worked on the project of the 421 nm laser light. We decided that a home made SHG cavity would have been the best solution to produce a high power laser light in the near UV region with a narrow linewidth. Our SHG system indeed generates more than 1.5 W with a conversion efficiency above 60%. We also designed and realized a very stable locking setup that has been proved to lock below 0.3Γ of the transition of interest. In the second year of my thesis we realized the UHV setup and performed first characterization measurements. At the end of the experimental work for my thesis we obtained signatures of the proper working of the Zeeman slower stage.

The experiment recently moved to the CNR research area where the locking system for the 626 nm laser is almost completed and the MOT will be realized soon.

I had an important part in the design and realization of all the electronics for the detection and the stabilization of the lasers and the SHG cavity. I designed and aligned the optical setup for the 421 nm light, I mounted the UHV system and performed the baking. I actively took part in the measurements. These activities improved my experimental skills by a lot. I also personally took care of some orders and also took part in a poster session presenting my work.

I faced the everyday problems that happen in a laboratory such as the power drop of the pump laser for the 421 nm light generation, the room temperature control failure or the unexpected coating of the transverse cooling windows, and I personally handled some of them. Unfortunately it is impossible to complete the construction of such an experiment within a thesis work, so I do not have any "frontier physics" measurements to show to the reader. However the amount of work that hides behind the design and the testing of the experiment involves many fields of the applied and theoretical physics I learned in the university courses, beginning from the study of quantum dipolar interaction to understand the requirements of the experiment, passing through light-matter interaction and non linear optics, to conclude with electronics, optics and vacuum theory that have been fundamental to realize the different parts of this experiment.

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