REALIZATION OF A DIPOLE TRAP
FOR ULTRACOLD ATOMS OF RUBIDIUM
AND STUDY OF THE COLLECTIVE EFFECTS

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Chapter 1

Introduction

1.1 Collective Atomic Recoil Laser (CARL)

Collective atomic recoil lasing (CARL) is a process in which atoms interacting with a strong optical pump field spontaneously self-organize to form a highly-ordered density grating with a periodicity of half of the optical wavelength. The atoms organized in such compact bunches amplify coherently and exponentially an initially small optical probe field counterpropagating with the strong pump beam. CARL works as a laser in the sense that an initially small field, originated from fluctuations or spontaneously emitted photons, grows exponentially to reach a saturation value. However, it differs from a laser in many aspects: it has no threshold and it does not normally reach a stationary state. The process of CARL was first predicted in 1994 by R. Bonifacio and coworkers [1, 2, 3]. It may be regarded as the atomic analogue to the self-amplified spontaneous emission (SASE) process, which is at the basis of the free electron laser (FEL). Several attempts have been undertaken to see this effect in experiment with hot atomic vapors [4, 5]. However the first conclusive experimental observations of CARL, using cold atoms in a high finesse cavity, have been obtained in 2003 by the group of Ph. Courteille in Tübingen [6]. These landmark experiments have shown that cold atoms provide a unique, perfectly controlled environment to investigate nonlinear physics at the borderline between classical and quantum behaviour. The classical CARL process has also been studied in the presence of a friction force due to optical molasses fields [7, 8]. Another important aspect of CARL is related to its quantum regime realized with atomic Bose-Einstein condensates (BEC) (the first experiment has been reported by the group of W. Ketterle [9, 10, 11]). Whereas in the original CARL experiments the atoms, although cold, could still be described as classical particles, for atoms colder than the recoil limit, CARL exhibits a wave behaviour, in which the atoms are not described as classical particles but as de-localised quantum-mechanical waves [12, 13]. Moreover, the momentum exchange
between the light and the atoms, which is essentially continuous in a classical regime, becomes discrete in ultra-cold gases when the recoil frequency is larger than the CARL gain bandwidth. The theoretical analysis of the quantum regime has shown that the atom-photon system exhibits a robust entanglement, interesting for quantum information applications. Although CARL has originally been predicted for atomic ensembles at room temperature, collective gain could until now only be observed in ring cavities or with ultracold atoms.

Our current interest of CARL is then the possibility to observe spontaneous self-organization in a cloud of cold atoms (not condensed) in free space, i.e. without the boundary conditions imposed by external cavities. Large and dense clouds of cold atoms should allow for self-ordering of the atoms even without external cavity. If a cloud of cold atoms with a large optical thickness and high spatial density is used, interference of rescattered photons with the incident beam will create an interference pattern modifying the local intensity, similar to the situation in optical binding experiments. In the absence of an optical cavity, the radiation escapes rapidly from the atomic sample and a superradiant description is needed. This regime can be described using a “mean field” model (as in the regime with the cavity), in which a phenomenological damping term inversely proportional to the length of the atomic sample is added to the field equation. However a rigorous description of this superradiant regime requires the solution of the exact propagation equation of the radiation field [14].

1.2 CARL at INLN

The cold atom group at INLN (Institut Non Linéaire de Nice) has been developing techniques to produce large clouds of cold atoms [15, 16] with the goal of studying Anderson localization of light (together with other mechanisms to trap photons in a sample of atoms) and random lasers. These phenomena are based on multiple scattering of light by cold atoms and can be studied by analyzing the properties of the photons after their interactions with the atoms. At present, there are four cold atom experiments in Nice, two of them are dedicated to multiple scattering of light and its mechanical impacts on the atoms. The mechanical counterpart on the atoms of multiple scattering of light can be described by a long range force similar to the Coulomb interaction between charged particles. The optical forces induced by multiple scattering of light also limit the densities one can realize in a magneto-optical trap (MOT). A detailed understanding of these phenomena might allow for efficient spatial compression of the atomic cloud, important for Anderson localization as well as for Bose-Einstein condensation. For all these goals many atoms and high density are required and collective effects are expected; superadiance is one of these: it is in this contest that the study of classical CARL
without cavity is performed. The basic idea is to prepare a sample of cold atoms trapped in an optical dipole trap with high density and to illuminate it with a powerful beam detuned from the atomic resonance. The (collective) mechanical effects are investigated either measuring the position and the velocity of the atoms after the interaction or by detecting the light scattered back by the atoms.

This work is organized as follow: in section 2 we give a small theoretical background of the mechanical effects of the radiation on the atoms and in particular of the CARL process, in section 3 we describe the experimental setup and in section 4 we present our experimental observations.
Chapter 2

Mechanical effects of the radiation on the atoms

2.1 Radiation forces

Other than energy, an atom interacting with electromagnetic radiation exchanges also momentum, and is hence submitted to a force affecting the dynamics of its center of mass. This force can be used to stop, cool and trap atoms. The momentum exchange between atom and radiation originates, in a quantum framework, from emission and absorption of photons with momentum \(\hbar \vec{k}\). In the classical description, the radiation electric field induces on the atom an electric dipole moment. Hence, the Hamiltonian of an atom interacting with a radiation electric field is [17]:

\[
H(\vec{r}_e, \vec{p}_e, \vec{R}, \vec{P}, t) = \frac{p_e^2}{2m} + H_0(\vec{r}_e, \vec{p}_e) - \vec{d} \cdot \vec{E}(\vec{R}, t),
\]

(2.1)

where \((\vec{r}_e, \vec{p}_e)\) are electron variables (we consider hydrogen like atoms, with a single electron in the external shell), \((\vec{R}, \vec{P})\) are center-of-mass variables and \(\vec{d} = e \vec{r}_e\) is the dipole moment of the atom. The internal dynamics of the atoms is described by \(H_0\), with eigenstates \(H_0 \psi_n = E_n \psi_n\). In the dipole approximation the electric field \(E\) is uniform on the atom size (with wavelength \(\lambda \gg a_0\), Bohr radius) and depends on the center of mass coordinate \(\vec{R}\) only.

Since in quantum mechanics \((\vec{R}, \vec{P})\) are quantum operators, with \([R_j, P_k] = i\hbar \delta_{jk}\), their Heisenberg equations are:

\[
\frac{d\vec{R}}{dt} = \frac{1}{i\hbar} [\vec{R}, H] = \frac{\vec{p}}{m} \quad (2.2)
\]

\[
\frac{d\vec{P}}{dt} = \frac{1}{i\hbar} [\vec{P}, H] = -\nabla_{\vec{R}} H = \nabla_{\vec{R}} (\vec{d} \cdot \vec{E}) \quad (2.3)
\]

where \(\nabla_{\vec{R}}\) is the gradient with respect to the center of mass coordinate. Evaluating the average value on the atomic state \(\psi\), and defining \(\vec{x} = \langle \vec{R} \rangle_\psi\) and \(\vec{p} = \langle \vec{P} \rangle_\psi\),
the equations become

\[ \frac{d\vec{x}}{dt} = \frac{\vec{p}}{m} \]  
(2.4)

\[ \frac{d\vec{p}}{dt} = \langle \nabla_R (\vec{d} \cdot \vec{E}) \rangle_\psi. \]  
(2.5)

If \( \vec{E} = \hat{\epsilon}E \) has the polarization \( \hat{\epsilon} \), then \( \vec{d} = \vec{d} \cdot \hat{\epsilon} \). Since \( E \) does not depend on the internal variable \( \vec{r}_e \), it can be factorized out of the quantum average and Eq.(2.5) is approximated by:

\[ \frac{d\vec{p}}{dt} \approx \langle d \rangle_\psi \vec{\nabla} E = \vec{F}. \]  
(2.6)

Hence, the radiation force is proportional to the average atomic dipole and to the electric field gradient. If the field is a monochromatic plane wave, the force is directed along the wave vector \( \vec{k} \). Moreover, if the radiation intensity varies also, a force directed along the direction of variation occurs. For instance, a focused Gaussian laser beam produces strong transverse and a weak longitudinal forces which can be focusing or defocusing depending on the the atom response, i.e. on the average atomic dipole

\[ \langle d \rangle_\psi = d_{12}(\rho_{12} + \rho_{21}), \]  
(2.7)

where \( d_{12} \) is the dipole element of the transition (we consider only two-level atoms) and \( \rho_{12} \) is the coherence of the atomic density matrix.

### 2.1.1 Scattering force

For a plane wave, \( E = E_0 \cos(kx - \omega t) \), the radiation force is:

\[ F_x = -k d_{12}E_0(\rho_{12} + \rho_{21}) \sin(kx - \omega t) \]
\[ = -k\Omega(\rho_{12} + \rho_{21}) \sin(kx - \omega t), \]  
(2.8)

where \( \Omega = d_{12}E_0/\hbar \) is the Rabi frequency. The density matrix elements evolve following the equations:

\[ \dot{\rho}_{12} = -i(\omega - \omega_0 - kv_x)\rho_{12} - (\Gamma/2)\rho_{12} + (i/2)\Omega(\rho_{22} - \rho_{11}) \]  
(2.9)

\[ \dot{\rho}_{22} = -\dot{\rho}_{11} = -\Gamma\rho_{22} + (i/2)\Omega(\rho_{12} - \rho_{21}) \]  
(2.10)

where

\[ \rho_{12} = \rho_{12}e^{i(kx - \omega t)} \]  
(2.11)

(we consider a closed transition, with \( \rho_{11} + \rho_{22} = 1 \), with only radiative decay, i.e. \( \gamma_2 = \Gamma \) and \( \gamma_{12} = \Gamma/2 \)). When the relaxation time \( \Gamma^{-1} \) is much shorter than
the characteristic time of variation of the atomic momentum (typically \( \text{ns} \) against \( \mu \text{s} \)), then

\[
\tilde{\rho}_{12} \approx -\frac{i\Omega}{\Gamma} \frac{\rho_{11} - \rho_{22}}{1 + i(2\Delta/\Gamma)} \tag{2.12}
\]

\[
\rho_{22} \approx \frac{(\Omega/\Gamma)^2}{1 + (2\Delta/\Gamma)^2 + 2(\Omega/\Gamma)^2}. \tag{2.13}
\]

and

\[
\rho_{11} - \rho_{22} \approx \frac{1 + (2\Delta/\Gamma)^2}{1 + (2\Delta/\Gamma)^2 + 2(\Omega/\Gamma)^2}. \tag{2.14}
\]

where \( \Delta = \omega - \omega_0 - kv_x \). From (2.11) and (2.12) we obtain

\[
\rho_{12} + \rho_{21} = \frac{-2\Omega/\Gamma}{1 + (2\Delta/\Gamma)^2 + 2(\Omega/\Gamma)^2} \left\{ \frac{2\Delta}{\Gamma} \cos(kx - \omega t) + \sin(kx - \omega t) \right\} \tag{2.15}
\]

Inserting (2.15) in the expression (2.8) of the force and averaging in time, taking into account that

\[
\sin(kx - \omega t) \cos(kx - \omega t) = 0 \quad \text{and} \quad \sin^2(kx - \omega t) = 1/2,
\]

we obtain

\[
F_x = \frac{\hbar k}{\Gamma} \frac{\Omega^2}{1 + (2\Delta/\Gamma)^2 + 2(\Omega/\Gamma)^2}. \tag{2.16}
\]

Finally, defining the saturation intensity \( I_s \) such that:

\[
\frac{2\Omega^2}{\Gamma^2} = \frac{I}{I_s}, \tag{2.17}
\]

then the scattering force (2.16) assumes the form:

\[
F_{\text{scatt}} = \hbar k \frac{\Gamma}{2} \frac{I/I_s}{1 + (2\Delta/\Gamma)^2 + I/I_s}. \tag{2.18}
\]

The scattering force equals the rate at which the absorbed photons impart momentum to the atom, i.e. \( F_{\text{scatt}} = (\hbar k)R \), where \( R = \Gamma \rho_{22} \) is the scattering rate, which is maximum at resonance, \( \Delta = 0 \). The maximum force is for \( I \gg I_s \), \( F_{\text{max}} = \hbar k (\Gamma/2) \), when \( \rho_{22} \approx (1/2) \). It has the following interpretation in terms of photons: each absorbed photon gives the atom a kick in the direction of the incident photon and spontaneously-emitted photons result in momentum transfer in all directions. As a result, the average force after many cycles of absorption/spontaneous emission is directed along the propagation of the incident laser beam and can be used to slow down atoms. The scattering force is able to deflect or stop an atomic beams, producing an acceleration \( a_{\text{max}} = v_{\text{rec}}/2\tau \), where \( v_{\text{rec}} = \hbar k/m \) is the recoil velocity and \( \tau = \Gamma^{-1} \) is the lifetime of the excited state. Typically for alkali atoms \( v_{\text{rec}} \) is few mm/s (6mm/s for Rb) and \( \tau \) tens of ns, so \( a_{\text{max}} \sim 10^5 \) times the gravitational acceleration \( g \). The typical stopping distance
of an alkali atomic beam with a velocity \( v_0 \sim 1000 \text{ m/s} \) is \( L \sim 1 \text{ m} \). However, the atoms experiencing such a force are only those in a narrow range of velocities, \( \Delta v \sim \Gamma/k \), for atoms having a range of Doppler shift approximately equal to the homogeneous width of the transition, \( \Gamma \). Atoms that interact strongly with the laser light slow down until the change of their Doppler shift takes them out of the resonance with the light. It is possible to compensate this change in order to keep the force close to its maximum throughout the slowing process (The first realization of a MOT has been obtained using a slowed atomic beam, with either a so-called Zeeman slower or a chirped slowing technique [18]).

### 2.1.2 Dipole force

We have seen that the mechanism of scattering force relies on the absorption of photons incident on the atom. The analogous classical force is the radiation pressure force \( F = IA/c \), when light of intensity \( I \) is absorbed by a surface \( A \). In fact, the momentum carried by the light is \( E/c \) where \( E \) is the light energy and, since the power is \( P = dE/dt \), the light transfers momentum to the surface at rate \( d(E/c)/dt = IA/c \). However, classically it exists an other kind of radiation force not related to absorption, arising from the deflection of the light by a dispersive medium: for instance, when light passes through a prism of refraction index \( n \), it is deviated by an angle proportional to \( n \). The angular deviation can be seen as a change of the light momentum in direction and not in modulus, and induces a reaction force on the prism itself. More specifically, if the change of the light direction is \( \theta \), the change of the wave vector is \( \Delta k = 2k \sin(\theta/2) \) and the force caused by the light on the prism is \( (IA/c)2 \sin(\theta/2) \). This force is known as dipole (or gradient) force, and has interesting applications, for instance in biology (‘optical tweezers’). When a small dielectric sphere is illuminated by a focused laser beam, it diffracts light on opposite sides of the sphere with different strengths, proportional to the spatial varying intensity of the laser. This causes a net force that pull the sphere toward the region of high intensity, i.e. near the focus. With this method, it is possible to drag micro-organisms in water, as for instance biological cells, without perturbing them [18].

In order to derive the expression of the dipole force on a two-level atom, we return to the radiation force (2.6), for a plane wave in which the electric field amplitude varies also on \( x \). Then, Eq.(2.8) becomes more generally:

\[
F_x = (\rho_{12} + \rho_{21}) \left[ -\left( \hbar k \right) \Omega \sin(kx - \omega t) + \hbar \frac{\partial \Omega}{\partial x} \cos(kx - \omega t) \right] = F_{\text{scatt}} + F_{\text{dip}}.
\]

Combined with the atomic dipole (2.15) and after the average in time, the dipole
force is:

\[
F_{\text{dip}} = - \frac{1}{\Gamma^2} \frac{\partial^2 \Omega^2}{\partial x^2} \frac{\hbar \Delta}{1 + (2\Delta/\Gamma)^2 + 2(\Omega/\Gamma)^2} + \frac{1}{2I_s} \frac{\partial I}{\partial x} \frac{\hbar \Delta}{1 + (2\Delta/\Gamma)^2 + I/I_s}
\]  

(2.20)

The dipole force is proportional to the intensity gradient and is zero on resonance. For \( \Delta \gg \Gamma \) (and an intensity such that \( \Delta \gg \Omega \)), the dipole force is:

\[
F_{\text{dip}} \approx - \frac{\hbar \Omega^2}{4\Delta} \left( \frac{\partial}{\partial x} \frac{\hbar \Omega^2}{4\Delta} \right)
\]

(2.21)

Hence, the force derives from a potential and more generally, in three dimensions

\[
\vec{F}_{\text{dip}} = - \vec{\nabla} U_{\text{dip}},
\]

(2.22)

where

\[
U_{\text{dip}} \approx \frac{\hbar \Omega^2}{4\Delta} = \hbar \Gamma \frac{\Gamma}{8\Delta} \frac{I}{I_s}.
\]

(2.23)

### 2.2 Laser Cooling and Trapping

#### 2.2.1 Magneto-Optical Trap (MOT)

Whereas an atomic beam can be slowed with a single laser beam, a gas needs three orthogonal pairs of lasers to reduce its temperature (cooling). They use the Doppler effect to imbalance the forces of two counter-propagating laser beams with the same frequency.

Let’s consider initially a laser beam on an atom moving with a velocity \( v \) along the lasers direction. The detuning is \( \Delta = \Delta_0 \mp kv \) where \( \Delta_0 = \omega - \omega_0 \) and the minus (plus) sign holds for atoms moving in the same (opposite) direction of the laser beam. The force, for \( I \ll I_s \) and low velocities, \( kv \ll \Gamma \), is:

\[
F_{\text{scatt}} = \hbar k \Gamma \frac{I/I_s}{2 \left( 1 + (2\Delta_0/\Gamma)^2 (\Delta_0 \mp kv)^2 \right)} \approx F_0 \mp \alpha v
\]

(2.24)

where

\[
F_0 = \hbar k \Gamma \frac{I/I_s}{2 \left( 1 + (2\Delta_0/\Gamma)^2 \right)}
\]

and

\[
\alpha = 4 \hbar k^2 \frac{I}{\Gamma} \frac{I/I_s}{(1 + (2\Delta_0/\Gamma)^2)^2}
\]

(2.25)

The second term acts as a viscous force (proportional to the velocity) if \( \Delta_0 < 0 \) (red-shifted detuning). If on the atoms act two counter-propagating laser beams with the same frequency and intensity, the net force is

\[
F_{\text{molasses}} = (F_0 - \alpha v) - (F_0 + \alpha v) = -2\alpha v
\]

(2.26)
which is a friction force if $\alpha \propto -\Delta_0 > 0$. With three orthogonal pairs of lasers red-detuned from the atomic resonance it is possible to cool an atomic gas to very small temperature (hundreds or tens of $\mu K$)). In the reference frame of the atom moving toward the right, the Doppler effect leads to an increase in the frequency of the laser beam propagating in the direction opposite to the atom’s velocity. This Doppler-shift brings the light closer to resonance with the atom and thereby increasing the rate of absorption from this beam. This leads to a resultant force that slows the atom down. The light induces a frictional, or damping, force on the atoms just like that on a particle in a viscous fluid (like honey). For this analogy this way to cool atoms is called *Optical Molasses technique* [19].

The optical molasses technique accumulate cold atoms in the region where the three orthogonal laser beams intersect, because it takes a long time for atoms to diffuse out. With some change in the polarization of the laser beams and adding a magnetic field gradient, this configuration can be turned into a trap (Magneto-Optical Trap)[20]. The magnetic field is created by two pair of coils around the atoms, with current in opposite directions, producing a quadrupole magnetic field, which is zero at the center of the coils and whose magnitude increases linearly in every direction for small displacements from the zero point. The magnetic field does not confine atoms by itself, but causes an unbalance in the scattering forces of the laser beams and it is the radiation force which strongly confines the atoms, producing a variable Zeeman shift of the atomic hyperfine levels. For a $J = 0 \to J = 1$ transition, the constant magnetic field gradient splits the three sub-level with $M_J = 0, \pm 1$ of $J = 1$ with a separation depending on the atom’s position. The counter-propagating laser beams have circular polarization and red-shifted frequency with respect to the atomic resonance. The Zeeman shift causes an imbalance of the radiation forces. The frequency shift caused by the magnetic field can be incorporated in the detuning of the scattering force, $\Delta = \omega - kv - (\omega_0 \pm \beta z)$, so that for two laser beams with opposite circular polarization,

$$
F_{\text{MOT}} = F_{\text{scatt}}^+(\Delta_0 - kv - \beta z) - F_{\text{scatt}}^-(\Delta_0 + kv + \beta z) \\
\approx -2 \frac{\partial F}{\partial \Delta_0}(kv + \beta z) = -\alpha v - \frac{\alpha \beta}{k} z.
$$

(2.27)

We call $\beta = (g \mu_B / h)(dB/dz)$, where $\mu_B = 9.27 \times 10^{-24} J/T$ is the Bohr magneton, $h = h/2\pi = 6.64 \times 10^{-34} / 2\pi J s$ is the Planck constant and $g$ is the Landé factor. The unbalance of the radiation force caused by the Zeeman effect leads to a restoring force with spring constant $(\alpha \beta / k)$. Atoms that enter the region of intersection of the laser beams are slowed and the position-dependent force pushes the cold atoms to the trap center, providing an efficient and easy method to load a large number of cold atoms (up to $10^{10}$), to be used in laser cooling experiments.
2.2.2 Optical Dipole Trap

If we look at the expression (4.1) for the dipole potential we can see that when $\Delta$ is positive ($\omega > \omega_0$) this potential has a maximum where the intensity is highest, i.e. the atoms are repelled from regions of high intensity. But in the opposite case, of $\Delta$ negative ($\omega < \omega_0$), the dipole force acts in the direction of increasing $I$ and $U_{dip}$ is an attractive potential: atoms in a tightly-focused laser beam are attracted towards the region of highest intensity, both in the radial direction and along the axis of the beam. This dipole force confines the atoms at the focus of a laser beam to create an optical dipole trap. Let’s note that in this case the force is conservative, i.e. there is not a dissipation mechanism, like in the optical molasses technique, where the exceeding atomic momentum is carried away by the spontaneously-emitted photons. We note that the dipole potential is proportional to $I/\Delta$ whereas the radiation pressure is proportional to $I/\Delta^2$: increasing the detuning of the laser the dipole force became more and more important with respect to the radiation pressure.

It is thus possible to trap atoms also thanks to the dipole force illuminating, for example, a cloud of cold atoms with a gaussian beam. The intensity of a gaussian beam propagating along $z$ is:

$$I(x, y, z) = \frac{2P}{\pi \omega(z)^2} \exp\left(-\frac{x^2 + y^2}{\omega(z)^2}\right) \quad \text{with} \quad \omega(z) = \omega_0 \sqrt{1 + \left(\frac{z}{z_r}\right)^2}$$ \quad (2.28)

where $\omega_0$ is the beam waist, $z_r = \frac{\pi \omega_0^2}{\lambda}$ its Rayleigh length and $P$ its power. If we put the expression for $I$ (2.28) into (4.1) we get at the following expression for the dipole potential:

$$U_{dip}(x, y, z) = -\frac{U_0}{1 + \frac{z^2}{z_r^2}} \exp\left(-\frac{2x^2 + y^2}{\omega(z)^2}\right)$$ \quad (2.29)

where $U_0$ is the depth of the dipole trap:

$$U_0 = \frac{2P}{\pi \omega_0^2 \frac{1}{I_{sat}}} \frac{\hbar \Gamma}{8 \Delta}$$ \quad (2.30)

In an optical dipole trap the atoms must be gently set in the potential well and must have a sufficiently low kinetic energy, otherwise they escape from the potential. From this reason normally the atoms are transferred from a Magneto-Optical trap to a dipole trap and are not directly caught in the dipole trap from room temperature [25]. After a process of evaporation the temperature of the atoms reaches the value for which [21]:

$$U_0 = \eta k_B T$$ \quad (2.31)

and normally $\eta = 5 \div 10$. The atoms are thus set at the bottom of the dipole well and we can develop the expression for the potential:

$$U_{dip}(x, y, z) = U_0 \left(-1 + \frac{z^2}{z_r^2} + 2\frac{x^2 + y^2}{\omega(z)^2}\right)$$ \quad (2.32)
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This is the expression for an harmonic trap. The oscillation frequencies in the atomic position are:

\[ \omega_{x,y} = \sqrt{\frac{4U_0}{m\omega_0^2}} \]  

\[ \omega_z = \sqrt{\frac{2U_0}{mz_r^2}} \]  

2.3 Collective Atomic Recoil Laser

Let’s now turn to a closer look at the processes which lead to the CARL emission. In the collective atomic recoil laser the atoms, under the action of an intense laser pump, cooperate to amplify a counter-propagating probe beam, forming simultaneously an optical lattice which traps the atoms in their periodic wells.

The study of CARL have been developed in the case of presence of an external cavity which supports the emission. For this reason the derivation of the CARL’s equations is made in the mean field approximation. In the cavity in fact the field can be averaged over the atomic sample. Our current interest, however, is the possible realization of CARL without cavity. In this limit the CARL process is better described as superradiance and the mean field approximation is no more rigorously valid. As far as the CARL without cavity is concerned, we thus need to take into account the exact propagation of the field in the atomic sample as we will see below.

2.3.1 Optical lattices

The optical lattice is another another way of trapping atoms based on the dipole force. It is in fact possible to obtain a relatively strong gradient force in a standing wave formed by two counter-propagating laser beams: in this case the intensity changes from a maximum (at the anti-nodes) to zero (at the nodes) over a distances \( \lambda/2 \) to give a high gradient of intensity. Let’s consider two counter-propagating fields:

\[ E(x,t) = \frac{1}{2} \left[ E_{01} e^{ik(x-ct)} + E_{02} e^{-ik(x+ct)} + c.c. \right] \]  

creating two polarization waves

\[ \langle d \rangle = P(x,t) = \frac{d_{12}}{2} \left[ P_{01} e^{ik(x-ct)} + P_{02} e^{-ik(x+ct)} + c.c. \right] \]  

where

\[ P_{0i} = i \frac{\Omega_i/\Gamma}{1 - i(2\Delta/\Gamma)} = \frac{\Omega_i}{\Gamma} \frac{i + (2\Delta/\Gamma)}{1 + (2\Delta/\Gamma)^2} \quad (i = 1, 2). \]  

We have assumed \( \Omega << \Gamma \) and \( \rho_{11} \approx 1 \). The resulting force is

\[ F_x = P(x,t) \cdot \frac{\partial E(x,t)}{\partial x} = F_{scatt} + F_{dip}, \]
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where

\[ F_{\text{scatt}} = \frac{i\hbar k}{4} \{-P_{01}\Omega_{01}^* + P_{012}\Omega_{02}^* - \text{c.c.}\} \]

\[ = \frac{\hbar k}{2\Gamma} \frac{1}{1 + (2\Delta/\Gamma)^2} \left\{ |\Omega_{01}|^2 - |\Omega_{02}|^2 \right\} \]

(2.39)

and

\[ F_{\text{dip}} = \frac{i\hbar k}{4} \left\{ P_{01}\Omega_{02}^* e^{2ikx} + P_{02}\Omega_{01}^* e^{2ikx} - \text{c.c.} \right\} \]

\[ = \frac{2\hbar k}{\Gamma^2} \frac{\Delta}{1 + (2\Delta/\Gamma)^2} |\Omega_{01}\Omega_{02}| \sin(2kx) \]

(2.40)

The dipole force is zero on resonance \( \Delta = 0 \) and is maximum for \( \Delta = \Gamma/2 \). For \( \Delta \gg \Gamma \) it derives from the potential:

\[ U_{\text{dip}} = \hbar \frac{\Omega_{01}\Omega_{02}}{4\Delta} \cos(2kx) = \hbar \frac{\Omega_{01}\Omega_{02}}{2\Delta} \cos^2(kx) + \text{cost.} \]

(2.41)

We thus have that for a red detuned couple of beams a standing wave of light traps the atoms at the anti-nodes. This regular array of microscopic dipole traps is called optical lattice. The physical explanation of this dipole force is stimulated scattering of radiation. In a standing wave, an atom absorbs light with wave vector \( \vec{k} \) from one beam and the laser beam in the opposite direction stimulates the emission with wave vector \( \vec{k}' = -\vec{k} \); this gives the atoms an impulse of \( 2\vec{k} \).

We have here considered two counter propagating beams with the same frequency. However we can generalize what we have seen in the case of two beams having a small phase difference \( \phi = (\omega_2 - \omega_1)t \). It is enough to add \( \phi \) in the sinus argument in eq. (2.40). We have that in this case the optical lattice moves at the velocity \( v = (\omega_2 - \omega_1)/(k_1 + k_2) \).

2.3.2 CARL’s equations

In CARL the two beams which make up the optical lattice are a pump of frequency \( \omega_2 \) and Rabi frequency \( \Omega_2 \), constant, real and intense, and a ‘probe’ with frequency \( \omega_1 \) and complex Rabi frequency \( \Omega_1 \), variable in strength and in phase. The probe field, originated from spontaneously emitted photons or seeded by a prebunching beam, is fed by the pump photons back-scattered by the atoms. The atoms interact through the dipole force with these counter-propagating beams and the force on them is given by Eq.(2.38). When \( \Delta \gg \Gamma \), the scattering force (2.39) can be neglected and the dipole force is:

\[ F \approx \frac{\hbar k}{2\Delta} |\Omega_1|\Omega_2 \sin(2kx + (\omega_2 - \omega_1)t + \phi) \]

(2.42)

If we define the atom’s phase \( \theta_j = 2kx_j \) for each atom with \( j = 1, \ldots N \), the atom’s momentum \( p_j = mv_j \) and the ‘pump-probe detuning’ \( \delta = \omega_2 - \omega_1 \), the motion’s
equations for the atoms are:

\[
\frac{d\theta_j}{dt} = \frac{2k}{m} p_j \tag{2.43}
\]

\[
\frac{dp_j}{dt} = \frac{\hbar k}{2\Delta} \Omega_2 |\Omega_1| \sin(\theta_j + \delta t + \phi) \tag{2.44}
\]

These motion equations are coupled with the equation for the complex scattered field:

\[
\frac{d\Omega_1}{dt} = i \frac{\omega d^2}{2 \epsilon_0 \hbar} \eta (P_1 + P_2 e^{-i(\theta + \delta t)}) - \kappa \Omega_1 \tag{2.45}
\]

where \(\Omega_1 = |\Omega_1| \exp(i\phi)\), \(P_{1,2}\) are the amplitudes of the polarization waves relative to the two fields and \(\kappa = cT/L_{cav}\) is cavity damping. Assuming \(\Delta >> \Gamma\), so that, from (2.37), \(P_1 \sim -\Omega_1/2\Delta\), and \(\Omega_2 \gg \Omega_1\), Eq.(2.45) becomes:

\[
\frac{d\Omega_1}{dt} \approx -\frac{i\Omega_2}{2\Delta} \omega_p^2 (e^{-i(\theta + \delta t)}) - \kappa \Omega_1 \tag{2.46}
\]

where \(\omega_p = \sqrt{\omega d^2 n/2 \epsilon_0 \hbar}\) is the plasma frequency. In Eq.(2.46) the average is on all the \(N\) atoms, i.e.

\[
\langle e^{-i\theta} \rangle = \frac{1}{N} \sum_{j=1}^{N} e^{-i\theta_j} \equiv b \tag{2.47}
\]

where \(b\) is the bunching, i.e. the coherence factor of the emission. At the beginning, the phases \(\theta_j\) are random and \(b \approx 0\), but when the atom’s phases become correlated, as it occurs in CARL, the bunching factor becomes near unity, enormously enhancing the emission process.

Equations (2.43),(2.44) and (2.46) can be set in a more compact form redefining the variables as follows. We introduce a dimensionless parameter \(\rho\) and we define the dimensionless variables \(\bar{t} = 2\omega_{rec} \rho t\) (where \(\omega_{rec} = 2 \hbar k^2 / m\) is the recoil frequency) and \(\bar{p}_j = p_j/(2\hbar k) = kv_j/\omega_{rec} \rho\). In this way the equations become:

\[
\frac{d\theta_j}{d\bar{t}} = \bar{p}_j \tag{2.48}
\]

\[
\frac{d\bar{p}_j}{d\bar{t}} = \frac{2C_1}{\rho^2} |\Omega_1| \sin(\theta_j + \bar{\delta} \bar{t} + \phi) \tag{2.49}
\]

\[
\frac{d\Omega_1}{d\bar{t}} = -\frac{iC_2}{\rho} (e^{-i(\theta + \bar{\delta} t)}) - \bar{\kappa} \Omega_1 \tag{2.50}
\]

where \(\bar{\delta} = \delta/2\omega_{rec} \rho\), \(\bar{\kappa} = \kappa/2\omega_{rec} \rho\), \(C_1 = \Omega_2/8\omega_{rec} \Delta\) and \(C_2 = \Omega_2 \omega_p^2 / 4\omega_{rec} \Delta\). Finally, we eliminate the coefficients re-defining the scattered field such that \(iC_1 \Omega_1 / \rho^2 \equiv A \exp(-i\bar{\delta} \bar{t})\) in (2.49) and \(i\rho \Omega_1 / C_2 \equiv A \exp(-i\delta t)\) in (2.50) and we obtain:

\[
\frac{d\theta_j}{d\bar{t}} = \bar{p}_j \tag{2.51}
\]

\[
\frac{d\bar{p}_j}{d\bar{t}} = -\left( A e^{i\theta} + A^* e^{-i\theta} \right) \tag{2.52}
\]

\[
\frac{dA}{d\bar{t}} = \langle e^{-i\theta} \rangle + i\bar{\delta} A - \bar{\kappa} A. \tag{2.53}
\]
Since $C_1C_2 = \rho^3$, the $\rho$ parameter is:

$$\rho = \frac{1}{2} \left( \frac{\Omega_2}{2\Delta} \right)^{2/3} \left( \frac{\omega_p}{\omega_{rec}} \right)^{2/3} \cdot (2.54)$$

The scaled scattering field amplitude is $|A| = |\Omega_1|/(\omega_p\sqrt{2\rho})$ (where $\omega_p = \sqrt{\omega d^2 n/2\epsilon_0 \hbar}$ is the plasma frequency, with $n = N/V$). Since $(\epsilon_0 E_1^2)V/\hbar \omega$ is the average number of photons in the volume $V$, $\langle N \rangle_{\text{photon}}$, then

$$|A|^2 = \frac{\langle N \rangle_{\text{photon}}}{\rho N}. \quad (2.55)$$

Hence, $\rho|A|^2$ can be interpreted as the average number of photons scattered per atom.

We note that eqs. (2.51)-(2.53) have the same form of the equations for the free electron laser (FEL). In CARL the scattered radiation energy is provided by the kinetic energy of the atoms, whereas the pump laser plays the role of a unlimited energy source providing the photons to be scattered by the atoms. In this approximation, we neglect the decrease of the pump energy due to the scattering by the atoms. In the case of negligible field losses one of the constant of motion that Eqs.(2.51)-(2.53) admit is:

$$\langle \vec{p} \rangle + |A|^2 = \text{constant}, \quad (2.56)$$

i.e. the scattered field intensity grows only when the average atomic momentum decreases. More specifically, the decrease of average momentum equals the average number of scattered photons times the two-photon recoil momentum, $2\hbar k$.

In figure (2.1) we plot the numerical solution of Eqs.(2.51)-(2.53) for the field amplitude and for the bunching in the case of good cavity ($\kappa = 0$). We note the oscillatory behaviour of the system and that a stationary state it is not reached. If there are not losses the peaks of the CARL emission maintain the same amplitude.

The amplification bandwidth of CARL

We can investigate the conditions which allow the CARL process to start. We consider the initial conditions of field $|A| = 0$ and bunching $b = 0$, and small perturbations of this steady solution of the equations. It is thus possible to linearize the equation system and find that $|A|$ and $b$ grow in time, in this linear regime, as $\exp(i\lambda t)$, where $\lambda$ is the complex root of the characteristic equation:

$$\lambda^2(\lambda - \delta - i\kappa) + 1 = 0. \quad (2.57)$$

Exponential growth occurs when the solution of Eq.(2.57) is complex and with negative imaginary part. In the ’good-cavity’ limit, $\kappa \ll 1$, the maximum growth
Chapter 2. Mechanical effects of the radiation on the atoms

Figure 2.1: Intensity and bunching in CARL, in the good-cavity regime $\bar{\kappa} = 0$.

is on resonance, $\delta = 0$, with $\lambda = (1 - i\sqrt{3})/2$. The intensity grows as $\exp(\sqrt{3}t) = \exp(Gt)$, where $G = 2\sqrt{3}\omega_{\text{rec}}\rho$ is the exponential gain.

Furthermore, the solution of (2.57) shows that $\lambda$ is imaginary only for $\bar{\delta} < 2$, i.e. for $(\omega_2 - \omega_1) < 4\omega_{\text{rec}}\rho$: this define the CARL bandwidth, showing that only a limited range of frequencies around the pump frequency $\omega_2$ are amplified. With an initial distribution of the atomic velocity the detuning is modified as $\delta = (\omega_2 + kv) - (\omega_1 - kv) = \omega_2 - \omega_1 + 2kv$; the scattered intensity can thus be amplified only if the CARL bandwidth is larger than the initial Doppler broadening, i.e. for $2k\sigma_v < 2\omega_{\text{rec}}\rho$. Since $\sigma_v = \sqrt{k_B T/m}$, a maximum limit to the temperature of the atomic gas to experience the CARL amplification is set. Since typically $\omega_{\text{rec}} \sim (2\pi)10$kHz and $\rho \sim 10^3$, the temperature for CARL must be normally below hundreds of $\mu$K, i.e. the atoms must be very cold. For this reason, CARL has not been clearly observed experimentally until when a dense sample of cold atoms, as the one provided by a MOT, became available.

2.3.3 The superradiant regime of CARL

An interesting collective effect in CARL is superradiance, a phenomenon which shares some similarity with the better known superradiance or superfluorescence observed in two-level atoms. It occurs when the cavity radiation damping $\kappa$ is larger than the gain rate $G_{SR}$, or equivalently when the time-of-life of the photon in the cavity $\tau_c = 1/\kappa$ is shorter than the build time of the superradiant signal $\tau_{SR} = G_{SR}^{-1}$. In this limit the radiation amplitude follows adiabatically the time
evolution of the atoms and Eq.(2.53) is approximated by:

\[ A \approx \frac{1}{\bar{\kappa} - i\delta} e^{-i\theta} \]  

(2.58)
i.e. the radiation amplitude is proportional to the atomic bunching \( b = \langle \exp(-i\theta) \rangle \).

By substituting (2.58) in Eq.(2.52) and averaging, we obtain:

\[ \frac{d\langle \bar{p} \rangle}{dt} \approx -\frac{2\bar{\kappa}}{\bar{\kappa}^2 + \delta^2} |b|^2 < 0, \]  

(2.59)
the modulus of the average momentum is continuous increasing in time: the atoms scatter the pump photons into the reverse mode, whose photons are in average not scattered back to the pump. This is typical for superradiance, in which the photons are only 'emitted' and the reabsorption is inhibited by the fast escape of the light from the atomic sample (i.e. large \( \kappa \)). The maximum of emission occurs at resonance (\( \delta = 0 \)) and with a large bandwidth, approximately equal to \( \bar{\kappa} \).

For \( \delta = 0 \), the maximum radiation intensity is \( |A|_{\text{max}}^2 \sim 1/\bar{\kappa}^2 \), so that the maximum number of photons (from eq. (2.55)) is \( N_{\text{photon}} \approx (2\omega_{\text{rec}}/\kappa)^2 N \rho^3 = (\Omega^2/\Delta)^2(\omega_p/\kappa)^2(N/8) \propto N^2 \): the radiated intensity is proportional to the square of the atom number, indicating a cooperative emission of the atoms, which at the moment of the maximum intensity scatter the photons in phase.

The superradiant gain \( G_{SR} \) can be calculated as before from a stability analysis. The adiabatic elimination consists in assuming \( \bar{\kappa} \gg \lambda \) in the dispersion equation (2.57). At resonance, the gain is \( G_{SR} = (2\omega_{\text{rec}}\rho)\sqrt{2/\bar{\kappa}} = (\Omega^2/2\Delta)\omega_p\sqrt{2\omega_{\text{rec}}/\kappa} \propto \sqrt{N} \), smaller than in the 'good-cavity' limit by a factor \( \sqrt{2}/3\bar{\kappa} \). We note the different dependence on \( N \) of the gain in the superradiant regime (\( \sqrt{N} \)) and in the good-cavity regime (\( N^{1/3} \)).
Figure 2.3: **Superradiance in CARL in the free space (self-similar solution).**

In figure (2.2) the numerical solutions of the CARL’s equations in the bad cavity case ($\bar{\kappa} = 1$) for the field amplitude and for the modulus of the average velocity of the atoms are shown. We note important differences between this superradiant case ad the case of CARL in good cavity. The peaks of emission which follow the first one are in the superradiant limit more and more smaller and the atomic modulus of the average velocity does not oscillate but it increases along all the process.

Eqs.(2.51)-(2.53) are in general valid only in an optical cavity, in which the field can be averaged over the atomic sample and the transmission of the cavity described by a simple damping term in the field amplitude equation. However, to describe more accurately the process in free space, the exact propagation of the field in the atomic sample has to be taken into account and the equations are modified:

\[
\begin{align*}
\frac{\partial \theta_j}{\partial t'} &= \bar{\rho}_j \\
\frac{\partial \bar{\rho}_j}{\partial t'} &= -\left( A e^{i\theta_j} + A^* e^{-i\theta_j} \right) \\
\frac{\partial A}{\partial z'} &= \langle e^{-i\theta} \rangle,
\end{align*}
\]

where $t' = 2\omega_{rec}\rho(t - z/c)$ and $z' = 2\omega_{rec}\rho(z/c)$ are retarded time and position in the atomic sample. It exists the following not trivial self-similar solution of Eqs.(2.60)-(2.62) describing superradiance: $A(z', t') = z'A_1(y)$, $\theta_j(z', t') = \theta_1j(y)$ and $\bar{\rho}_j(z', t') = \sqrt{z'}p_{1j}(y)$, where $y = \sqrt{z'}t'$ and $A_1$, $\theta_1j$ and $p_{1j}$ are solutions of
the following equations:

\[
\frac{d\theta_{ij}}{dy} = \bar{p}_{ij} \tag{2.63}
\]

\[
\frac{dp_{ij}}{dy} = -\left( A_{1}e^{i\theta_{ij}} + A_{1}^{*}e^{-i\theta_{ij}} \right) \tag{2.64}
\]

\[
\frac{y}{2} \frac{dA_{1}}{dy} + A_{1} = \langle e^{-i\theta_{1}} \rangle \tag{2.65}
\]

The previous description of superradiance in a cavity is recovered approximating \( z' \) with \( 1/\bar{\kappa} \) and neglecting the term \( (y/2)(dA_{1}/dy) \) in (2.65). However, this term is responsible for further secondary peaks after the principal peak, as can be observed in the figure (2.3). After the first peak, atoms are dephasing and reabsorb radiation, to emit later further smaller superradiant peaks. The modulus of the average velocity of the atoms too, does not continue to increase in time but has some peaks. This effect is named ’ringing’ in the two-level superfluorescence.
Chapter 3

Experimental setup

3.1 MOT

The biggest part of the experimental setup is devoted to realize the magneto-optical trap (MOT) i.e. to cool the Rb atoms and trap them in the center of a vacuum cell. In the figure (3.1) the atomic levels for the two isotopes $^{85}$Rb and $^{87}$Rb are reported. Cooling and trapping is performed on the D2 line of the Rubidium $^{85}$Rb ($5S_{1/2} \rightarrow 5P_{3/2}$) at the wavelength $\lambda = 780\text{nm}$. We call the beam used to cool the atoms MOT beam. The light of the MOT laser is red-detuned (2 or 3 $\Gamma$) from the hyperfine transition $F = 3 \rightarrow F' = 4$, which is a closed transition, but there is a probability different from zero to excite the open transition $F = 3 \rightarrow F' = 3$, and so it is possible to have a change of the hyperfine ground state via a spontaneous Raman transition. This happens for each $10^3$ photons exchanged between one atom and the MOT laser. The atoms which fall to the F=2 state do not see anymore the MOT light and so they are lost. For this reason we need a repump laser tuned, in our case, to the $F = 2 \rightarrow F = 3'$ transition of the D2 line. It pumps back the atoms into the F=3 hyperfine ground state and so they are ready to be cooled and trapped again.

The setup for the two beams is similar: they are obtained from two DFBs (Distributed FeedBack laser) and they are frequency-stabilized on a cross-over transition of the saturated absorption. Then they are frequency-shifted on the blue side of the respective resonances via a double pass Acousto-Optic Modulator (AOM). Other two AOMs, in single pass configuration, are used to have the repumper in resonance with its transition and the MOT laser slightly red detuned. All the AOMs are controlled by independent Voltage Controlled Oscillators (VCO). We can easily set the frequency of the two beams changing the frequency of the double pass AOM, and the power changing the voltage of the single pass AOM. The single pass AOMs are also used as switches, i.e. to turn on and off the laser light.
Figure 3.1: Atomic energetic levels for $^{85}\text{Rb}$ and $^{87}\text{Rb}$.
The power obtained from the DFB is not enough to trap the atoms and so, before the single pass AOM, we injected the MOT beam into a MOPA (Master Oscillator/Power Amplifier). The output of the MOPA is a beam with the same wavelength but far more powerful (up to 1W). Using polarizing cubes, the light is then split into six beams with the same power, directed in counter-propagating pairs into the 10cm sized cubic vacuum chamber. The six beams have been each expanded by a telescope to a waist of about 2.5cm before entering the vacuum chamber. At the end the intensity for each of the six beams used to cool the atoms is about 3mW/cm$^2$.

For the repump beam the power obtained from the DFB is enough so we don’t need any amplification. After the single pass AOM it is split in two parts with the same power by a polarizing cube, it is expanded by the same telescope as the MOT beam and it is directed on the atoms. The intensity on the atoms is about 1mW/cm$^2$ for each beam.

The magnetic field gradient is obtained by two coils in an anti-Helmholtz configuration and it is typically of about 10Gauss/cm (using a current of 4.6A in the coils). The magnetic field can be rapidly switched on and off. In figure (3.2) the laser setup for cooling and trapping atoms is shown.

At INLN there are two similar experimental setups for cooling Rubidium atoms. All the experiments presented in this report have been performed on the new one. The first MOT with this setup has been realized in April 2008 and not all the parameters are yet optimized. At the moment we typically trap $10^8$ atoms at 100$\mu$K, the atomic cloud is spherical and its diameter is about 3mm.

### 3.2 Cloud preparation

As we will see below we need to transfer the atoms into an optical dipole trap. This procedure depends critically on initial temperature, size and density of the cloud [22]. The parameters of the atoms trapped in the MOT are not good for the transfer, so we have tried several protocols. The final sequence we have used for all the experiments discussed in this report uses three different steps: MOT, dark MOT and molasses.

The MOT size, and so its density, is limited by repulsive interaction between the atoms. This interaction is based on multiple scattering of photons. A well known technique [23] to reduce this effect is to pump the atoms into a different hyperfine level of Rubidium: this is done by reducing the intensity of the repump laser (to about 5% of the maximum power). After this dark MOT we add 30 ms of molasses phase where the gradient of the magnetic field is switched off and the power of the repumper is reset to the maximum. Laser cooling is more efficient in such a molasses than in a MOT. The parameters of a typical sequence are noted
Figure 3.2: Laser setup for cooling and trapping atoms
in table (3.1).

<table>
<thead>
<tr>
<th>Typical sequence</th>
<th>Time (ms)</th>
<th>$I_{MOT}$ (mW/cm$^2$)</th>
<th>$\delta_{MOT}$ (Γ)</th>
<th>$I_{rep}$ (mW/cm$^2$)</th>
<th>$\delta_{Rep}$ (Γ)</th>
</tr>
</thead>
<tbody>
<tr>
<td>MOT</td>
<td>1000</td>
<td>3</td>
<td>-3</td>
<td>1</td>
<td>$\approx 0$</td>
</tr>
<tr>
<td>Dark MOT</td>
<td>35</td>
<td>3</td>
<td>-3</td>
<td>0.05</td>
<td>$\approx 0$</td>
</tr>
<tr>
<td>Molasses</td>
<td>30</td>
<td>3</td>
<td>-2</td>
<td>1</td>
<td>$\approx 0$</td>
</tr>
</tbody>
</table>

Table 3.1: Typical sequence used to prepare the atomic cloud

After the molasses we switch on the dipole laser used for CARL. As we are interested in collective effects, one of the most important parameters is the number of atoms. We can control the final number of atoms trapped in the dipole trap by changing the duration of the loading time for the MOT.

We did not make a systematic study of all the parameters of the sequence which can surely be improved.

3.3 Dipole trap-CARL beam

We call CARL beam the laser beam that is used as pump in the CARL process. In figure (3.3) the setup for this laser beam is illustrated. It is obtained from a DFB laser, which allows to have a narrow spectral width (of the order of 5MHz). It’s possible to control the wavelength of the emitted light with the laser current and with the temperature of the diode: $\lambda = \lambda(I,T)$ increases with $T$ and $I$. In graph (3.4) we can see the calibration curves for the laser wavelength. If we keep the temperature of the laser fixed, we can choose the detuning of the CARL beam by changing the laser current. To know the absolute value of the laser wavelength, we need to know the current corresponding to the atomic resonance: for this reason we take a small part of the beam, pass it through a small Rubidium cell at room temperature and detect the transmitted beam. In this way, modulating the current of the laser, we make the absorption spectroscopy of the beam. When the frequency of the laser is close to the atomic resonance we can see on the oscilloscope the four peaks which correspond to the transitions from $F=2$ and $F=3$ for $^{85}$Rb, and from $F=1$, $F=2$ for $^{87}$Rb. Another small part of the beam enters a Fabry-Perrot cavity and we detect the transmission through this cavity on a photodiode. Comparing the transmission through the Fabry-Perrot cavity with the different transitions of Rubidium, we determinate the Free Spectral Range of the cavity (FSR=810MHz). Then we increase the current and, starting from the peak which corresponds to the D2 line of $^{85}$Rb (see image (3.5)), we count the peaks of the Fabry-Perrot in order to calibrate the curve $\delta$ vs $I$. 
The power that can be obtained from the DFB is not enough for our experiments and we thus need to amplify it preserving the same wavelength. For this reason we use a MOPA. The output of the MOPA is a beam with a power that can be up to 1W according to the power of the injected beam (which can be varied turning a \( \lambda/2 \) put in front of a polarizing beam splitter) and to the current of the power supply of the MOPA. In graphs (3.6) and (3.7) we can see the dependence of the output power on the input power and on the current.

In order to be able to rapidly switch on and off the beam we use an AOM controlled by a VCO. The first order of the AOM is our \textit{CARL beam}. The variation in the frequency induced by the AOM (80 MHz) can be neglected because we use the beam far detuned (tens of GHz) from resonance.

In order to focus the laser beam on the cloud of atoms, we put a lens at the focal distance from the MOT. We choose a lens with \( f=500\text{mm} \) in order to have a beam waist of 200\( \mu\text{m} \). The dimension of the beam has been measured outside the cell at the same distance as the atoms with a razor blade and with a CCD camera and the two values are in good agreement (see the figures (3.8) and (3.9)). We also tried to measure it via the fluorescence of the hot atoms but we didn’t get the same result. We obtained a value which is about the double of what we have measured outside the vacuum cell. This difference is not understood yet, but it might be due to saturation of the hot atoms.
Chapter 3. Experimental setup

Figure 3.4: Calibration curves for the detuning of the CARL beam. Different lines correspond to different temperatures. The slope of the curve is the same for all the temperatures and it is $d\delta/dI = 2\text{GHz}/\text{mA}$

Figure 3.5: We sweep the current of the DFB (and thus the frequency) with a ramp and we look, at the same time, at the four peaks of the absorption of the Rubidium cell (blue) and at the peaks of the Fabry-Perrot cavity (green).
Figure 3.6: Dependence of the output power of the MOPA on the injected power

Figure 3.7: Dependence of the output power of the MOPA on the current of the power supply of the MOPA
Figure 3.8: Measurement of the beam waist with the razor blade method: a straight razor blade is drawn across the beam profile and the intensity is recorded for several position of the blade (on the left). The derivative of this curve (on the right) is a gaussian beam profile. The value for the beam waist obtained from the fit is 190 µm.

Figure 3.9: Measurement of the beam waist with the CCD camera: on the left the picture of the beam; on the right the fits on the integration of the image along the two axis. The values for the beam waist obtained from the fit are reported.
In the path from the MOPA to the MOT we lose some power because of the AOM and the different mirrors and cubes used to bring the beam up to the atoms and to have the good polarization. In front of the vacuum cell, we add a polarization cube to filter the good polarization and a $\lambda/4$ plate to make it circular. This allows us to obtain a well controlled circular polarization of the laser incident on the atomic cloud. Also, light backscattered from the atoms will pass twice through the $\lambda/4$ plate and leave the polarization cube in the orthogonal channel, which allows the detection of a small amount of reflected light from the atoms.

At the end the beam which illuminates the atoms has a waist of 200$\mu$m, a power that can be tuned up to 200mW and a frequency which can be varied between the atomic resonance and a detuning of -200GHz.

The first part of my master thesis at INLN consisted of setting up this dipole laser.

### 3.4 Detection

#### 3.4.1 Taking and analyzing pictures

The most part of our data are extracted from images taken with a CCD camera (DragonFly2). We can easily set the gain and the shutter time of the camera (typically about 10ms). To take the image we switch on again the six beams of the MOT and the repumper (the detuning is typically a little different) and we look at the fluorescence of the atoms. The figure (3.10) is an example of the typical pictures we obtain.

The size of the picture in pixel is 1064 X 768 and the size of a pixel is 3.9$\mu$m. To image the MOT onto the CCD, we use a commercial objective, which allows us to obtain a sharp picture of the atoms. We calibrated the camera taking the picture of a ruler put at the same distance from the camera as the MOT. We thus obtain the calibration of the imaging which corresponds to 39$\mu$m per pixel at the MOT position. We also measured the resolution of the camera looking at the image of a razor blade put at the same distance from camera as the MOT. We found that we are not able to resolve objects which are smaller than 200$\mu$m.

The camera does not only detect the light scattered by the cold atoms but also the stray light from the beams and from the room, for this reason we always need to take a picture of the background (without the cold atoms). If we are only interested in the atoms trapped by the dipole laser we take a background picture by switching off the CARL laser.

We noted that interference fringes (possibly due to the CCD window) can affect the quality of our pictures. Also stray light (from dust on the vacuum cell)
limits the performance of our imaging technique.

Image analysis

We analyze the images by MATLAB. We subtract the background from the image and we select the part of the figure we are interested in. The program makes the sum of the pixels in the same row and in the same column, projecting the image along the two axis. It then makes a gaussian fit for the two axis and gives the values of the center and of the width of the two gaussians. From here we can extract many information. We report an example of a picture analyzed by this program in figure (3.11).

3.4.2 Position, velocity and temperature measurement

Position

The most important information we have extracted from the recorded figures is the displacement of the center of mass of the atomic cloud and its dependence on different parameters. This displacement is due to the interaction between the CARL beam and the atoms themselves. The initial position is the position of the MOT (center of the gaussian fit on the MOT). It’s necessary to pay attention not to change the region for the fit during one set of data.
Figure 3.11: Image analyzed by MATLAB: we extract information on the position and on the size of the atoms trapped by the dipole beam.
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Figure 3.12: Picture after different time of flight: we extract information on the velocity and on the temperature of the atoms trapped by the dipole beam. In the images a modulation in the fluorescence detected by the CCD is clearly present. This modulation is, almost partly, due to the presence of interference fringes (possibly due to the CCD window) that affect the quality of our pictures.

Velocity

We also make measurements of the average value of the final velocity of the atoms after the interaction with the CARL beam. We take different images for increasing time of flight (TOF) i.e. we switch off the CARL beam and we take the image after some milliseconds, where the atoms are no more subjected to any force except gravity and they thus keep their final velocity. We measure the positions of the center of mass of the atomic cloud and, from a linear fit (figure (3.13)), we get its velocity:

$$\langle x(t) \rangle = \langle x(0) \rangle + \langle v \rangle \text{TOF}$$ (3.1)

In the figure (3.12) we can see the images taken for different time of flight. The cloud falls because of gravity. Due to the finite temperature of the atoms, the cloud is also expanding after being released from the laser. Furthermore the non zero longitudinal velocity of the atoms results in a horizontal displacement of the center of mass of the cloud.
Figure 3.13: Position of the atoms as function of the TOF. The slope of the linear fit is the average longitudinal velocity

**Temperature**

The protocol for the measurement of the temperature of the cold atoms is very similar to the one for the velocity. We take the same kind of images but we look at the width of the gaussian instead of the position. If we plot the square of the gaussian width as function of the square of the TOF we obtain a linear dependence (figure (3.14)). From the linear coefficient of the fit we get the temperature of the atoms:

\[
T = \frac{a m_{\text{Rb}}}{k_B}
\]

where \(a\) is the linear coefficient and \(m_{\text{Rb}}\) is the atomic rubidium mass.

In theory the temperature along the two axis can be different. We always only measure the transverse one because the longitudinal dimension of the cloud is too large so that an additional expansion after 12ms of TOF is not detectable. The cloud is too dilute, which makes a quantitative observation more delicate.

### 3.4.3 Atoms number measurement

We are interested in collective effects: it is thus important to be able to know how many atoms are trapped by the dipole beam in order to study the atomic behaviour as a function of the number of atoms. If we know the number of atoms
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Figure 3.14: Transverse size of the atomic cloud as function of the TOF. From the slope of the linear fit we calculate the temperature.

in the MOT we can calibrate the fluorescence detected by the CCD camera from it. We always consider the integral of the gaussian fit of the cloud along the y axis as value for the fluorescence detected.

We calculated the number of atoms in the MOT from its density (measuring the optical thickness as explained below) and its size. Typical values for the MOT density and number of atoms are \( n = 10^{10} \) cm\(^{-3} \) and \( N = 10^8 \); in the dipole trap we usually have \( N = 10^5 - 10^6 \) atoms.

**Optical thickness**

It is possible to deduce the optical thickness from the measurement of the transmission of a probe laser through the sample of atoms:

\[
T = \frac{I_{\text{out}}}{I_{\text{in}}} = e^{-\frac{L}{\ell_{\text{sc}}}} = e^{-b(\delta)}, \tag{3.2}
\]

where \( I_{\text{in}} \) and \( I_{\text{out}} \) are the intensity of the probe beam respectively before and after the sample, \( b(\delta) \) is the optical thickness, \( \ell_{\text{sc}} \) is the scattering mean free path and \( L \) is the length of the sample.

We can derive the expression for \( b(\delta) \) in (3.2) from the steady state equation of the field in the Maxwell-Bloch equations:
Figure 3.15: Transmission of the probe beam as function of its detuning after the dark MOT phase. From the full width of the curve at half maximum ($\Delta \nu_{\text{FWHM}}$) we can extract the value of $b_0$. In this case $\Delta \nu_{\text{FWHM}} = 5.8$ and then $b_0 = 24$.

\[
\frac{\partial \Omega}{\partial z} = -\frac{\omega_p^2 T_2}{c} \frac{\Omega}{1 + \delta^2 T_2^2 + T_1 T_2 \Omega^2} \tag{3.3}
\]

where $\omega_p = \sqrt{\frac{n d \omega_0}{2 \hbar c}}$ is the plasma frequency, the $\delta = \omega - \omega_0$ is the detuning and $\Omega = \frac{4 E}{\hbar}$ is the Rabi frequency. Neglecting the collision ($T_1 = T_2 = \frac{1}{\Gamma}$), and introducing the saturation intensity $\frac{I}{I_{\text{sat}}} = T_1 T_2 \Omega^2 = \frac{2 \Omega^2}{\Gamma^2}$, the equation (3.3) becomes:

\[
\frac{\partial I}{\partial z} = -\alpha \frac{I}{1 + 4 \frac{\delta^2}{\Gamma^2} + \frac{I}{I_{\text{sat}}}}, \tag{3.4}
\]

where $\alpha = \frac{4 \omega_p^2}{\Gamma}$ is the absorption coefficient. Using the proper definition of $d$ ($d = \sqrt{\frac{3 \pi \sigma_0 h \hbar}{k^2}}$) in the expression for $\omega_p$, we have that $\alpha = n \sigma_0$ where $\sigma_0$ is the well known cross section at resonance:

\[
\sigma_0 = \frac{3 \lambda^2}{2 \pi} \tag{3.5}
\]

For $I \ll I_{\text{sat}}$ the equation (3.4) yields:

\[
I(L) = I(0) \exp \left(-\frac{\alpha}{1 + 4 \frac{\delta^2}{\Gamma^2} L} \right), \tag{3.6}
\]
Comparing the equations (3.2) and (3.6) we have that:

\[ b(\delta) = \frac{b_0}{1 + 4\frac{\Gamma^2}{\nu^2}} \]  

\[ \text{where} \quad b_0 = n\sigma_0 L. \]  

(3.7)

We measure the optical thickness of the MOT at resonance \( b_0 \), by a weak tunable probe laser obtained from the zero order of the single pass AOM for the MOT beam. We can modulate its detuning by a ramp on the VCO for the double pass AOM and thus vary the probe frequency around the atomic resonance. We detect the probe beam behind the cloud of atoms by a photodiode. We obtain a graph (see the figure (3.15)) of the transmission as function of the detuning of the probe and we extract the value of \( b_0 \) from the full width of the curve at half maximum (\( \Delta\nu_{FWHM} \)). In fact, from equations (3.2) and (3.7), considering \( T = \frac{1}{2} \), we get at:

\[ b_0 = \ln(2)(1 + \frac{\Delta\nu_{FWHM}^2}{\Gamma^2}). \]  

(3.8)

The previous equations are true only for a monochromatic laser source. The spectral width of our laser is 5MHz, we thus need to introduce a correction for the optical thickness. The figure (3.16) shows that the value for the optical thickness is overestimated if we do not take into account the laser spectral width.

Typically in our MOT, after the dark MOT phase, \( b_0 = 20 - 25 \). If we perform the measurement after the molasses phase we find a smaller value (\( b_0 = 8 - 10 \)) because the density of the atomic cloud decreases when the magnetic field is switched off.
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3.5 Characterization of the dipole trap

In order to better investigate the possible effects which determine the atomic behaviour we need to have some more precise information on the dipole trap and on the cloud of trapped atoms. In particular, because of the detection limit, we do not have a measurement of the beam waist inside the vacuum cell and we can not directly measure the transverse size of the cloud too, and thus we can not calculate the atomic density. However we can get to these information by indirect measurement approximating the dipole potential with an harmonic one.

Oscillation of the atoms in the dipole potential

As we have already seen in 2.2.2, the atoms in the dipole trap are situated at the bottom of the dipole well and we can develop the expression for the potential obtaining the equation for an harmonic trap. If we thus displace the atomic center of mass from the minimum of the potential the cloud begins to oscillate with the frequency:

$$\omega_y = \sqrt{\frac{4U_0}{m_{Rb}w_0^2}}$$  \hspace{1cm} (3.9)

We measure the period of the transverse oscillation along the vertical direction. Once the atoms are trapped in the dipole potential, we trigger the oscillations suddenly switching off the beam and switching it on again at off = 2ms later with a bigger power. During this time, the cloud is displaced by $\delta y = g t_{off}^2 = 20\mu m$ by the gravitation force. This amplitude of oscillation is too small to be directly detected. We thus measure the vertical oscillation of the cloud after a time of flight of 10ms, in this way it is as if we were measuring the oscillations in the mean transverse velocity. The period of the oscillation in position and in velocity in the dipole trap is obviously the same and the measurement of the second one is possible with our camera resolution. In the picture (3.17) the experimental data are reported.

From these data we extract a period of oscillation $T_{osc} \approx 9 ms$. In these measurements the power of the beam is $P=110mW$ and its detuning is $\delta=-80GHz$. With the equation (3.9) we calculate, considering a beam waist $w_0 = 200\mu m$, a vertical oscillation frequency $\omega_y/2\pi = 240Hz$ which corresponds to a period of 4.5ms. The measurement is in qualitative agreement with the expected value.

One has to note that if the atoms are in different magnetic sublevels, the average force acting on the atoms has to be estimated by taking into account the various Clebsch-Gordon coefficients. This tends to reduce the oscillation frequency compared to a 2 level atom. We can use this experiment as an indirect measurement of the beam waist of the dipole beam. We have, in fact, no way to get to this value with a direct measurement inside the vacuum cell.
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Figure 3.17: Position of the atoms t_{dpi2} ms after that the dipole beam is switched on again, and after 10ms of TOF.

Temperature, size and density

With the protocol described in 3.4.2 we can measure the temperature of the atoms trapped in the dipole potential. For a beam with \( P=110 \text{mW} \) and \( \delta=-80 \text{GHz} \) we measure a temperature \( T=23 \mu \text{K} \). The depth of dipole well for this beam, calculated with equation (4.1) is \( U_{\text{dip}} = 4 \times 10^{-27} \text{J} \) which corresponds (dividing by \( k_B \)) to a temperature \( T_{\text{dip}} = 280 \mu \text{K} \). We thus have that \( \eta = U_0/k_B T \) is bigger than 10 and it is thus true that the atoms are located at the bottom of the dipole well and the harmonic trap approximation is good.

As we have already told, we are not able to resolve objects which are smaller than 200 \( \mu \text{m} \) with the CCD camera. We thus cannot have a direct measurement of the transverse size of the cloud of atoms in the dipole trap. We can however get to this value from the measurement of the temperature and of the period of oscillation. In fact, if we always approximate the gaussian potential with an harmonic trap, we have that:

\[
\frac{1}{2} m_{\text{Rb}} \omega_y^2 \sigma_y^2 = \frac{1}{2} m_{\text{Rb}} \sigma_{v_y}^2 = \frac{1}{2} k_B T. \tag{3.10}
\]

The transverse size of the cloud can thus be calculated by:

\[
\sigma_y = \sqrt{\frac{k_B T}{m_{\text{Rb}} \omega_y^2}}. \tag{3.11}
\]

If we consider \( \omega_y = 1500 \text{s}^{-1} \) and \( T = 20 \mu \text{K} \) we have that \( \sigma_y \approx 30 \mu \text{m} \).
From the size of the atomic cloud and the number of atoms in the dipole potential (measured as explained in 3.4.3) we can calculate the atomic density. The longitudinal size of the cloud can be obtained by a gaussian fit on the integration along the x axis of the picture taken by the camera; we typically have $\sigma_x = 3\text{mm}$. If we consider $10^6$ atoms in the dipole trap, we have that the typical densities are $10^{10}-10^{11}$ atoms/cm$^3$. It is important to know the atomic density in the dipole trap because there are many effects that critically depend on the density of the sample. For example, if we have high densities, we can not neglect the collisions between atoms.

### 3.6 Time sequence

The main part of the experiment is easily controlled by the graphical interface of a program in MATLAB (see the figure (3.19)). This program loads a big matrix which is then transferred to an AOI (Analog Output Input) card. The different output channels of the card are connected to the VCOs, which control the double and single pass AOMs, and to the current supply for the MOT coils. The VCOs have been calibrated and it is enough to write the detuning or the power (with respect to the maximum) that we want to have in the graphical interface, and the program calculates the proper voltage to send to the VCOs. One of the channel is used as trigger for the CCD camera.

![Figure 3.18: Typical time sequence for our experiment.](image-url)
Thanks to this program we have a good flexibility in our experiment, we can rapidly change the parameters and the time sequence. Nevertheless some improvements are still necessary: the main limitation is that it is not possible to have different phases with very different time duration because the time step is fixed for the complete matrix (typically 500µs or 1ms). If we use a time step too little the matrix is too big and can not be charged, on the other side we can not have a phase shorter than one time step and thus the time step has not to be bigger than the shortest phase.

In the most part of the measurement the time sequence used is the one reported in the picture (3.18).

Figure 3.19: Graphical interface of the program in MATLAB. The time step, the duration of the phases and the main part of the parameters for the beam can be easily varied from here. We note that the calibrations for the VCOs are not yet optimized and the values typed in this interface (e.g. for the detuning of the repumper) do not always correspond to the real values transmitted to the AOI card.
Chapter 4

Observations

Once the cold atomic sample is realized, we illuminate it with the CARL beam. This beam acts on the atoms with two different forces: the radiation force and the dipole force (see chapter 2). The first pushes the atoms along the direction of propagation, the second tends to trap them in the center of the beam (*optical dipole trap*).

The simplest expressions to keep in mind to have a first understanding of the behaviour of the atoms are:

\[ U_{dip} = \frac{I}{I_{sat}} \frac{\Gamma}{8\delta} \hbar \Gamma \]  
(4.1)

\[ F_{rad} = \hbar k\frac{\Gamma}{2} s \text{ where } s = \frac{I/I_{sat}}{4(\delta/\Gamma)^2} \]  
(4.2)

which give the depth of the dipole potential and the force on one single atom due to the radiation force. We indicate with \( \delta \) the detuning between the beam frequency and the atomic resonance. Equation (4.2) is an approximate expression, valid for \( I/I_{sat} \ll (\delta/\Gamma)^2 \) and \( s \ll 1 \). The intensity at the center of the laser beam is deduced from the measured total power \( (P_{dip}) \) and the measured beam waist \( (w_0) \):

\[ I = \frac{2P_{dip}}{\pi w_0^2} \]  
(4.3)

We can see that there are three parameters of the beam which play an important role: the power, the waist and the detuning from the atomic resonance. We control the detuning of the beam by the current of the DFB laser (see Fig. (3.4)), the power by turning a \( \lambda/2 \) plate in front of a polarizing beam splitter, and we fix the waist at 200\( \mu \)m. The focus of the gaussian beam is at the center of the MOT. Typical values for the radiation pressure are \( F_{rad} = 10^{-25} - 10^{-24} \text{N} \) which correspond to 1000 - 10 000 photons exchanged per second per atom. The depth of the dipole trap is usually \( 10^{-26} - 10^{-27} \text{J} \) which corresponds (dividing by the Boltzmann constant \( k_B \)) to a temperature of a few hundreds of \( \mu \text{K} \), which is larger than the kinetic energy of the initial cloud of cold atoms. We have observed
that with a reduced depth of the dipole potential (e.g. lower power) the dipole potential becomes too weak and all atoms are accelerated in free fall due to gravity.

We performed measurements of the displacement and the velocity of the center of mass of the atomic cloud in the dipole potential as function of different parameters. The protocols used for these kinds of measurement are reported in 3.4.2 and 3.4.3.

4.1 First measurements: what we expected and what we observed

In order to extract the information about the radiation force, we measure the spatial displacement of the atoms (or better of the center of mass of the atomic cloud). Assuming a constant force acting on the atoms during all the interaction with the laser beam, the final position of the center of mass of the atomic cloud is given by:

$$\langle x \rangle = \langle x_0 \rangle + \langle v_0 \rangle t_{dip} + \frac{1}{2} \frac{F_{mRb}}{mRb} t_{dip}^2$$

with $t_{dip}$ the time of interaction between the dipole beam and the atoms. We take $\langle x_0 \rangle = 0$ as the center of mass of the MOT and $\langle v_0 \rangle = 0$, as the atomic cloud has a zero average velocity. This can be seen from the vertical free fall of the atomic cloud which is not trapped by the dipole potential. From the first measurements we noted some differences in the atomic behaviour from what is expected if there was only the radiation pressure acting on them. For this reason we compare the value of the radiation pressure calculated by (4.2) with the value of the hypothetical constant force which gives the measured displacement calculated by:

$$F_{meas} = \frac{2\langle x \rangle m_{Rb}}{t^2}$$

The presence of a CARL effect would give an enhanced backscattering of photons, resulting in a larger displacement of the atomic center of mass. It is thus convenient to compare the total force to the radiation pressure force expected for single atoms ($F_{meas}/F_{rad}$ should be bigger than 1). In fact in the CARL process the photons absorbed by the atoms are not spontaneously reemitted, as for the radiation pressure, but by stimulated emission they are scattered in the opposite direction of the incoming beam. Indeed, if all light was scattered exactly in backward direction, the momentum exchange per fluorescence cycle would be $2\hbar k$, instead of an average of $\hbar k$ if the emitted photons would be scattered in $4\pi$ steradian. This leads to an increment in the pushing effect of the dipole beam on the atoms. As the enhanced backscattering is due to a collective effect, we expect this effect to increase with the number of atoms trapped in the dipole potential and thereby involved in the process.
We made measurements investigating the dependence of the displacement on the different parameters (power and detuning of the laser, interaction time and number of atoms).

**Dependence on the interaction time**

If the force which acts on the atoms is constant during all the interaction with the dipole beam (in particular if there is only the radiation pressure acting on the atoms), the curve of the displacement as function of the interaction time is parabolic. As the starting position and velocity are zero, the best fit on the experimental data is expected to be without the linear term. The quadratic term is the acceleration.

![Graph showing the displacement of the center of mass of the atoms in the dipole trap as function of the interaction time (t\textsubscript{dip}).](image)

Figure 4.1: *Displacement of the center of mass of the atoms in the dipole trap as function of the interaction time (t\textsubscript{dip}).* We can not have interaction times shorter than 20 ms because we can not distinguish the atoms in the dipole trap from the atoms which are falling due to gravitation.

The graph (4.1) shows the dependence of the displacement of the atoms on the interaction time. The value of the acceleration due to radiation pressure for the parameters used in this experiment (P=100mW δ=-81GHz) is: \( a\text{rad} = \frac{F\text{rad}}{m\text{Rb}} = 17m/s^2 \). If we fit the experimental data with a quadratic function without the linear term we do not get to a good agreement with the data. In order to reach a better agreement we need to introduce in the fitting formula either a
negative initial displacement (about 3mm), which seems incompatible with the origin of the MOT, or a negative linear term i.e. a negative initial velocity \( v_i = -0.13 \text{m/s} \). In all the cases we deduce an acceleration of about \( 8 \text{m/s}^2 \). We checked that there is no initial velocity in the absence of the dipole laser by verifying that the atoms trapped in the MOT fall straight once the trap is switched off. At this point we do not have any explanation or theoretical model to explain the origin of such a negative initial velocity.

**Dependence on the power of the dipole laser**

The displacement due to a constant force (if the initial velocity is zero) is proportional to the force. The radiation pressure is proportional to the power of the beam and its dependence on the detuning is \( 1/\delta^2 \).

In Figure (4.2) we show the displacement of the center of mass of the atomic cloud as a function of the power of the dipole laser. The linear dependence we observe is consistent with what one expects from simple radiation pressure. A direct comparison with the expected radiation pressure force give a correction factor of 0.5. The force is thus half of what would expect from a simple model.
One should note however that this correction factor might be explained by the non homogeneous intensity profile of the laser beam across the atomic cloud and by a reduction factor due to the Clebsch-Gordon coefficients of the transition. Both effects give rise to a smaller force than for two level atoms at the center of the laser beam. A quantitative comparison would thus be required to check whether this smaller value could be explained with a single atom model.

In figure (4.3) we show the normalized force \( F_{\text{meas}}/F_{\text{rad}} \) extracted from figure (4.2) assuming a constant force (see equation (4.5)). Even though in Fig. (4.2), the displacement of the atomic center of mass and hence the force appears to depend linearly on the laser power, the normalized curves indicate a slight oscillatory behavior. A more systematic study with better signal to noise ratio and statistics would be required to confirm and investigate this deviation from the single atom radiation pressure expectation.

**Dependence on the detuning of the dipole laser**

Let us now turn to the dependence of the atomic displacement on the laser detuning from the atomic resonance. As we will see, more striking features are observed. In figure (4.4) we show the displacement of the center of mass of the atomic cloud as a function of the detuning of the dipole laser.

The laser detuning has been changed via the current of the pump laser (DFB).
Figure 4.4: Displacement of the center of mass of the atoms in the dipole trap as function of the detuning of the laser: we took the same data three times and we note that the error is smaller than the amplitude of the oscillation.

Figure 4.5: Plot of the normalized force extracted from figure (4.4) assuming a constant force.
Figure 4.6: Plot in log-log scale of the measured force. The slope of the linear fit indicates a detuning dependence of $1/\sqrt{|\delta|}$.

As this also affects the output power of the master laser and hence the power of the dipole beam, we have carefully adapted the power incident on the atomic cloud by turning a $\lambda/2$ plate in front of a polarizing beam splitter. We have thus maintained a constant power for each point of Fig. (4.4). As one can clearly see in the figure, a pronounced oscillation appears.

This oscillation has been very robust and we have systematically observed this oscillation for a large range of parameters (laser power, polarization, atom number). This oscillation can not be explained by single atom radiation pressure force. This new feature (more striking than in the power dependence) has thus been at the center of many speculations and experimental investigations.

In order to get some further insight into a possible explanation, we plot, in figure (4.5) the normalized force $F_{meas}/F_{rad}$ (see eq.(4.5)), as a function of detuning. Beyond the oscillatory behaviour, this normalized curve shows a qualitative difference from what one expects from single atom radiation pressure force. Indeed, the measured force does not decrease as fast as $1/\delta^2$. This shows up as an increase of the normalized force when the detuning is increased. A plot in log-log scale of the measured force (see figure (4.6)) indicates a detuning dependence of $1/\sqrt{|\delta|}$. The measured force thus decreases as the detuning is increased (which can be seen on fig (4.4)), but not as fast as the radiation pressure force (which explains the increase of the normalized force in fig (4.5)).
The period of these oscillation is 30GHz. We looked for variations in the period of oscillation or in their amplitude changing the power and the polarization of the beam or the time of interaction, and using a beam with a wavelength closer or further from resonance, but we did not note any change. The peaks are always at the same position and the distance between two peaks is constant (see the graph (4.7)). We also tried to vary the detuning varying the temperature instead of the current and we detected the same behaviour.

![Graph showing displacement of the atoms as function of the detuning of the dipole beam for different interaction times and different powers. The peaks are always at the same place, and the distance between two peaks is always 30GHz.]

**Figure 4.7: Displacement of the atoms as function of the detuning of the dipole beam for different interaction times and different powers: the peaks are always at the same place, and the distance between two peaks is always 30GHz.**

### Dependence on the number of atoms

As we have not been able to find any simple explanation of the above described oscillatory behaviour in terms of single atom radiation pressure force, we have turned to a study of the displacement of the center of mass of the atomic cloud as a function of the number of atoms in the dipole trap. We record a diminution in the pushing force (and thus in the final position of the atoms) if there are more atoms in the dipole trap (see the graph (4.8)). This behaviour too has been very robust. We have systematically observed a reduction in the position...
for big atom numbers in the dipole trap, for a large range of parameters (laser power and detuning, polarization, interaction time). We looked for a correlation between these two effects and, for this reason, we traced the curve as function of the number of atoms for different detunings (maximum and minimum of the oscillation), maintaining the same power of the beam and the same interaction time. We did not note any qualitative difference between these curves (see Fig. (4.8)).

We note that in the CARL process the enhanced backscattering is due to a collective effect and thus we expect this effect to increase with the number of atoms trapped in the dipole potential and thereby involved in the process. The pushing effect too should thus increase with the number of atoms. This is not the case in what we observed.

This reduced pushing effect is a collective effect in the sense that it depends on the total number of atoms and it is the confirmation that the interaction between the atoms and the dipole beam can not be understood in terms of a single atom effect. We can extrapolate the single atom pushing effect from the graph and we find that this value is in good agreement with the radiation pressure.

![Figure 4.8](image.png)

Figure 4.8: Displacement of the atomic center of mass as function of the number of atoms in the dipole trap. The lines are only traced to guide the eye and do not correspond to any particular explication. We note a change in the slope when the atom number is larger than about $10^6$. 
4.1.1 First conclusions: collective effect

From these measurements of the displacement of the center of mass of the atomic cloud trapped in the dipole beam we can extract some interesting information. If we suppose that there is a constant force which acts on the atoms, we know that its magnitude is compatible with the radiation pressure force. Anyway the observed experimental results cannot be explained by single atom behaviour and thus we have to suppose the presence of one or more other effects. We can summarize the most striking points we can extract from our measurements and that can not be explained by the radiation pressure:

- Diminution in the pushing effect with increasing the number of atoms: this reduced pushing effect is a collective effect in the sense that it depends on the total number of atoms

- Presence of oscillation in the final position of the atoms varying the detuning of the dipole beam. The constant period of these oscillation is 30GHz and it is independent on the power of the beam, on its polarization and on the time of interaction.

- Diminution in the pushing effect, compared with the one due to the radiation pressure, for a frequency of the beam closer to the atomic resonance.

4.2 Transient effect

In the data analysis in the previous section 4.1, we considered the presence of a constant force which acts on the atoms during all the interaction between the dipole beam and the atoms, and we compare this force with the radiation pressure. However we can not exclude the possibility that a transient process is present, i.e. we can have an effect which acts on the atoms only for a time shorter than the period in which the dipole beam is on, and then, after this transient, only the single atom radiation pressure force remains. We have thus turned to the investigation of the possible presence of this kind of effect.

If we consider a transient effect the formula (4.4) becomes:

\[
\langle x \rangle = \frac{1}{2} \frac{F_{\text{rad}}}{m_{\text{Rb}}} t_{\text{dip}}^2 + \int_0^{t_{\text{dip}}} \int_0^t \frac{F_{\text{tran}}(t')}{m_{\text{Rb}}} dt' dt
\]

if we consider \( \langle x_0 \rangle = 0 \) and \( \langle v_0 \rangle = 0 \). The additional term can have the opposite sign of the radiation force, as it is suggested by the fact that the pushing effect decreases if the number of atoms increases, the fit of the curve of the displacement as function of the time of interaction gives a negative initial velocity and the displacement measured is typically a little smaller of what is expected for the radiation pressure.
In order to experimentally investigate the presence of a transient effect, it is necessary to look at the behaviour of the atoms in the first hundreds of microseconds of interaction with the beam. This is not possible because we are not able to distinguish the atoms trapped in the dipole potential from the others if they are not sufficiently pushed. For this reason we need to create another initial transient once the atoms in the dipole trap are already well separated from the falling MOT due to gravitation. We load the atoms in a shallow dipole potential for about 30ms and then we suddenly increase the power of the dipole beam (and thus the depth of the dipole well) by a factor five. In this way we do not recreate exactly the initial condition (e.g. for the temperature) but anyway we have a transient that we are able to trigger and study.

We perform the measurements of the final velocity of the atoms after few milliseconds of interaction with the strong beam. If there is only a constant force acting on the atoms, the dependence of the velocity on the time is linear. In particular the effect of the radiation pressure, when the power is increased by five, is an increase of five in the slope of the line too. We choose to measure the velocity instead of the displacement because it allows us to have a more direct understanding of what is happening. Moreover the deviation in the displacement can be too small to be easily directly detect. The effect of a transient force on the velocity is more evident than on the position.

In the graph (4.9) the experimental data are reported. We have to note that the error in the measurement is big. The error bars shown in the figure have been calculated measuring three times the same velocity. A more systematic study with better signal to noise ratio and statistics would be required, at the moment we can anyway say that we have never recorded a bigger velocity after that the power is increased. We always have a “transient”, of about 1ms, in which the velocities are smaller.

We have not yet performed a systematic study of this interesting effect. Anyway we made a first attempt to investigate a possible connection between this transient diminution in the velocity, the reduced pushing effect with big number of atoms and the oscillation as function of the detuning. At the moment we have no conclusive observation because the systematic and statistical errors are too large. Moreover we are limited in the choice of the parameters as we need to trap the atoms with a beam having a low power and so we need a small detuning too. We thus can not vary a lot the depth of the first dipole potential and, consequently, the number of atoms. However we did not note a qualitative difference in the curve obtained with a beam detuning of -30GHz (maximum of the oscillation) and -15GHz (minimum of oscillation). In summary from this data a reduction in the velocity of the atoms seems to be present even though, at this moment, the systematic and statistical errors are too large to proof it.
CARL emission has been predicted and observed to be a transient effect in the absence of velocity damping [6]. In the experiments performed in a high finesse cavity, the change in velocity as the atoms scatter light in the backwards direction induces a change in the Doppler shift of the scattered photons. When this Doppler shift is larger than the spectral width of the cavity, the enhancement due to the cavity is no longer active and the self-induced modulated potential vanishes. As a result, the CARL emission becomes a transient effect.

In the absence of a build-up cavity, this Doppler shift does no longer have the same impact on the collective emission. Nevertheless, one can compare the gain width of the CARL emission to the Doppler broadening. Preliminary theoretical calculations and simulations [24] predict that CARL emission in absence of cavity is better described as superradiance and it is necessary for the superradiant gain to overcome the Doppler broadening. First simulations predict that the backscattered light (and thus the mechanical effect on the atomic cloud) is still transient in nature, even in the absence of a cavity. The experimental effort on the study of transient effects have been motivated by such considerations. At present, further experimental and theoretical investigations are required to both understand the features observed in the experiment and the feasibility of CARL and superradiance.

Figure 4.9: Velocity of the center of mass of the atoms. We indicate with $t_{\text{dip2}} = 0$ the instant in which the power of the dipole beam is increased.
effects in a cold cloud of atoms without external built-up cavity.

### 4.3 Other measurements

In order to further investigate the atomic behaviour and to exclude the presence of some possible parasite effects we performed many other measurements. Some of these gave us the possibility to advance in the understanding of the processes which occur during the interaction between the atoms and the CARL beam. For other further measurement are needed to extract conclusive observations.

**Detection of the backscattered light**

We have concentrated our efforts in the detection of the effects of the interaction between the atoms and the CARL beam via the measurement of the positions and velocities of the atomic cloud. However we have also tried to detect the light backscattered by the atoms. If bunching occurs as it is expected in the CARL process, the atoms self-organize in a periodic lattice, and the incoming light should be backscattered. This fact affects the atomic velocity. We have calculated from the radiation pressure that each atom typically exchanges with the beam about 1000 photons per second and in our sample there are about $10^6$ atoms. If we consider a variation in the velocity of about 10% from the value due to the radiation pressure only, we thus expected that about $10^8$ photons per second are backscattered by the CARL process and this corresponds to a power of few tens of picoWatt.

In the CARL beam path, in front of the vacuum cell, there is a polarization cube and a $\lambda/4$ plate. The light backscattered from the atoms passes twice through the $\lambda/4$ plate and leaves the polarization cube in the orthogonal channel, which allows the detection of the reflected light from the atoms. We have aligned the CARL beam retroreflected on itself by a mirror on a photomultiplier (PM). The PM is very sensitive and it is able to detect small power (less than nW), but it is necessary to pay a lot of attention in the reduction of the background. We worked hard on this and with a well closed box around the PM and a good extinction of the light of the CARL beam back reflected we have reached a background of few thousands of photon-counts par second.

Until now we have not detect any backscattered light, but we can not conclude that there is no bunching. It is in fact necessary to perform further measurements and to better focus the light on the PM. The detector is in fact placed far away from the atoms to reduce the background but in this way we lose a lot of photons of the backscattered light too, especially if this light is diverging. Moreover, if the CARL process is expected to be transient and to last few hundreds of microseconds, we need to look for the backscattered light only during the first milliseconds of
interaction with the beam whereas, up to now, we have only investigated a larger temporal scale.

The light backscattered in the CARL process is expected to be very weak. We have thus also investigated the presence of a lattice by a probe laser beam in resonance with the atomic transition as the Bragg scattering is in fact more effective for resonant light. In this case too, we have not detected any light backscattered by the atoms.

Possible presence of residual molasses

The MOT beam and the repumper are switched off by the single pass AOMs but, even when we switch the voltage control to zero, there is a small part of the incoming beam which is scattered by the AOM in its first order. For this reason it is possible that a small friction force, due to a residual molasses, acts on the atoms during the interaction with the CARL beam.

To exclude this possibility we put a mechanical shutter (Clic-Clac) in the MOT beam path, before it is divided into the six beams. We detect no changes in the measurement with and without clic-clac and we can thus conclude that, even if a residual molasses is present, it does not have any detectable effect on the interaction.

Possible presence of light on atomic resonance

The CARL beam is always far detuned from the atomic transition, but we cannot exclude the presence of a small, but maybe not negligible amount of light on resonance with the atoms. This light can come either from the MOPA’s spontaneous emission which has a very large band width, or from the DFB itself which can also have some smaller peaks of emission besides the main one. If this light is really present we need to take it into account in the data analysis. In particular we need to modify the equation for the radiation pressure considering the effect of all the spectral frequencies present in the beam. A small amount of light on resonance can have the same effect on the atomic position as a far bigger amount of light far detuned. The optical thickness is in fact very different in the two cases. Furthermore, as the optical thickness depends on the density of the sample, and thus on the number of atoms, it is also possible that the pushing effect of the photons on resonance depends on the number of atoms too.

We performed a fit on the experimental points for the oscillation in the final position as function of the laser detuning. For this fit we suppose that in the DFB emission there are some small peaks at the frequencies of $n \times 30$ GHz and we estimate the radiation pressure due to this kind of emission. The free parameters of this fit are the amplitudes of the peaks whereas their width is fixed at 10GHz. We
can see in the figure (4.10) that the experimental data can be well approximated by a fit of this kind.

Figure 4.10: Fit for the oscillatory dependence of the radiation pressure on the laser detuning.

This kind of collateral effect could explain some of our measurements: it is thus strictly necessary to remove from the beam all the frequencies which can affect the atomic behaviour and repeat the measurement. We tried to filter the beam with an etalon. It works as a Fabry-Perrot cavity and the transmitted beam has a narrow spectral width around the frequency we are interested in and in particular the photons on resonance are filtrate. We used a 0.5mm thick etalon with a free spectral range of 200GHz and a Finesse of 60. Unluckily the beam power has been too much reduced by the etalon and we are hardly able to trap the atoms in the dipole potential. It is thus impossible to perform the experiment in this condition and at the moment we do not have conclusive observations. It is also possible to filter the beam by a cell of hot Rubidium atoms. In this way we get rid of the photons on resonance as they are absorbed by the atoms. Doppler broadening allows us to filter a wide line of frequencies around the resonance (600MHz).

To investigate the effects of an additional radiation force, we have added on the atoms a weak probe beam in resonance with the atomic transition and aligned with the CARL beam, and we looked for the dependence of the atomic velocity on
the power of the probe. In this case too, a more systematic study with better signal
to noise ratio and statistics would be required to extract some useful information
from this type of measurement.

**Experiment with constant number of atoms**

The most difficult parameter to control in our experiment is the number of atoms
trapped in the dipole potential. If we repeat the same measurement several times
without changing any parameter we note that the number of atoms is not constant.
Furthermore it obviously depends on the depth of the dipole potential and, for
this reason, when we perform the measurement varying the power or the detuning
of the beam we can not keep constant the number of atoms (see Fig.(4.11)). We
can thus not well distinguish the dependence of the atomic displacement on the
different parameters.

![Figure 4.11: Dependence of the atom number on the laser detuning.](image)

We performed some measurements of the dependence of the final displacement
of the atoms on the beam detuning keeping constant the ratio $P/\delta$ (and thus the
potential depth $U_0$) instead of the beam power. In this way the atom number is
nearly constant for a large range of $\delta$. As in the previous measurement we have
carefully adapted the power incident on the atomic cloud by turning a $\lambda/2$ plate
in front of a polarizing beam splitter. We have thus maintained a constant $U_0$ for
each point. We have not a good enough statistic for this type of measurement and further measurements are necessary to confirm the result. The oscillations seem however to be still present even thought they are less pronounced (see Fig. (4.12)).

![Graph](attachment:image.png)

**Figure 4.12:** *In this measurement we keep constant the depth of the dipole potential. The oscillations in the position as function of the laser detuning are still present and the atom number is constant.*

**Hypothesis investigated**

During the last months several hypothesis have been formulated to find an explication of the experimental observations. Some of these are briefly reported in this chapter. They have not been deeply studied and investigated because further measurements always carried the discussion towards a new hypothesis.

One of these hypothesis is the possible presence of collisions in the dipole trap. The attractive force between two atoms in the ground state (van der Waals’s force) decreases as \(1/r^6\) (where \(r\) is the distance between two atoms) and we thus create molecules if the density is very high because the atoms need to be very close to each other to experience the mutual attraction. However, if one of the two atoms is in the excited state, the attractive force decreases slower \((1/r^3)\) and lower densities are enough. If in our atomic sample the creation of molecules occurs, we have new
energetic molecular levels. The energetic gap between these levels can be of tens of GHz [25]. The cross section of the dipole beam on the atoms thus depends on the wavelength of the laser. When the dipole beam is in resonance with one of the molecular transitions, the probability of scattering increases and we can thus have an increase in the pushing effect. We have wondered if the oscillations in the displacement of the atoms as a function of the detuning of the laser beam could depend on the excitations of the different energetic levels of the molecules created by an atom in the ground state and an atom in the excited one. For the moment we have excluded this possibility because the distance between the peaks we have detected is constant, whereas the molecular energetic levels become more and more closer to each other approaching the atomic resonance. Moreover the density in our sample is probably too small to allow this kind of effect.

Other effects which might occur in a hight density sample of cold atoms is the creation of atomic clusters. At the moment we do not have any evidence of this kind of effects, and we attribute the fact that the images of the dipole trap are not homogeneous to a detection problem. Furthermore in this case too the densities of our sample are probably too small to allow the creation of atomic structures [26].

Other hypothesis are related to the geometrical characteristics of the atomic cloud, in particular to its cigar shape. From qualitative observations we have noted that the longitudinal size of the cloud is not constant and in particular it seems to slightly oscillate as function of the laser detuning. This has lead us to the study of Mie scattering. Preliminary calculations have however shown that, considering the gaussian density distribution of the cloud, our measurements cannot be explained by this kind of process.

None of the hypothesis we have so far investigated are able to fully explain our experimental measurements and the discussion is still open.
Conclusion

In the last months I have had the opportunity to work in an advanced laboratory for cold atoms. I took part in the experiment performed on a new experimental setup for Rubidium atoms which ended to be placed during my staying at INLN. The first part of my work at INLN consisted in the participation of the implementation of this new experiment. I have in particular been in charge of the realization and control of the dipole laser. I had thus the opportunity to directly study and perform the most common techniques to cool and trap atoms, the magneto-optical trap and the optical dipole trap, and the delicate transfer of the atoms from one trap to the other.

The experiment is devoted to the study of the collective effects induced on cold atoms by a powerful far detuned beam. Our interest is, more specifically, the superradiant regime of CARL i.e. the possibility to observe spontaneous self-organization in a cloud of cold atoms (not condensed) in the absence of an optical cavity.

The most part of our experimental observations are measurements of the displacement and of the velocity of the center of mass of the atomic cloud trapped in the dipole potential after the interaction with the CARL beam. We can say, in summary, that the observed experimental results cannot be explained by single atom behaviour with the large detuning and power we have measured and they suggest the presence of collective effects. During the last months several hypothesis have been formulated to find an explication of the experimental observations but none of these hypothesis are able to fully explain our experimental measurements and the discussion is still open.

We can summarize the most striking points which cannot find their explications in the radiation pressure as follow. The best fit on the displacement measured as function of the interaction time gives an initial negative velocity as shown in figure (4.1). As regard the dependence of the displacement on the laser frequency we observed an oscillation in the final position of the atoms. The period of this oscillation is 30GHz and it does not depend on the power or the polarization of the beam and on the time of interaction. Furthermore we detect a diminution in the pushing effect, compared with the one due to the radiation pressure, for a
frequency of the beam closer to the atomic resonance (Fig. (4.5)). Moreover we observed a reduced pushing effect when the number of atoms is increased (Fig. (4.8)): this reduced pushing effect is thus a collective effect in the sense that it depends on the total number of atoms. In the measurements of the final velocity we have been able to look for transient effects: we first trap the atoms in a shallow dipole trap and then we increase by a factor of five the laser power. We measured the atomic velocity after few milliseconds of interaction with the powerful beam and we detect a reduction in the mean velocity of the atoms (Fig. (4.9)).

The future experimental and theoretical study of these effects would be able to lead to a further understanding of the mechanical effect induced by a powerful far detuned beam on a sample of cold atoms trapped in a dipole potential.
Riepilogo dei principali contenuti di questa tesi

Ho svolto il mio lavoro di tesi presso l’istituto INLN a Nizza. Il gruppo di lavoro sugli atomi freddi dell’INLN sta sviluppando in questi anni tecniche per produrre grandi campioni di atomi freddi [15, 16] con lo scopo di studiare i diversi meccanismi per intrappolare fotoni in un campione di atomi (la localizzazione di Anderson della luce) e il random laser. Per raggiungere questi obbiettivi è necessario avere campioni di atomi con alte densità e per questa ragione ci si aspetta di incorrere nello studio di effetti collettivi: la superradianza è uno di questi effetti. È in questo contesto che si colloca lo studio del laser a rinculo atomico collettivo (CARL) in assenza di cavità, argomento di questa tesi.

Il CARL è un processo in cui gli atomi, interagendo con un forte campo di pompa, si auto-organizzano formando un reticolo di densità con periodicità di mezza lunghezza d’onda. Gli atomi ordinati in questo reticolo amplificano coerentemente e esponenzialmente un campo di sonda inizialmente molto piccolo che si propaga in direzione opposta al campo di pompa. Il CARL è un laser nel senso che in questo processo un campo inizialmente molto piccolo, originato da fluttuazioni o da fotoni emessi spontaneamente, cresce esponenzialmente fino a raggiungere un valore di saturazione. D’altra parte è diverso da un processo laser classico in molti aspetti: non ha una soglia e non raggiunge, se non in particolari condizioni, uno stato stazionario. Il processo CARL è stato predetto da R. Bonifacio et al. nel 1994 [1, 2, 3] e può essere visto come l’analogico atomico del laser a elettroni liberi (FEL) con il quale condivide lo stesso tipo di equazioni. Sono stati fatti molti tentativi per verificare sperimentalmente questo processo con vapori di atomi caldi in cavità ad anello [4, 5], ma solo nel 2003 a Tübingen il gruppo di Ph. Courteille, ha ottenuto i primi risultati sperimentali utilizzando atomi freddi in una cavità ad alta finesse [6]. Un altro aspetto importante del CARL è legato al suo regime quantistico realizzato con i condensati di Bose-Einstein (BEC) [9, 10, 11]. In questo limite il momento scambiato tra la radiazione e gli atomi, che è essenzialmente continuo nel regime classico, diventa discreto [12, 13].

Il nostro interesse attuale è rivolto alla possibilità di osservare la spontanea
organizzazione degli atomi nel reticolo, e quindi l’effetto CARL, in un campione esteso e denso di atomi freddi (ma non condensati) in assenza di cavità. L’esperimento al quale ho partecipato consiste dunque nel preparare una campione di atomi di Rb in una trappola magneto-ottica (MOT) e quindi trasferito in una trappola ottica dipolare. Questo denso campione di atomi freddi è poi illuminato con un potente fascio laser fuori risonanza rispetto alla frequenza atomica. Gli effetti meccanici (collettivi) della radiazione sugli atomi possono essere studiati sia misurando la posizione e la velocità degli atomi stessi dopo l’interazione con questo fascio laser, sia rilevando la luce diffusa dagli atomi. Questo esperimento è stato realizzato su un nuovo apparato sperimentale per atomi di Rb. La prima MOT con questo setup è stata realizzata ad aprile 2008. La prima parte del mio lavoro all’INLN è consistita nella partecipazione all’implementazione di questo nuovo esperimento (mi sono in particolare occupata della realizzazione e del controllo del laser di dipolo); nella seconda parte ho partecipato alle prime misure realizzate con questo nuovo apparato i cui risultati sono riportati in questa tesi.

**Effetti meccanici della radiazione sugli atomi - Forza di scattering e forza di dipolo**

Oltre che energia, un atomo che interagisce con un campo elettromagnetico cambia anche quantità di moto ed è sottoposto quindi ad una forza (la forza di radiazione) che influenza la dinamica del suo centro di massa. Lo scambio di quantità di moto è originata, dal punto di vista quantistico, dall’assorbimento ed emissione di fotoni di momento \( \hbar \vec{k} \). Nella descrizione classica la radiazione induce un momento di dipolo sugli atomi che si accoppia con il campo elettrico della radiazione.

La forza di radiazione è quindi proporzionale al valor medio del momento di dipolo dell’atomo \( \langle d \rangle_\psi \) e al gradiente di campo elettrico [17]:

\[
\frac{d\vec{p}}{dt} \approx \langle d \rangle_\psi \vec{\nabla} E = \vec{F}.
\]

Cosiderando un atomo a due livelli, il valor medio del momento di dipole è \( \langle d \rangle_\psi = d_{12}(\rho_{12} + \rho_{21}) \), dove \( \rho_{12} \) e \( \rho_{21} \) sono gli elementi fuori diagonale della matrice di densità e \( d_{12} \) è il dipolo elettrico.

Se il campo è un’onda monocromatica \( E = E_0 \cos(kx - \omega t) \), la forza di radiazione è diretta lungo il vettore d’onda \( \vec{k} \). Approssimando le equazioni per l’evoluzione degli elementi della matrice di dipolo nel caso in cui il tempo di vita dello stato eccitato dell’atomo \( \Gamma^{-1} \) è molto minore del tempo caratteristico di variazione del momento atomico, si ha che la forza di radiazione, che prende in questo caso il nome di *forza di scattering* o pressione di radiazione, assume la
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La forza di scattering può essere interpretata, in termini di fotoni, in questo modo: ogni fotone assorbito da una spinta all’atomo nella direzione di incidenza; il fotone è poi riemesso spontaneamente dall’atomo in una direzione casuale. La forza media dopo molti cicli di assorbimento ed emissione è quindi diretta lungo la direzione di propagazione del laser incidente e può essere usata per rallentare gli atomi. È proprio questa forza che sta alla base della trappola magneto-ottica.

Se consideriamo un campo di intensità non costante si ha anche una forza diretta lungo la direzione di variazione dell’intensità che viene chiamata forza di dipolo. Nel nostro esperimento il campione di atomi viene illuminato con un fascio laser gaussiano focalizzato. Gli atomi sperimentano quindi, oltre alla forza di radiazione, un’intensa forza trasversale e una debole forza longitudinale. Introducendo una dipendenza dalla posizione del campo elettromagnetico nell’equazione della forza di radiazione e eseguendo le opportune approssimazioni per $\Delta \gg \Gamma$ e $\Delta \gg \Omega$, si ha che la forza di dipolo assume la seguente forma:

$$F_{dip} \approx -\frac{\partial}{\partial x} \left( \frac{\hbar \Omega^2}{4\Delta} \right).$$

Notiamo che la forza di dipolo è proporzionale al gradiente dell’intensità e deriva da un potenziale:

$$\vec{F}_{dip} = -\vec{\nabla} U_{dip} \quad \text{dove} \quad U_{dip} \approx \frac{\hbar \Omega^2}{4\Delta} = \frac{\hbar \Gamma}{8\Delta} \frac{I}{I_s}.$$ 

Se consideriamo un fascio gaussiano si ottiene che il potenziale di dipolo ha anch’esso, trasversalmente, un profilo gaussiano e il minimo di potenziale è dato da:

$$U_0 = \frac{2P}{\pi w_0^2} \frac{1}{I_{sat}} \frac{\hbar \Gamma}{8\Delta} \frac{I}{I_s}.$$ 

dove $w_0$ è il beam waist del fascio gaussiano. Notiamo che per intrappolare gli atomi in questo potenziale è necessario che il fascio sia disaccordato verso il rosso ($\Delta < 0$). Solo in questo caso infatti la forza agisce nella direzione in cui aumenta l’intensità e il potenziale è attrattivo. È inoltre necessario che l’energia cinetica degli atomi (ovvero la loro temperatura) sia sufficientemente bassa da non permettergli di uscire dal potenziale. Dopo un processo di evaporazione gli atomi nella trappola dipolare raggiungono una temperatura per la quale $U_0 = \eta k_B T$, dove tipicamente $\eta = 5 \div 10$. Gli atomi sono quindi disposti sul fondo della trappola di
dipolo ed è quindi possibile approssimare il potenziale con una trappola armonica per la quale la frequenza di oscillazione trasversale degli atomi è data da:

\[ \omega_y = \sqrt{\frac{4U_0}{m w_0^2}} \]

Notiamo infine che il potenziale dipolare è proporzionale a \( I/\Delta \) mentre la pressione di radiazione (per \( \Delta \gg \Gamma \)) è proporzionale a \( I/\Delta^2 \): aumentando il detuning del laser la forza di dipolo diventa quindi sempre più importante rispetto alla pressione di radiazione.

Valori tipici per la pressione di radiazione nel nostro esperimento sono: \( F_{\text{rad}} = 10^{-25} \div 10^{-24} N \) che corrisponde a 1000 \( \div 10 \) 000 fotoni scambiati per secondo. La profondità della trappola di dipolo è \( 10^{-26} \div 10^{-27} J \) che corrisponde (dividendo per la costante di Boltzmann \( k_B \)) a una temperatura di poche centinaia di \( \mu K \).

**MOT e trappola di dipolo - Apparato sperimentale**


Nel nostro esperimento il raffreddamento e il confinamento degli atomi è realizzato sulla linea D2 del \( ^{85}\text{Rb} \) di lunghezza d’onda \( \lambda = 780nm \) (si veda l’immagine (3.1)). La frequenza del laser di MOT è disaccordata di 2 o 3 \( \Gamma \) dalla trasizione iperfina \( F = 3 \rightarrow F' = 4 \) che è una transizione chiusa. Esiste però una probabilità non nulla di eccitare la transizione aperta \( F = 3 \rightarrow F' = 3 \) e si può quindi avere, in questo caso, un cambio nel livello iperfino fondamentale (da \( F=3 \) a \( F=2 \)) attraverso un’emissione Raman spontanea. Questo succede in media ogni \( 10^3 \) fotoni scambiani tra il fascio e gli atomi. Gli atomi che cadono nel livello \( F=2 \) non vedono più la luce del laser di MOT e sono quindi persi. Perciò è necessario illuminare gli atomi con un altro fascio laser, chiamato repumper, accordato sulla transizione \( F = 2 \rightarrow F = 3' \) in modo da “ripompare” gli atomi nel livello iperfino \( F=3 \).

In figura (3.2) è mostrato uno schema dell’apparato sperimentale per i due laser utilizzati per raffreddare e intrappolare gli atomi (laser di MOT e repumper). Non tutti i parametri e gli allineamenti di questo nuovo apparato sono ottimizzati e, al
momento, vengono intrappolati tipicamente $10^8$ atomi alla temperatura di \(100\mu K\) e il diametro del campione sferico di atomi è di circa 3 mm.

Per permettere il trasferimento degli atomi dalla MOT alla trappola dipolare è necessario che la densità del campione atomico sia alta e la sua temperatura sia bassa. Abbiamo provato diversi protocolli per permettere il trasferimento; il procedimento infine adottato e utilizzato per tutte le misure presentate in questa tesi prevede, dopo aver caricato gli atomi nella MOT, una fase di Dark MOT (per aumentare la densità) e una fase di *melassa ottica* (per abbassare la temperatura). I parametri di una sequenza tipica sono riportati in tabella (3.1).

Dopo la fase di melassa ottica il campione atomico viene illuminato con un potente fascio laser (laser di dipolo) che intrappola gli atomi in una trappola di dipolo e contemporaneamente svolge la funzione di campo di pompa per il processo CARL. Lo schema dell’apparato sperimentale per questo laser è riportato nella figura (3.3). Il fascio è ottenuto da un laser DFB (Distributed FeedBack) e in seguito amplificato da un MOPA (Master Oscillator/Power Amplifier). La larghezza spettrale della luce emessa dal DFB è stretta (circa 5 MHz) ed è possibile controllare la sua lunghezza d’onda variando la temperatura del diodo o la corrente del laser: \(\lambda = \lambda(I, T)\) aumenta con T e con I. La calibrazione della lunghezza d’onda del laser (riportata in figura (3.4)) è stata realizzata, per una temperatura fissata, a incremento di corrente e contando i picchi di trasmissione di una cavità Fabry-Perrot il cui Free spectral range è stato precedentemente misurato (\(FSR = 810 MHz\)). Per conoscere il valore assoluto della lunghezza d’onda è necessario conoscere a quale corrente corrisponde la risonanza atomica. Per questa ragione una piccola parte del fascio viene fatta passare attraverso una capsula di Rb a temperatura ambiente e la luce trasmessa dopo il campione viene rilevata. In questo modo, modulando la corrente del laser, viene eseguita la spettroscopia per assorbimento del fascio e, quando la luce laser ha una frequenza prossima a quella di risonanza, è possibile vedere sull’oscilloscopio quattro picchi di assorbimento relativi alle transizioni dai livelli iperfini fondamentali F=2 and F=3 per l’isotopo $^{85}Rb$, e F=1, F=2 per l’isotopo $^{87}Rb$. (si veda la figura (3.5)).

Per essere in grado di spegnere e accendere velocemente il fascio utilizziamo un AOM (Acousto-Optic Modulator) controllato da un VCO (Voltage Controlled Oscillator). È il primo ordine dell’AOM che viene poi focalizzato sugli atomi con una lente posta ad una distanza dalla MOT pari alla sua distanza focale (\(f = 500 mm\)). Il detuning introdotto dall’AOM (80 MHz) è trascurabile rispetto al detuning al quale viene utilizzato il fascio (decine di GHz). Prima della camera a vuoto contenente gli atomi poniamo un cubo polarizzatore e una lamina \(\lambda/4\) per ottenere una buona polarizzazione circolare della luce, in questo modo la luce scatterata indietro dagli atomi passa due volte attraverso la lamina e risulta essere
polarizzata perpendicolarmente rispetto alla luce incidente e viene quindi separata da questa dal cubo. In questo modo è possibile rilevare la luce scatterata dagli atomi anche se la sua intensità è molto bassa.

Il fascio che illumina gli atomi ha infine un waist di 200 $\mu m$, una potenza che può raggiungere i $200 mW$ e la sua frequenza può essere facilmente variata tra la risonanza atomica e un disaccordo verso il rosso di 200 $GHz$.

Gli atomi intrappolati nel potenziale dipolare creato da questo fascio sono $10^6$ e la loro temperatura è di poche decine di $\mu K$ (meno di un quinto della profondità del potenziale). Abbiamo anche misurato la frequenza delle oscillazioni trasversali degli atomi nella trappola (si veda l’immagine (3.17)) che risulta essere di circa 100Hz, in buon accordo con il valore calcolato approssimando il potenziale dipolare con un potenziale armonico.

**Protocolli di misura**

La maggior parte dei nostri dati sono estratti da immagini ottenute con una telecamera CCD. Il guadagno e il tempo di esposizione della CCD (tipicamente 10ms) possono essere facilmente variati. La telecamera registra la fluorescenza degli atomi (si veda l’immagine (3.10)) illuminati dai 6 fasci del laser di MOT e dal repumper. Dalla calibrazione della CCD otteniamo che un pixel della telecamera corrisponde, utilizzando un obiettivo commerciale, a 39 $\mu m$ nella posizione della MOT. La risoluzione della CCD, misurata analizzando l’immagine di una lama di rasoio, è di 200 $\mu m$. Le immagini sono analizzate in MATLAB dopo aver sottratto l’immagine di fondo senza atomi intrappolati, infatti la CCD non registra solo la luce scatterata dagli atomi ma anche la luce diretta dei fasci laser e dell’ambiente. Dopo aver selezionato la zona di interesse, il programma proietta l’immagine sui due assi e esegue il fit gaussiano delle curve che ottiene. Dal valore centrale e dalla larghezza calcolati da questo fit si estraggono molte informazioni. Un esempio di immagine analizzata dal programma è riportata in figura (3.11).

Il valore centrale del fit sulla proiezione lungo l’asse $x$ è il valore della posizione del centro di massa degli atomi intrappolati nel potenziale dipolare. Sottraendo questo valore alla posizione della MOT otteniamo di quanto si sono spostati mediamente gli atomi a causa dell’interazione con il laser di dipolo.

Per la misura della velocità e della temperatura degli atomi nella trappola dipolare è necessario prendere una successione di immagini per tempi di volo (TOF) sempre più grandi, cioè la CCD non è triggerata nel momento in cui viene spenta la trappola di dipolo ma qualche $ms$ più tardi. In questo intervallo di tempo gli atomi non sono soggetti a nessuna forza e mantengono costante la velocità che hanno al termine dell’interazione con il laser di dipolo.

Nella figura (3.12) sono riportate le immagini degli atomi per differenti tempi di volo. Gli atomi cadono a causa della gravità e a causa della loro temperatura finita.
Riepilogo

si espandono dopo essere liberati dalla trappola dipolare. Inoltre la velocità longitudinale non nulla degli atomi fa sì che ci sia anche uno spostamento orizzontale del loro centro di massa. Dal coefficiente lineare del fit sulle posizioni del centro di massa degli atomi per diversi TOF (immagine (3.13)) ricaviamo la velocità al termine dell’interazione con il laser dipolare poiché: $\langle x(t) \rangle = \langle x(0) \rangle + \langle v \rangle T_{OF}$. Per la misura della temperatura è invece necessario considerare la larghezza del fit gaussiano della proiezione dell’immagine lungo l’asse $y$. Il quadrato della larghezza trasversale del campione atomico dipende linearmente dal quadrato del tempo di volo. Dal fit lineare si ricava la temperatura, infatti: $T = a m_{Rb}/k_B$ dove $a$ è il coefficiente lineare e $m_{Rb}$ è la massa atomica del Rb (immagine (3.14)).

Poiché il nostro interesse si concentra su effetti collettivi degli atomi indotti dall’interazione con il laser di dipolo, un importante parametro delle nostre misure è il numero di atomi che partecipa al processo. Abbiamo quindi calibrato la fluorescenza registrata dalla telecamera CCD (che è proporzionale al numero di atomi) conoscendo il numero di atomi nella MOT. Come valore per la fluorescenza consideriamo l’integrale del fit gaussiano del campione atomico lungo l’asse $y$. Il numero di atomi nella MOT è calcolato a partire dalla densità atomica (calcolata a sua volta dallo spessore ottico misurato) e dalla dimensione del campione. Valori tipici di spessore ottico del nostro campione sono $b_0 = 20 \div 25$ dopo la fase di Dark MOT e $b_0 = 8 \div 10$ dopo la fase di melassa poiché quando si spegne il campo magnetico la densità atomica diminuisce.

La maggior parte dell’esperimento è facilmente controllata attraverso l’interfaccia grafica di un programma in MATLAB che carica una grossa matrice e la trasferisce ad una carta AOI (Analog Input Output). I canali di output di questa carta sono collegati ai diversi VCO che controllano gli AOM. In questo modo è possibile controllare da computer le frequenze e le potenze dei laser, l’accensione e lo spegnimento del gradiente di campo magnetico, la sequenza e la durata delle diverse fasi e triggerare la telecamera CCD. La sequenza temporale delle fasi utilizzata nella maggior parte delle misure è quella riportata in figura (3.18).

Osservazioni sperimentali

Abbiamo innanzitutto eseguito molte misure dello spostamento del centro di massa del campione atomico nella trappola di dipolo dovuto all’interazione con il fascio laser in funzione di diversi parametri (tempo di interazione, potenza, detuning e polarizzazione del laser, numero di atomi intrappolati). Abbiamo cercato di mettere in evidenza le differenze nel comportamento atomico rispetto a ciò che ci si aspetta se sugli atomi agisse solo la forza di scattering.

In presenza del processo CARL i fotoni assorbiti dagli atomi non sono riemessi spontaneamente come nel caso della forza di scattering, ma sono riemessi in modo stimolato nella direzione opposta al fascio di pompa. In particolare se tutta la
luce fosse scatterata all’indietro lo scambio di momento per ciclo di fluorescenza sarebbe $2\hbar k$, invece di una media di $\hbar k$ se il fotone riemesso è scatterato con uguale probabilità in tutto l’angolo solido. Per questo motivo la presenza dell’effetto CARL risulterebbe in un aumento dei fotoni scatterati all’indietro agli atomi, e quindi in uno spostamento maggiore del centro di massa. Inoltre, essendo il CARL un processo collettivo, questo effetto dovrebbe essere sempre più importante all’aumentare del numero di atomi coinvolti. Abbiamo quindi confrontato la pressione di radiazione con la forza estratta dallo spostamento misurato $\langle x \rangle$ (considerando una forza costante durante tutta l’interazione): $F_{\text{meas}} = 2 \langle x \rangle m_{\text{Rb}} / t^2$.

Se la forza è costante nel tempo (in particolare se sugli atomi agisce solo la pressione di radiazione), la curva dello spostamento in funzione del tempo di interazione è una parabola. In particolare, poiché la velocità e lo spostamento iniziale sono nulli, il fit migliore dei dati sperimentali dovrebbe avere il termine lineare e il termine noto nulli. Il termine quadratico è l’accelerazione. Il grafico (4.1) mostra la dipendenza dello spostamento degli atomi dal tempo di interazione $t_{\text{dip}}$. L’accordo migliore tra il fit e i dati sperimentali si raggiunge introducendo nel fit un termine lineare negativo, ossia una velocità iniziale negativa. Non siamo in grado, in questo momento, di fornire una spiegazione per questo fit.

In figura (4.2) è invece mostrato lo spostamento del centro di massa in funzione della potenza del laser. La dipendenza lineare che si osserva è consistente con quello che ci si aspetta per la semplice pressione di radiazione. In realtà la forza misurata risulta essere circa la metà della forza di scattering (come si osserva nella figura (4.3)), ma questo può essere dovuto al fatto che nel calcolo della pressione di radiazione non si è tenuto conto del profilo d’intensità del fascio che non è omogeneo nel campione e di un fattore di riduzione dovuto ai coefficienti di Clebsch-Gordon della transizione.

Effetti molto più interessanti sono risultati nello studio dello spostamento del centro di massa del campione atomico in funzione del detuning del laser. Abbiamo variato il detuning del laser variando la corrente del laser DFB e, poiché questo influenza anche la potenza del laser, abbiamo posto particolare attenzione a mantenere costante la potenza del fascio sugli atomi ruotando di volta in volta una lamina $\lambda/2$ posta davanti a un cubo polarizzatore. Come è possibile osservare nelle figure (4.4), (4.5) e (4.7) un marcatamente oscillatorio caratterizza la dipendenza dello spostamento dal detuning del laser. Non è possibile trovare una spiegazione a queste oscillazioni nella forza di radiazione e inoltre, confrontando la forza misurata con quella di scattering, notiamo che la forza misurata non dipende dal detuning come $1/\delta^2$. Abbiamo calcolato che la dipendenza da $\delta$ che meglio descrive i dati sperimentali è $1/\sqrt{\delta}$. Le oscillazioni sono risultate essere un segnale molto robusto e sono state sistematicamente osservate per una grande range di parametri (potenza e polarizzazione del laser, tempo di interazione, numero di
Abbinato inoltre osservato che il loro periodo (30GHz) e la loro ampiezza
non dipendono in modo critico da questi parametri.

Le osservazioni qua sopra elencate non trovano una semplice spiegazione in
termi di pressione di radiazione e di effetti a singolo atomo. La dipendenza dello
spostamento dal numero di atomi nella trappola dipolare è stata quindi studiata,
e i dati raccolti sono riportati in figura (4.8). Abbiamo osservato una riduzione
nella spinta indotta sugli atomi dal fascio laser (e quindi nella loro posizione dopo
l’interazione) all’aumentare del numero di atomi. Anche questo comportamento,
come le oscillazioni, è risultato molto robusto. Questo effetto può essere definito
collettivo nel senso che dipende dal numero totale di atomi coinvolti nel processo
ed è la conferma che l’interazione tra il fascio e gli atomi non può essere compresa
solo in termini di singolo atomo. Possiamo estrapolare dal grafico l’effetto del
fascio sul singolo atomo e notiamo che questo valore è in buon accordo con la
pressione di radiazione.

Abbiniamo finora analizzato i dati ipotizzando una forza costante durante il
tempo di interazione. Non possiamo però escludere la possibilità di essere in pre-
senza di un effetto transitorio, cioè di un processo che agisce sugli atomi solo per
un intervallo di tempo più breve. In particolare il processo CARL in cavità, in
assenza di attrito, è un processo transitorio; anche nel caso superradiante senza
cavità, le previsioni teoriche sono di un effetto transiente. Per questa ragione
abbiamo concentrato l’attenzione verso la ricerca di un effetto nei primi ms di
interazione: per investigare sperimentalmente la presenza di un processo di questo
genere è necessario osservare gli atomi nei primi istanti di interazione con il fascio.
Questo non è purtroppo possibile perché non si è in grado di distinguere gli atomi
intrappolati nel potenziale dipolare dagli atomi che cadono per effetto della gravi-
tazione. è quindi necessario creare un altro transiente iniziale dopo che gli atomi
sont già stati sufficientemente spinti e risultano quindi ben separati. Intrappo-
liamo quindi gli atomi in un potenziale debole e, dopo circa 30ms aumentiamo di 5
volte la potenza del laser e quindi la profondità del potenziale. Abbiamo misurato
la velocità degli atomi nei primi ms di interazione con il fascio potente (4.9). Se
la forza che agisce sugli atomi è costante, la velocità aumenta in modo lineare e in
particolare l’effetto della pressione di radiazione, quando la potenza è aumentata
di un fattore 5, è quello di aumentare di un fattore 5 anche la pendenza della
retta. Gli errori sistematici e statistici in questo genere di misure sono ancora,
al momento, troppo grandi per poter fare affermazioni conclusive sui risultati ed
è necessario lavorare per migliorare il rapporto segnale/rumore. Possiamo però
osservare che dalle nostre misure non risulta mai un incremento della velocità
quando la potenza del laser è aumentata ma, sembra essere presente al contrario,
un transiente di circa 1 ms in cui le velocità sono più basse.
In conclusione, i risultati sperimentali preliminari osservati non possono essere spiegati in termini di comportamento di singolo atomo e suggeriscono l’esistenza di effetti collettivi, il cui studio futuro sperimentale e teorico potrà portare alla comprensione di nuovi effetti fisici interessanti nel rinculo di atomi freddi in una trappola di dipolo creata da un intenso fascio di pompa.
Bibliography


    Science 286, 2309 (1999)


[17] N. Piovella
Appunti del corso di Fisica Atomica a.a. 2007/2008

[18] C. Foot
Atomic Physics, Oxford Univ. Press

[19] S. Chu, L. Hollberg, J. E. Bjorkholm, A. Cable, and A. Ashkin

[20] E.L. Raab, M. Prentiss, A. Cable, S. Chu, and D.E. Pritchard

[21] W. Guerin
Source atomique cohérent dans des pièges optique et magnétique: réalisation d’un laser à atomes guidé (chap.4)


[23] M.T. DePue, S. Lukman Winoto, D.J. Han and D.S. Weiss
Optics Communications 180, 73 (2000)

[24] N. Piovella
private communication


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