Cold atoms: modified radiative properties and evaporative cooling from optical traps

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Cold atoms:
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evaporative cooling from optical traps
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1 Introduction

1.1 Background

1.1.1 Optical dipole traps

Optical dipole traps (ODT) rely on the electric dipole interaction of atoms with (usually) far-detuned laser light. The interaction between the laser electric field and the laser-induced electric dipole results in a potential energy that is proportional to the laser intensity. For large detuning this potential is nearly conservative. The sign of the dipole potential can be positive or negative, depending on the color of the light, i.e., depending on whether the laser frequency is tuned above or below an atomic resonance line. Consequently, minima or maxima in the light intensity can serve as traps for cold atoms. These so-called optical dipole traps have proven to be very flexible.

The first proposal to use a focused laser beam as a trap was made in 1978 [40]. In 1986, S. Chu and coworkers succeeded in holding about 500 sodium atoms for several seconds in a tight focus of a red-detuned laser beam [41].

The far-off-resonant optical dipole trap (ODT) gained more attention after its demonstration by Miller et al. in 1993 [8], where the authors used a conventional MOT as a source for cold atoms. Since then, ODTs are widely used in the field of cold atoms.

The first evaporative cooling in a crossed beam optical dipole trap was demonstrated in 1995 [9]. In 2001 Barrett et al. demonstrated the all-optical formation of an atomic Bose-Einstein condensate [10].

ODTs do not rely on the magnetic moment of the trapped atoms. Therefore optical dipole traps opened up the possibility to trap nonmagnetic atoms, molecules, and several spin states of the same atom. These advantages of ODTs have been used in many recent experiments. ODTs are the standard way to trap species with weak magnetic moment such as ytterbium or strontium. Recently trapping and condensation of cold Yb atoms was demonstrated [20, 21] in a crossed-beam ODT. The use of ODTs also opens up the possibility to tune atomic scattering properties by a magnetic field near a Feshbach resonance. This has been demonstrated by many groups, and has led to recent success with the condensation of cesium and chromium [16–18]. Degenerate two-component Fermi gases were produced using direct evaporation in a CO$_2$ laser ODT [22].
Optical traps can also include evanescent waves (EW) as described in [18], where a highly anisotropic surface trap is realized with two evanescent light waves. Such ODT relies on a combination of a repulsive and an attractive EW field, both produced above the surface of a dielectric prism by total internal reflection of two laser beams (blue and red detuned). Such a configuration is a natural approach to build highly anisotropic, essentially two dimensional traps.

### 1.1.2 Atoms near dielectric surface

The behavior of atoms near a dielectric surface has been subject of many studies. The atom-surface interaction includes the electrostatic or Van der Waals shift, the Casimir-Polder shift, and the resonant radiative shift. The spontaneous emission rate of a multilevel atom is modified in the vicinity of a vacuum-dielectric interface [23], as a consequence of the modification of the vacuum field fluctuations in the vicinity of this interface.

Modification of spontaneous emission was first observed by Drexhage [24, 25]. Both inhibited and enhanced spontaneous emission have since then been observed by others in a variety of geometries and circumstances [28–31].

The use of cold atoms for the experimental study of quantum electrodynamics (QED) phenomena is relatively recent advantage. The situation of a particle in front of a distant mirror has recently been investigated using a single trapped ion. Both broadening of the radiative linewidth and energy level shifts have been reported for this system [26, 27]. Atom-surface interaction was recently studied in JILA [33]. Using a nearly pure BEC the authors report on measurements of a temperature dependence of the Casimir-Polder force.

### 1.2 This thesis

This thesis is devoted to the study of atom-light interaction near a surface and to trapping and cooling of atoms in a far off-resonance optical dipole trap.

A cloud of cold atoms is a good probe for studying QED phenomena due to its low temperature and its low density. We study the distance-dependent absorption linewidth of cold $^{87}$Rb atoms close to a dielectric-vacuum interface. We drop a cloud of cold atoms on the dielectric surface. We measure the absorption of EW light by cold atoms that happen to be in the EW field. An EW appears when a beam undergoes total internal reflection part a surface. This method is selectively sensitive to atoms very close to the surface [37]. These experiment allows us to observe the behavior of atoms in vacuum, the radiative linewidth of an atom has never been investigated experimentally in a vacuum at a distance of the order of an optical wavelength from a surface. We compare our experimental results to calculations based on theories of linewidth broadening and level shifts.

We also study the loading of ODTs directly from dark molasses. The direct loading of an ODT from a MOT, dark molasses or dark MOT is a common way to load an ODT. The loading of ODTs with wavelength of 10 µm proved to be rather
1.2 This thesis
efficient, and led to high atomic densities in ODTs. Various groups reported on the achievement of an all optical $^{87}$Rb BEC [10, 13, 14] by straightforward evaporation from a 10 $\mu$m ODT. However, similar experiments using 1 $\mu$m light were less successful [9]. Finally, one had to employ a compressible ODT to produce an all-optical $^{87}$Rb BEC [15] using a near infrared laser. 2 $\mu$m light was never used to produce an ODT for cold atoms. However this wavelength promised a compromise between very far detuned 10 $\mu$m, and more practical near infrared wavelengths.

Evaporative cooling from an ODT is also studied. Although ODTs are commonly used in the field of cold atoms, straightforward evaporation from an ODT is not typical. Evaporative cooling from an ODT suffers from decompression due to a decrease of the trap frequencies. In contrast to evaporative cooling from a magnetic trap, it is unknown which initial conditions are needed to reach BEC by evaporative cooling from an ODT. There are also no criteria defining an efficient evaporation strategy.

This thesis is organized in the following way. In chapter 2 we compile important theoretical expressions useful for later chapters of this thesis. We start with a brief description of atom-light interaction. Then we discuss ODTs and their possible geometries. We discuss evaporative cooling in the simplest ODTs and their difference from evaporative cooling in magnetic traps. Finally, we switch to the description of an atom near a dielectric surface. We discuss the modification of the radiative linewidth. Then we describe the shift of energy levels near a dielectric surface. We present an equation for the shift of the ground and excited levels of $^{87}$Rb atoms near a surface.

In chapter 3 various aspects of the experimental setup are described. We describe in detail the changes that are made from the previous version of the experimental setup, which is described in [34–36]. We briefly mention the diode laser setup and the vacuum system. Also the MOT, molasses and dark molasses, which were used as sources of cold atoms for ODTs, are discussed. Emphasis is put on the high-power laser setup. We briefly describe the high-power lasers. Alignment of ODTs is key element of this experiment. Trapping beams have to align with respect to MOT, and each other. Careful alignment of ODTs is needed to a good loading and high trap frequencies. The trap frequencies is an important parameter of an ODT. Two methods for the measurements of oscillation frequencies of atoms trapped in an ODT are described.

In chapter 4 we concentrate on measurements of radiative properties of cold atoms in vacuum near a dielectric surface. We describe the EW as a tool to study the atomic behavior near a dielectric surface. The optical field on the vacuum side decays exponentially with the distance $z$ to the surface, $\propto \exp(-z/\xi)$. Atoms can absorb light from the EW if their distance to the surface is on the order of the decay length $z \lesssim \xi \sim \lambda$. We adjust the decay length by changing the angle of incidence, and thus we can vary the distance scale at which the atoms interact with the probe light. We describe the experiment and the data extraction in detail. We discuss the contribution of the level shifts and the QED broadening to the overall profile of the absorption line.

In chapter 5 we describe the loading of an ODT from dark molasses. An efficient
loading is needed to reach a high initial density of atoms in an ODT. We discuss the
difference of ODTs with a wavelength of 1 µm, 2 µm and 10 µm. We describe the role
of dark molasses in our loading procedure. Then we experimentally compare ODTs
of different wavelength and various geometries. We summarize parameters such as:
the number of trapped atoms, the collision rate and the phase-space density of atoms
trapped in various ODTs. We discuss the dependence of the loading efficiency for
various parameters of the ODT.

In chapter 6 we describe evaporative cooling of cold atoms from an ODT. We
present experimental results of evaporative cooling from different ODTs. We discuss
the possibility of entering the run-away evaporation regime during evaporative cool-
ing from an ODT. One of our experimental results is that we observe a phase-space
density just above the critical. We observe an anisotropic expansion of an atomic
cloud after evaporative cooling from a crossed-beam ODT, crossed in the vertical
plane. We analyze our absorption images and we calculate the number of atoms in
the condensed fraction. Our experimental results with evaporative cooling from the
single beam ODT are compared with numerical calculations. We demonstrate the
possibility of a compressible single beam ODT. Finally, we discuss possible ways to
further improve evaporative cooling from an ODT.
In this chapter we provide a theoretical background for the thesis. In the first section we describe the interaction of atoms with light in general and the use of laser fields as optical dipole traps in particular. In the second section we describe far-off-resonance optical dipole traps and discuss their use in the field of cold atoms. We discuss the equations that describe the key parameters of an optical dipole trap. In the third section we describe the theoretical aspects of forced evaporation from an optical dipole trap. Special attention is paid to the difference between evaporation from optical dipole traps and evaporation from magnetic traps. Finally, the fourth section deals with neutral atoms near a dielectric surface. We describe the influence of a dielectric surface on the linewidth and the shift of energy levels of an atom.
2.1 Atom-light interaction

The optical dipole force is a consequence of the dispersive interaction of the induced atomic dipole moment with the intensity gradient of an external light field [38]. The force can be derived from a potential, because of its conservative character. The minima of this potential can serve as traps of atoms. The absorptive part of the dipole interaction with far-detuned light leads to residual photon scattering of the trapping light, which leads to the heating of trapped atoms, and thus, sets limits to the performance of dipole traps.

In the framework of the oscillator model [39] for a two-level atom with a resonance frequency $\omega_0$, one can derive expressions for the dipole potential and for the scattering rate. The former is associated with the interaction of the laser electric field with the induced electric dipole of the atom. The latter is associated with the scattering of photons. When an atom is placed in a laser beam with light intensity $I$, the light induces an atomic dipole moment, that oscillates at the driving frequency $\omega$. The dipole potential and the scattering rate can be calculated as [39]:

$$U_{dip}(r) = -\frac{3\pi c^2}{2\omega_0^3} \left( \frac{\Gamma}{\omega_0 - \omega} + \frac{\Gamma}{\omega_0 + \omega} \right) I(r), \quad (2.1)$$

$$\Gamma_{sc}(r) = \frac{3\pi e^2}{2\hbar \omega_0^3 \omega_0^3} \left( \frac{\Gamma}{\omega_0 - \omega} + \frac{\Gamma}{\omega_0 + \omega} \right)^2 I(r), \quad (2.2)$$

where $\hbar$ is Planck’s constant divided by $2\pi$, $c$ is the speed of light in vacuum and $\Gamma$ is the spontaneous decay rate of the excited level. If the laser is tuned relatively close to the atomic resonance at $\omega_0$ such that the detuning $\Delta = \omega - \omega_0$ fulfills $\Delta \ll \omega_0$, the counter-rotating term containing $\omega/\omega_0 \approx 1$, can be neglected. This approximation is known as the rotating-wave approximation.

In this case (which is of great practical interest), the general expressions for the dipole potential and the scattering rate simplify to

$$U_{dip}(r) = \frac{3\pi c^2}{2\omega_0^3} \frac{\Gamma}{\Delta} I(r), \quad (2.3)$$

$$\Gamma_{sc}(r) = \frac{3\pi e^2}{2\hbar \omega_0^3} \left( \frac{\omega}{\omega_0} \right)^3 \left( \frac{\Gamma}{\Delta} \right)^2 I(r). \quad (2.4)$$

These equations provide two important points for dipole trapping: (1) Below an atomic resonance (red detuning, $\Delta < 0$) the dipole potential is negative and therefore, the interaction attracts atoms into the light field. For red detuning potential minima correspond to positions with maximum intensity. Above resonance (blue detuning, $\Delta > 0$) the dipole interaction repels atoms out of the field, and potential minima correspond to minima of the intensity. So, dipole traps are usually divided into two classes, red-detuned traps and blue-detuned traps. (2) The dipole potential scales as $I/\Delta$, whereas the scattering rate scales as $I/\Delta^2$. Therefore, optical dipole traps need to use large detunings and high intensities in order to keep the scattering rate as low as possible at a certain potential depth.
2.1 Atom-light interaction

These results can be generalized for multi-level atoms. We use the scalar polarizability throughout this work, i.e. the polarizability is the same for all m-components. For a calculation of the ground-state dipole potential \( U_{\text{dip}} = \Delta E_i \), one has to sum up the contributions of all coupled excited states, taking into account the relevant line strengths \( c_{ij}^2 \) and detunings \( \Delta_{ij} \). The result of the summation can also be expressed in terms of the polarizability \( \alpha(\omega) \) of a certain atomic species for a certain frequency:

\[
\Delta E_i = \frac{3\pi c^2 \Gamma}{2\omega_0^3} I \times \sum_j \frac{c_{ij}^2}{\Delta_{ij}} = \frac{\alpha_i(\omega)}{2\varepsilon_0} I. \tag{2.5}
\]

Here the summation is carried out over all electronically excited states \( |e_j\rangle \), that are dipole-coupled with the ground state [39]. The polarizability of state \( i \) is,

\[
\alpha_i(\omega) = \frac{3\pi \varepsilon_0 \Gamma c^3}{\omega_0^3} \times \sum_j \frac{c_{ij}^2}{\Delta_{ij}}. \tag{2.6}
\]

The spatial intensity distribution of a focused Gaussian beam with power \( P \), that propagates along the \( z \)-axis is described by:

\[
I(r, z) = \frac{2P}{\pi w^2(z)} \exp\left[-\frac{2r^2}{w^2(z)}\right], \tag{2.7}
\]

where \( r \) denotes the radial coordinate. The \( 1/e^2 \) radius of the beam, \( w(z) \), depends on the axial coordinate, \( z \), via

\[
w(z) = w_0 \sqrt{1 + \left(\frac{z}{z_R}\right)^2}, \tag{2.8}
\]

where the minimum radius \( w_0 \) is called the beam waist and \( z_R = \pi w_0^2 / \lambda \) denotes the Rayleigh length.

As already mentioned the focus of a red detuned laser beam can be used as a trap for neutral atoms. The trap depth for a laser beam of power \( P \) is:

\[
U_0 = \frac{\alpha}{\pi \varepsilon_0 w_0^2} \frac{P}{w_0^2}, \tag{2.9}
\]

where \( \varepsilon_0 \) is the vacuum permittivity and \( \alpha \) is the polarizability of the ground state of the atom. An example for realistic experimental parameters with \(^{87}\text{Rb} \) atoms: for an optical power of 10 W at wavelength 1030 nm (which has been used in our experiment) and a beam waist of 50 \( \mu \)m, the polarizibility is \( \alpha = 13.6 \times 10^{-39} \text{ Cm}^2/\text{V} \) and we calculate a trap depth of 440 \( \mu \text{K} \). The corresponding photon scattering rate is 1.3 s\(^{-1} \). This trap depth is larger than the typical MOT temperature and substantially larger than the temperature which can be obtained by using sub-Doppler cooling of \(^{87}\text{Rb} \) atoms. Therefore, such a trap can be used in practice for the trapping of cold atoms. The polarizibility for a wavelength of 2 \( \mu \)m (the other important wavelength used in the experiment) is \( \alpha = 6.18 \times 10^{-39} \text{ Cm}^2/\text{V} [34], \) which is two times smaller.
Both red and blue detuned laser beams are commonly used for manipulating cold atoms. More specifically they are used to produce optical dipole traps (ODT). The most popular types of ODTs are single beam (mentioned above) and crossed beams ODTs. A single beam ODT is the easiest to create, but it has the disadvantage of weak confinement in one direction. Crossed beam ODTs consist of two focused laser beams crossed at one point. Such traps can provide a tight confinement in all directions, but they demand more delicate alignment. More complicated optical traps can include a three beam configuration, where a tightly focused beam is overlapped with a large crossed beam ODT as in [17]. It is also common to use optical lattices as part of an ODT. Optical traps can also include evanescent waves (EW) as described in [18], where a highly anisotropic surface trap is realized with two evanescent light waves. Such a trap relies on a combination of a repulsive and an attractive EW field, both produced above the surface of a dielectric prism by total internal reflection of two laser beams (blue and red detuned). Such a configuration seems to be a natural approach to create a two dimensional trap, although the prism surface roughness became an important issue [34]. Another approach for EW traps was proposed in [51], where a trap is generated by two EW beams with different polarization.

### 2.2 Evaporative cooling from optical dipole trap

Forced evaporative cooling was the decisive step in reaching Bose-Einstein condensation of cold atoms in 1995. Evaporative cooling has been used in all BEC experiments ever since.

The idea of evaporative cooling is based on the preferential removal of atoms with an energy higher than the average energy per atom (invented by H. F. Hess [42]). Subsequent thermalization by elastic collisions leads to a thermal energy distribution with a lower temperature than before the removal of the hot atoms. The main drawback of evaporative cooling is the decrease of the number of atoms in the trap. A key parameter for evaporative cooling is the ratio of the trap depth and the atomic temperature usually written as \( \eta \). The evaporation rate per atom can be written as [43]:

\[
\Gamma_{ev} = \sqrt{\frac{2}{\pi}} n_0 \sigma \bar{v} \eta e^{-\eta},
\]  

(2.10)

where \( n_0 \) is the peak atomic density and \( \sigma = 8\pi a^2 \) is the elastic collision cross section. \( a \) is the scattering length (\( a_0 \) is the Bohr radius, \( a = 99 a_0 \) for \(^{87}\)Rb). \( \bar{v} = \sqrt{k_B T/m} \) is the thermal atomic velocity (note that in Ref. [43] \( \bar{v} \) is defined differently: \( \bar{v} = \sqrt{8k_B T/(\pi m)} \), which leads to a different prefactor in this equation). The product \( \sqrt{2/\pi} n_0 \sigma \bar{v} \) is the cloud-averaged collision rate \( \gamma \). This is another key parameter for evaporative cooling.

Assuming a Boltzmann distribution for the trapped atoms the peak density can
be calculated as:

\[ n_0 = \frac{N\bar{\omega}^3}{(2\pi k_B T/m)^{3/2}}. \]  

(2.11)

With a few approximations (mentioned below), the change of the number of atoms in the trap and their temperature can be written as [44]:

\[ \dot{N} = -N\Gamma_{ev} - N\Gamma_{vac}, \]  

(2.12)

\[ \dot{T} = -\Gamma_{ev} \frac{\eta - 2}{3} T + \frac{\dot{\omega}}{\omega} T, \]  

(2.13)

where \( \Gamma_{vac} \) is the background vacuum collision rate, and \( \bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3} \) is the mean trap frequency. Note that the last term in Eq. (2.13) describes adiabatic cooling due to decompression. This term is usually absent when evaporation is performed from a magnetic trap. These equations are accurate if: 1. there is no heating (due to trap instability, light scattering etc.), 2. the \( \eta \) parameter is much larger than one, 3. an atom with an energy larger than the trap depth can leave the trap without collisions with other atoms (the opposite situation is known as the hydrodynamic regime). These equations are applicable for an atomic cloud trapped in any conservative trap, i.e. they can be equally well applied for a magnetic as for a far-off-resonance optical dipole trap.

The phase-space density of atoms of the same state in a trap can be calculated as:

\[ \rho = \frac{N}{1.2} \left( \frac{\hbar \bar{\omega}}{k_B T} \right)^3. \]  

(2.14)

The phase-space density defined here in such a way, that the BEC phase transition occurs at \( \rho = 1 \). In following chapters we use an additional factor 1/3 due to the fact, that atoms are distributed over the 3 different spin states.

### 2.2.1 Scaling laws: magnetic vs. optical dipole traps

The atomic velocity is proportional to the square root of temperature \( \bar{v} \sim \sqrt{T} \). Therefore, using Eq. (2.11), the collision rate is proportional to \( \gamma \sim N/T \). This scaling law is correct as long as the trap frequency \( \bar{\omega} \) is constant. In practice, a decrease of trap depth in combination with a constant trap frequency can be achieved by using a RF knife. A simple and obvious conclusion is: the collision rate in a magnetic trap grows if the temperature of the atoms in the trap falls. This is possible because the increase of the density beats the decrease of the atomic velocity. Therefore, in the course of evaporative cooling it might be possible to enter the so-called "runaway" regime, where, despite of the particle loss, the density of atoms increases while the collision rate \( \gamma \) remains constant or even increases. The efficiency of evaporative cooling is limited by loss mechanisms of the trap such as collisions with the background gas or three-body collisions. The figure of merit in this case is the ratio of "good" to "bad" collision rates, \( R = \gamma/\Gamma_{vac} \). The runaway regime can
be reached if the ratio $\gamma/\Gamma_{\text{vac}}$ is more than $\sim 240$. If the runaway regime is reached, a scalable solution $\rho \sim N^{\alpha}$ exists (with $\alpha$ known as efficiency of evaporation). In practice this means that if one remains in the runaway regime and keeps $\gamma$ and $\eta$ in the proper range, the BEC transition will be reached sooner or later. Evidently the final number of atoms in the BEC depends on the initial phase-space density, the initial number of atoms and the efficiency of evaporation $\alpha$.

The situation in an ODT is different, because one has to decrease the optical power and therefore the trap frequency in order to decrease the trap depth. Unfortunately, the "RF-knife" technique does not work for an ODT, since an ODT traps atoms in all magnetic states. Therefore unavoidably trap frequencies of an ODT decrease during evaporation. In the simplest case of a single beam ODT: $\bar{\omega} \sim \sqrt{U}$, assuming that gravity can be neglected. This statement is also correct for more complicated ODTs as long as the optical power of all beams is ramped down proportionally, and the geometrical configuration of the ODT is not changed.

Therefore, the expressions given above for magnetic traps, need to be modified for optical traps. Assuming that the $\eta$ parameter is constant, we obtain: $n \sim N$, and $\gamma \sim N \sqrt{T}$. Evaporation from an ODT leads to decompression (opposite to the runaway regime in a magnetic trap), and it is not possible to achieve the runaway regime in a single beam ODT (or in any ODT as long as the optical power of all beams is reduced proportionally, and the geometrical configuration is fixed). In practice, this means that one has to deal with a decrease of the collision rate, and subsequently a decrease of the efficiency of forced evaporative cooling during the evaporation from an ODT. This fact is known, and experimentally observed in [10], [20], [14], [48]. It also follows from the scalable solution for evaporation from an ODT [47]. This drawback can be compensated by good initial conditions for evaporative cooling such as a high initial collision rate and a high phase-space density [10, 14, 20, 48]. For example, M. Barrett et al. [10] report an extremely high initial collision rate of 12 kHz and phase-space density of about 1/600, while the typical parameters in a magnetic trap for $^{87}\text{Rb}$ are a collision rate of about 100 Hz and phase-space density of about $10^{-6}$.

There are several ways to avoid decompression during evaporation. One possibility is a dynamic compression of the ODT volume during evaporation [15]. In this paper, atoms are first cooled and loaded into a crossed dipole trap of 1.06 $\mu$m light. Then the density is increased by dynamically reducing the trap size, while the atoms are evaporatively cooled by reducing the light intensity. The trap size was decreased by reducing the waist of the trap beams at their crossing point by moving one of the lenses of a telescope. In this way T. Kinoshita et al. obtained a nearly pure BEC with $3.5 \times 10^5$ $^{87}\text{Rb}$ atoms.

Another way to prevent a decrease of the atomic collision rate during evaporation from an ODT is magnetic tuning of scattering properties [17], [22].

Despite a drop of the collision rate it is possible to achieve BEC by straightforward evaporative cooling from an ODT [10], [14], [48], [20], if the initial conditions are good enough i.e. a high initial ratio of good-to-bad collisions and high initial phase-space density.
2.3 Bose-Einstein condensation

We present some basic theoretical results required for the description of trapped Bose gases. Expressions for the critical temperature ($T_c$) and number of condensed atoms ($N_0$) are shown. We also briefly mention the equations for ballistic expansion of a BEC.

The density of trapped non-condensed, non-interacting atomic gas can be written as [53]:

$$n(r) = \frac{1}{\Lambda_T^3} g_{3/2}[\exp[\frac{\mu - V(r)}{k_B T}]],$$  \hspace{1cm} (2.15)

where $\mu$ is the chemical potential ($\mu < 0$ ), and $\Lambda_T$ is the thermal de Broglie wavelength, defined as:

$$\Lambda_T = \sqrt{\frac{2\pi \hbar^2}{mk_B T}}.$$  \hspace{1cm} (2.16)

The polylogarithm functions $g_a(x)$ are defined as:

$$g_a(x) \equiv \sum_{n=1}^{\infty} \frac{x^n}{n^a}.$$  \hspace{1cm} (2.17)

Far from the phase transition the fugacity $\tilde{z} = \exp[\mu/kT] \ll 1$, and the polylogarithm function can be approximated by $g_{3/2}[\exp[(\mu - V)/k_B T]] \simeq \exp[(\mu - V)/k_B T]$. The density distribution then reduces to the Boltzmann case with a central density given by Eq. (2.11).

Below a given temperature the population of the lowest state becomes macroscopic and this corresponds to the onset of Bose-Einstein condensation. This corresponds to a critical value of the phase-space density $\rho = 1$ in Eq. (2.14). For a given trap potential this can be expressed in terms of a critical number of atoms $N_c$ (i.e. atoms will start to macroscopically populate the lowest state for $N > N_c$).

For a three-dimensional harmonic potential the critical number of atoms can be written as:

$$N_c = \zeta(3) \left( \frac{k_B T}{\hbar \omega} \right)^3,$$  \hspace{1cm} (2.18)

where $\zeta(n)$ is the Riemann zeta-function, which corresponds to the polylogarithm function with its argument equal to 1: $\zeta(n) = g_n(1)$. The numerical value is $\zeta(3) = 1.202$.

Therefore the transition temperature $T_c$ for a three-dimensional harmonic-oscillator potential can be calculated as [53]:

$$k_B T_c \simeq 0.94 \hbar \omega N^{1/3}.$$  \hspace{1cm} (2.19)

The number of condensed atoms can be calculated from,

$$N - N_0 = \int_0^{\infty} \frac{\rho(\varepsilon)d\varepsilon}{\exp[\beta \varepsilon] - 1},$$  \hspace{1cm} (2.20)
The theoretical background

where \( \rho(\varepsilon) \) is the density of states, and \( \beta = 1/k_B T \).

For a three-dimensional harmonic-oscillator potential \( \rho(\varepsilon) = \varepsilon^2/(2\hbar \bar{\omega})^3 \), and the fraction of condensed atoms \( N_0 \) for a \( T < T_c \) is:

\[
N_0/N = 1 - \left( T/T_c \right)^3.
\]  

(2.21)

The previous statements (Eqs. (2.18),(2.19),) are correct in the thermodynamic limit \( (N \rightarrow \infty) \). However the number of atoms that can be put into the traps is not truly macroscopic. As a consequence, the thermodynamic limit is never fully reached. The first finite-size correction for the condensate fraction can be evaluated analytically by studying the large-\( N \) limit.

\[
\frac{N_0}{N} = 1 - \left( \frac{T}{T_c} \right)^3 - 3\omega_{tr} \zeta(2) \left( \frac{3}{2\bar{\omega} \zeta(3)^{2/3}} \right)^2 N^{-1/3}.
\]  

(2.22)

In lowest order, finite-size effects decrease as \( N^{1/3} \) and depend on the ratio of the arithmetic (\( \omega_{tr} \)) and geometric (\( \bar{\omega} \)) averages of the trap frequencies. Finite-size effects reduce the condensate fraction and thus result in a decrease of the transition temperature compared to the \( N \rightarrow \infty \) limit. One can estimate the shift of the critical temperature as: \( \delta T_c/T_c \approx -0.73 (\omega_{tr}/\bar{\omega}) N^{-1/3} \). We will neglect this term during calculations of the critical temperature, since this effect is well below 10%.

The condensate in a trapped weakly interacting gas is usually described by a macroscopic wavefunction. Its steady state is determined by the Gross-Pitaevskii equation:

\[
\left( -\frac{\hbar^2 \Delta}{2m} + U(r) + g|\Psi_0(r)|^2 \right) \Psi_0(r) = \mu \Psi_0(r),
\]  

(2.23)

where \( g = 4\pi\hbar^2 a/m \) is the coupling constant, \( a \) is the scattering length. In many cases the mean field energy term \( g|\Psi_0(r)|^2 \) becomes dominant compared to the kinetic energy term. The latter can then be neglected, leading to the Thomas-Fermi approximation. The Thomas-Fermi approximation is of great practical use, since it is well applicable to most experimentally produced BECs. The density profile in the framework of this approximation can be written as:

\[
n_c(r) = \max\left\{ \frac{\mu - U(r)}{g}, 0 \right\}.
\]  

(2.24)

For a harmonic trap this approximation leads to a parabolic shape of the density profile of the condensate, which is different from a gaussian profile of a thermal cloud.

In the Thomas-Fermi approximation the chemical potential can be calculated as:

\[
\mu = \frac{(15\hbar^2 \sqrt{mN_0 \bar{\omega}^3 a})^{2/5}}{2},
\]  

(2.25)

where again \( a \) is the scattering length, \( m \) is the mass, \( \bar{\omega} \) is the mean trap frequency.

The dynamics of the expansion of a BEC is an important issue, because much information on these Bose condensed gases is obtained experimentally from images.
of expanded atomic clouds. The Bose condensed cloud shows an anisotropy during the expansion: the condensed gas expands faster in the direction of stronger confinement. In the Thomas-Fermi regime the parabolic shape of the density is preserved during the expansion. The radii \( R_i \) \((i = \{x, y, z\})\), where the density vanishes (Eq. (2.26)), depend on time as [53]:

\[
R_i(t) = R_i(0) b_i(t) = \sqrt{\frac{2\mu}{m\omega_{0i}^2}} b_i(t),
\]

where \( \omega_{0i} \) are the trap frequencies \( \omega_{0x}, \omega_{0y}, \omega_{0z} \).

The scaling parameters \( b_i(t) \) obey a set of three coupled differential equations [53]:

\[
\ddot{b}_i + \omega_i^2 b_i - \frac{\omega_{0i}^2}{b_i b_x b_y b_z} = 0.
\]

(2.27)

For a cigar-shape strongly elongated trap (aspect ratio \( \lambda = \omega_\parallel / \omega_\perp \ll 1 \)), these equations can be solved analytically [52]. The atomic cloud expands as \( b_\perp(\tau) = \sqrt{1 + \tau^2} \) in the direction of the strong confinement, and \( b_z(\tau) = 1 + \lambda^2 (\tau \arctan[\tau] - \ln[\sqrt{1 + \tau^2}]) \) in the weakly confined direction, where \( \tau = \omega_\perp t \) is a dimensionless time (\( \omega_\perp \) is the strong frequency of the trap). Therefore the cloud of condensed atoms changes its shape from a cigar to a pancake during the expansion. The anisotropic expansion is a telltale signature of the BEC, which is widely used as an experimental proof of a condensate.

### 2.4 Atoms near a dielectric surface

An interface between vacuum and a simple dielectric medium modifies both the external and internal dynamics of a multilevel atom [23] in its vicinity. Indeed, the electromagnetic field surrounding an atom becomes modified when the atom is located close to a surface. This is even true for the vacuum fluctuations. In this way, the presence of a surface changes the natural lifetime and the energy of atomic levels. These two effects are usually discussed separately in the literature. We start with the modification of the natural lifetime.

A modification of the natural lifetime leads to a modification of the natural linewidth and of the photon scattering rate. Consider a downward transition of an atom between two levels separated by energy \( \hbar \omega_0 \), with the emission of a photon with an energy \( \hbar \omega_0 \). If the electric transition dipole moment matrix element \( \mu \) is oriented parallel to the cartesian \( j \)-axis, the spontaneous emission rate based on first-order time-dependent perturbation theory is [55]

\[
\Gamma_j = \frac{2\pi}{\hbar^2} \mu^2 \sum_f |\langle 0|\hat{E}_j(r)|f\rangle|^2 \delta(\omega - \omega_0),
\]

(2.28)

where \( \hat{E}_j(r) \) is the electric field operator, \( |0\rangle \), \( |f\rangle \) are initial and final states of the radiation field, and \( |\langle 0|\hat{E}_j(r)|f\rangle|^2 \) describe the vacuum field fluctuations. Therefore, any modifications of vacuum field fluctuations automatically leads to a modification
of the spontaneous emission rate. Generally speaking, the presence of vacuum field fluctuations provides a non-zero value for the spontaneous emission rate and leads to a finite life-time of an excited state.

The vacuum fluctuation field obeys Maxwell equations and the boundary conditions just as any other electric field. This explains the complicated behavior of the vacuum fluctuations field near an interface between vacuum and a dielectric medium, including oscillatory and evanescent wave (EW) contributions [55]. The formulas of the modified atomic linewidth close to a dielectric surface are obtained in [23] and extended in chapter 4 of this thesis.

The vicinity of a dielectric surface causes an increase of the spontaneous emission rate in a vacuum due to EW contributions from a dielectric medium.

2.4.1 Level shifts near a dielectric surface

As we show in chapter 4 the measured linewidth is also influenced by inhomogeneous broadening due to level shifts. Here we describe those level shifts in detail. The nature of the interaction of an atom with a surface depends on the distance between them. If the atom is close to the surface \((\pi z/\lambda \ll 1)\) the interaction is dominated by electrostatics, i.e., the atomic dipole experiences Coulomb forces due to the image charges at the other side of the surface, the so called Van der Waals interaction. However, as first worked out by Casimir and Polder [58], retardation becomes important for atoms further away from the surface [57].

The Van der Waals potential for an atom in vacuum at a distance \(z\) from a flat medium with index of refraction \(n\) is given by

\[
U_{vdW} = -\frac{n^2 - 1}{n^2 + 1} \frac{1}{48\pi\varepsilon_0} \frac{D^2}{z^3},
\]

which is derived in [23]. The factor \(D^2\) is the variance of the atomic electric dipole of the atomic state under consideration. An expression for \(D^2\) is derived in [23] section III.A.3. Using this, the quantum mechanical expression approaches the common expression for the Van der Waals potential for a two level atom, which is given by

\[
U_{vdW}(z) = -\frac{3}{16} \frac{n^2 - 1}{n^2 + 1} \left(\frac{1}{k_0z}\right)^3 \hbar\Gamma,
\]

to approximately 10%. For \(k_0\) and \(\Gamma\) the values of the dominant \(D_2\) transition of Rb atoms should be used (\(\lambda = 780\,\text{nm}\)). As it is visible from this equation, a characteristic distance for the van der Waals potential is \(\lambda/(2\pi) \approx 0.1\,\mu\text{m}\).

The Casimir-Polder potential is widely recognized as an example of a fully quantum electrodynamical (QED) effect in that its derivation relies on the quantization of the electromagnetic field. At distances larger than the typical Van der Waals potential distance, atom-surface separations \((2\pi z/\lambda \gg 1)\), retardation effects (due to the finite speed of light) of the virtual photons passing between the atom and its image causes the scaling behavior of the attractive force to change from \(1/z^3\) to \(1/z^4\). The Casimir-Polder potential can be written as:
\[ U_{CP}(z) = -\frac{1}{4\pi\varepsilon_0} \frac{3\alpha_0 \hbar c}{8\pi z^4} \],

where \( \alpha_0 \) is the DC polarizability of the atom. This is the original Casimir-Polder result [58].

Recently there were attempts to generalized the atom-surface interaction for multilevel atom, such as [54]. E. A. Hinds et al. provide equations, where simple, recognizable potentials emerge as a limiting case. These are the Van der Waals, Casimir, and resonant radiative interactions. The authors also provide formulas for the line shift of the lowest \( S \) and \( P \) levels of sodium near a surface.

We have extended the formulas for the lowest \( S \) and \( P \) levels for the \( ^{87}\text{Rb} \) atoms, by taking also the hyper-fine structure into account. The calculation follows Hinds and Sandoghdar [54]. The shift of a level is given by Eq. (1) in [54], assuming that the surface is a metallic mirror. In our experiment we have a dielectric surface with refractive index \( n \). Therefore we multiply the expression by \( (n^2 - 1)/(n^2 + 1) \). According to [57], this approach should be correct in the nonretarded limit, which gives the dominant contribution in our experiment.

The authors [54] perform the calculation for the lowest fine structure levels of Na: \( 3S_{1/2} \) and \( 3P_{J,M_J} \). We need to calculate the shifts for the \( ^{87}\text{Rb} \) hyperfine levels that are involved in the probe transition, i.e. \( 5S_{1/2} \) and \( 5P_{3/2}(F = 3, M_F) \). This requires a calculation of electric dipole matrix elements, starting from a given hyperfine state.

The dipole matrix elements can be decomposed in a reduced dipole matrix element \( D_{J,J'} \), Clebsch-Gordan coefficients and \( 6j \)-symbols. The reduced dipole matrix elements for the lowest transitions in Rb have been taken from Safronova et al. [56], who gives the dipole matrix elements in atomic units (which is \( e\alpha_0 \) for electric dipole moments).

Given \( D_{J,J'} \) in a.u. we can obtain the Einstein \( A \) coefficient for the transition as follows

\[ A_{J,J'} = \frac{e^2\alpha_0^2}{3\pi\varepsilon_0\hbar} \frac{D_{J,J'}^2}{2J'+1} \left( \frac{2\pi}{\lambda_{J,J'}} \right)^3 = 2.026 \times 10^{15} \frac{D_{J,J'}^2}{2J'+1} \frac{1}{\lambda_{J,J'}^3}, \]

see also Eq. (3) in [56]. If \( \lambda \) is given in nm, \( A \) is obtained in \( \text{s}^{-1} \). Note that \( J' \) is the upper level of the transition. For example, for the usual \( D \)-lines at 795 and 780 nm, the above equation gives \( A/2\pi = 5.72 \) and 6.03 MHz, respectively, which is reasonably accurate.

### 2.4.2 Ground state shift

The ground state \( 5S_{1/2} \) has only higher levels to connect to, so the shift is a summation of negative contributions. The summation runs over the upper levels \( nP_{1/2} \) (“D1-like” lines) and \( nP_{3/2} \) (“D2-like” lines) with \( n \geq 5 \). For a given level \( 5S_{1/2}(F_g, m_g) \), connected to an \( nP_{J_e} \) level, we replace \(|\langle a|d_z|j\rangle|^2\) in Eq. (1) in Hinds and Sandoghd-
dar [54] by a summation over the $F_e$ levels:

$$\frac{k_{a_j}^3}{4\pi\epsilon_0} |\langle a|d_z|j\rangle|^2 \rightarrow \frac{3\hbar A_{a_j}}{4} \sum_{F_e} \langle F_g 1 m 0|F_e m\rangle^2 (2F_g + 1)(2J_e + 1) \left\{ \begin{array}{ccc} I & J_g & F_g \\ 1 & F_e & J_e \end{array} \right\}^2.$$  \hspace{1cm} (2.33)

Depending on whether $J_e = 1/2$ or $3/2$, the range of summation for $F_e$ is 1,2 or 0,1,2,3. Similarly,

$$\frac{k_{a_j}^3}{4\pi\epsilon_0} |\langle a|d_\rho|j\rangle|^2 \rightarrow \frac{3\hbar A_{a_j}}{4} \sum_{F_e} \sum_{q=-1,1} \langle F_g 1 m q|F_e m+q\rangle^2 (2F_g+1)(2J_e+1) \left\{ \begin{array}{ccc} I & J_g & F_g \\ 1 & F_e & J_e \end{array} \right\}^2.$$  \hspace{1cm} (2.34)

Note that $I = 3/2$ is the nuclear spin of $^{87}\text{Rb}$, and the 1 comes from the photon; $\langle j_1 j_2 m_1 m_2 | j_3 m_3 \rangle$ denotes a Clebsch-Gordan coefficient, and \{\} a $6j$ symbol.

The result of the above is that all hyperfine ground states have the same shift (omitting here the refractive index correction),

$$\delta_{5S_{1/2}} = -\sum_{n=5}^{\infty} \frac{\hbar}{2} \left\{ [A(U^z + U^\rho)]_{nP_{1/2},5S} + [2A(U^z + U^\rho)]_{nP_{3/2},5S} \right\}. \hspace{1cm} (2.35)$$

This is a small extension of Eq. (15) in Hinds and Sandoghdar [54], accounting for hyperfine structure. The functions $U^z$ and $U^\rho$ contain the dependence of the level shift on the distance to the surface, as shown below.

### 2.4.3 Shift of the excited state

For the shift of the $5P_{3/2}(F = 3, m_F)$ levels, we must take into account the following transitions:

- down to $5S_{1/2}(F = 2)$,
- up to $nS_{1/2}(F = 2)$, for $n \geq 6$
- up to $nD_{3/2}(F = 2, 3)$ and $nD_{5/2}(F = 2, 3, 4)$, for $n \geq 4$

Similar to the calculation for the ground state, we decompose the dipole matrix elements in $6j$ symbols and Clebsch-Gordan coefficients.

This yields the following shifts:

$$\delta_{5P_{3/2}(F=3,m)} = \frac{1}{20} \left\{ [(18 - 2m^2)(U^z + V^z) + (6 + m^2)(U^\rho + V^\rho)] \hbar A \right\}_{nP_{3/2},5S} - \frac{1}{60} \sum_{n=6}^{\infty} \left\{ [(18 - 2m^2)U^z + (6 + m^2)U^\rho] \hbar A \right\}_{nS_{1/2},5S} - \frac{1}{300} \sum_{n=4}^{\infty} \left\{ [(9 + 4m^2)U^z + (33 - 2m^2)U^\rho] \hbar A \right\}_{nD_{3/2},5P_{3/2}} - \frac{3}{200} \sum_{n=4}^{\infty} \left\{ [(58 - 2m^2)U^z + (46 + m^2)U^\rho] \hbar A \right\}_{nD_{5/2},5P_{3/2}}. \hspace{1cm} (2.36)$$
This is an extension of Eqs. (17,18) in [54] to account for hyperfine structure. After averaging over \( m \) levels, it agrees with the Hinds and Sandoghdar results (after averaging over \( M_j \)). These are the equations that were used to plot the distance dependent level shifts in chapter 4. The summations were done up to \( n = 8 \) for the 5\( P \rightarrow nS \) transitions, up to \( n = 6 \) for the 5\( P \rightarrow nD \) transitions.

The functions \( U \) and \( V \) are defined as:

\[
U_{a_j}^z = \frac{1}{\pi} \left[ \frac{f(\phi_{a_j})}{\phi_{a_j}^3} + \frac{g(\phi_{a_j})}{\phi_{a_j}^2} \right],
\]

\[
U_{a_j}^\rho = \frac{1}{\pi} \left[ \frac{f(\phi_{a_j})}{\phi_{a_j}^3} + \frac{g(\phi_{a_j}) + 1}{\phi_{a_j}^2} - \frac{f(\phi_{a_j})}{\phi_{a_j}} \right],
\]

\[
V_{a_j}^z = \left[ -\frac{\cos(\phi_{a_j})}{\phi_{a_j}^3} - \frac{\sin(\phi_{a_j})}{\phi_{a_j}^2} \right],
\]

\[
V_{a_j}^\rho = \left[ -\frac{\cos(\phi_{a_j})}{\phi_{a_j}^3} - \frac{\sin(\phi_{a_j})}{\phi_{a_j}^2} + \frac{\cos(\phi_{a_j})}{\phi_{a_j}} \right],
\]

(2.37)

where \( \phi_{a_j} = 2k_{a_j}z \) and \( z \) is the atom-mirror distance. The functions \( f \) and \( g \) can be expressed as integrals:

\[
f(x) = \int_0^\infty \frac{\exp(-xt)}{1 + t^2} dt
\]

\[
g(x) = \int_0^\infty \frac{t \exp(-xt)}{1 + t^2} dt,
\]

(2.38)

Although these equations are rather complicated, the calculation of the level shift is straightforward. The above given potentials describe the interaction of an atom and a dielectric surface. In addition, since the shift of the ground and excited level are not identical, the level shifts also lead to a shift of the atomic resonance near a dielectric surface. Therefore the level shift pronounces itself as additional broadening near the surface, if one measures an absorption linewidth over an interval of distance from the surface. We use these results in chapter 4 for the calculated linewidth.
We discuss the most important components of our experimental setup. A substantial part of the used setup is already described [34–36]. The most important extension of the setup is the installation of two high-power lasers: a diode-pumped Yb:YAG Thin-Disk Laser (ELS, VersaDisk) and a Thulium Fiber Laser (IPG, TLM-50-2000-LP) working at wavelengths of 1 µm and 2 µm, respectively. Far off resonance optical dipole traps have been realized with these lasers. We converted to a tapered amplifier system as a source for the MOT light. An additional repumper was added and aligned to the MOT. We also mention improvements of the vacuum system.
3.1 Introduction

The experimental setup is distributed over two optical tables. All diode lasers are situated on an auxiliary table from where light is transported by single mode (SM) polarization maintaining optical fibers to the main table, which supports the vacuum setup, optics for the MOT, and lasers and optics for the optical dipole trap (ODT). The present experimental setup is based on an earlier one. The original setup described in [34], was used for the observation of modified radiative properties of cold atoms, which are reported in chapter 4 of this thesis. This previous setup had already some components that were required for an optical trapping experiment, for example, good optical access to the MOT chamber. However, some parts of the setup did not meet the new demands (such as a good vacuum, which defines the lifetime of trapped atoms). We have chosen for a few relatively fast but crucial changes to adapt our old experimental machine to the new demands.

In this chapter, we will discuss the high-power tapered amplifier (TA) system, which has been used for all MOT beams. We will describe a technique of frequency switching a locked laser, which is similar to the approach described in [64]. This technique helps us to change the frequency of the locked master laser by a substantial value (about 100 MHz), while only moderately changing the frequency of the acousto optical modulator (AOM) or even without any change of the AOM frequency.

We discuss a modification of the MOT setup and the installation of the dark repumper beam.

Finally, we describe the high-power lasers, that operate at wavelengths of 1 µm and 2 µm. The advantage of the application of these lasers, compared to near-infrared diode lasers is that the resulting potential is close to conservative. A drawback of the large detuning is the required high power. The frequencies of these lasers do not need to be set accurately, since they are far-off resonance. However, the linewidth of the output has direct consequences for the time dependence, in particular the stability of the light. The high power of these lasers needs to be handled in a proper way.

3.2 Diode laser park

Here we describe the used semiconductor diode lasers. We use our master diode lasers in external grating Littrow configuration. Use of the optical feedback by the first diffraction order from a grating allows us to decrease the laser linewidth. The frequencies of the external grating master laser can be accurately set, which is important for the MOT, the probe, as well as for the repumper beams. We do this in the usual way by frequency locking several diode lasers to the Rb transitions.
3.2 Diode laser park

3.2.1 Frequency locked diode lasers: the MOT, probe and repumper beams

There are three diode lasers locked to a $^{87}$Rb spectral feature, two of them to the $F = 2 \rightarrow F' = (1, 3)$ cross-over of the $D_2$ line (780 nm), and a third to the $F = 1 \rightarrow F' = 2$ transition of the $D_1$ line (795 nm). All three lasers are 50 mW diode lasers with external grating feedback [59, 60] in Littrow configuration [61].

In order to frequency lock these diode lasers to a spectral feature we perform FM spectroscopy. A small RF signal (20, 33, 40 MHz) is added to the laser current, which results in frequency sidebands on the laser light. A small fraction of the available optical power is used to obtain a Doppler-free saturation spectroscopy signal. The spectroscopy error signal is obtained by mixing the photodiode signal with the output of an RF oscillator [62, 63]. The resulting signal is used as feedback to the laser current in order to compensate for frequency fluctuations. Furthermore this error signal is integrated and sent to the piezo on which the grating is mounted, in order to compensate for long term drift. A typical linewidth of such a laser system is about 1 MHz. This is described in more detail in the next chapter.

Two master lasers are locked to the $D_2$ line. They are used as a tapered amplifier (TA) master and as a probe master. Although it is unnecessary to use two different master lasers for the probe and the MOT laser, we find this scheme more flexible. The probe master laser has been described in [34] and has not been changed ever since. The master for the TA will be discussed in the next sub-section.

The laser that is locked to the $D_1$ line is called the repumper, and is used to recycle atoms that fall into the $F = 1$ ground state during the MOT and molasses stages of the experiment. The output of the repumper laser is split in two parts, and coupled into two optical fibers. We were able to couple about 70 percent of the optical power into an optical fiber. A total power up to 3 mW can be obtained behind each of the SM optical fibers. All beams can be switched on and off independently within 100 $\mu$s by mechanical shutters (Vincent Associates, Uniblitz, LS2T2). The purpose of the two repumper beams is discussed later in chapter 5 in connection with the ODTs.

3.2.2 Master laser for the tapered amplifier: a power-efficient switching scheme

A TA laser system provides a simple way to combine a high frequency precision, a narrow linewidth ($\leq 1$ MHz) and a relatively high power (up to 500 mW) in an economical way. The master laser accounts for the high frequency resolution and the narrow linewidth, the TA issues the high power. A proper injection lock is the way to combine these advantages.

Generally, our MOT laser setup looks as follows. We split the beam of the external grating master laser into two parts. About 80 percent of the master light is used for injection locking of the TA chip and about 20 percent is used for spectroscopy in a Rb vapor cell.

The spectroscopy beam goes to a double pass AOM setup, and then to a Doppler-
Figure 3.1: Schematic overview of the TA laser system used for the MOT. The following abbreviations are used: L - lenses, M - mirrors, AOM - acousto-optic modulator, EGDL - external grating diode laser, NBS - nonpolarizing beam splitter, PBS - polarizing beam splitter, \( \lambda/4 \) - quarter-wave plate, OI - optical isolator, TA - tapered amplifier, PMF - polarization maintaining fiber, Rb cell - glass cell with rubidium vapor. The dashed box comprises a commercial TOptica system.

Experimental setup

free spectroscopy section (see Fig. 3.1). Our AOM can produce a frequency shift of the optical frequency varying from \( 2 \times 48 \text{ MHz} \) to \( 2 \times 105 \text{ MHz} \). However the use of AOM over its full range is difficult due to a frequency dependent AOM efficiency for a different frequency shift. Thus, it is hard to keep the laser locked during this change. We need to tune the MOT light from \( 2 \Gamma \) to \( 15 \Gamma \) below the \( F = 2 \rightarrow F' = 3 \) transition of the \( D_2 \) line (\( \Gamma \) is a natural linewidth \( \Gamma = 6.07 \text{ MHz} \)). This problem can be solved by buying AOM and RF drivers for a wider frequency range, but this solution has the drawback that one has to lock the master laser on the \( F=2 \rightarrow F'=2 \) transition, which has a relatively small intensity.

A practical solution is to send a feedforward signal to the modulation input of the laser. This feedforward signal leads to a current change, and subsequently to a frequency change in the master laser. The value of this signal is carefully adjusted to move the laser frequency from the \( F = 2 \rightarrow F' = (2, 3) \) cross-over to the \( F = 2 \rightarrow F' = (1, 3) \) cross-over. If this happens fast enough, the locking electronics do not notice that the locking transition has been changed. This approach gives us a frequency change of 78 MHz, so that the detuning becomes 15 \( \Gamma \) with respect to the MOT transition (together with the initial MOT detuning of \( \sim 2 \Gamma \)). Note that the AOM frequency remains the same during this procedure. Thus, we do not have a problem with frequency locking. In the following we will call this AOM-free frequency shifting. This AOM-free frequency shifting is one of the way to produce a large shift of a laser frequency. An obvious drawback of this AOM-free frequency switching technique is that one can shift the frequency only in fixed intervals, which are define by the Rb spectrum. However, by combing the AOM-free technique with
normal power-efficient frequency change [64], one can reach a substantial frequency interval.

3.3 A MOT and its variations

3.3.1 A MOT and molasses

Most of the MOT setup is similar to the one described in [34]. However we made a few essential improvements. We used the high power from the TA system to improve our MOT. We obtain up to 140 mW of MOT light after the single-mode (SM) optical fiber. This value can be increased by simply increasing the TA chip current, but this is not necessary in our case. We enlarged the MOT beams to diameter of 11 mm. A further beam increase is restricted by the diameter of our MOT optics.

By this straightforward improvement combined with a proper beam alignment, we increased the number of atoms in the MOT up to $100 \times 10^6$ (compared to $20 \times 10^6$ in [34]).

The MOT is loaded from background $^{87}$Rb gas from the Rb dispenser, which is on during 10 s. When the dispenser is on, the current through the dispenser is 4 A. During the loading stage two coils in anti-Helmholtz configuration produce a magnetic field gradient of approximately 12 G/cm. The 6 MOT beams are circularly polarized; the beams that propagate perpendicular to the coil axis are orthogonally polarized with respect to the beams that propagate parallel to the coil axis. The repumper light shines directly into the MOT region.

When the MOT loading is saturated, we switch off the dispenser. The temperature of the trapped cloud of atoms can be further decreased by a 3 ms short period of polarization-gradient cooling. During this stage the magnetic field coils are switched off and the detuning of the MOT light is increased to $\delta \approx -15 \Gamma$ with respect to the $F = 2 \rightarrow F' = 3$ transition of the $D_2$ line. We realize a temperature of 5 $\mu$K after this molasses stage. This was used in the QED measurements, described in chapter 4.

3.3.2 A MOT and a dark molasses

Unfortunately, a MOT or optical molasses do not form a good loading source for an ODT. The reason is photon rescattering; an issue which is discussed in details in the literature [10, 14].

To efficiently load an ODT, we use a dark MOT procedure. The MOT detuning is increased to $\delta \approx -15 \Gamma$ and one of the repumper beams is extinguished (we will called this the main repumper beam), while the second repumper beam (which we will call the dark repumper beam), is left on. The dark repumper beam is called dark, because it has a dark region in the center. The beam is passed through a glass plate with a dark spot on it. This spot is formed by a square piece of black paper, approximately 1 mm in size, which has been glued on the glass plate. This dark spot is imaged into the vacuum cell using a relay telescope. The task of the relay
telescope is to produce an image with minimal diffraction fringes. The image of the dark spot is spatially overlapped with the center of the MOT.

![Figure 3.2: The dark repumper optical setup. In the drawing the following abbreviations are used: PMF - polarization maintaining fiber, L - lens with a focal length \( f = 20 \) cm, GP - glass plate with a dark spot.](image)

In the shadow formed by the dark spot, the atoms do not scatter repumper light, and thus, by the optical pumping, they gradually fall into the \( F=1 \) state. Due to suppression of photon rescattering the atoms can temporarily reach a higher density. Unfortunately, the dark MOT loses its atoms rather fast, because the atoms spend most of the time in the \( F = 1 \) state, where they do not interact with the MOT light. We measure a decay of the number of trapped atoms close to exponential, with a \((1/e)\) decay time on the order of 50 ms. In general, the central peak density increases during the first 10 ms, after which it decreases. We measure an increase of the peak density to \(3 \times 10^{10} \) cm\(^{-3}\), which is approximately a factor of 2 higher than in a normal MOT. Although the atomic density does not increase dramatically, the dark MOT has a huge effect on the loading of an ODT. It improves the loading of atoms into an ODT by a factor of five, compared to direct loading from molasses. We did not succeed to reach the dark MOT densities reported in the literature [65]. A possible, in fact a probable reason for this is formed by the initial number of the atoms in the MOT. The mentioned authors use a Zeeman slower. Consequently, they have a much higher initial atomic number. Therefore, atoms in the MOT are deeply in the multiple scattering regime, which is believed to be a crucial criterion for an efficient dark MOT [65]. Nevertheless, the dark MOT phase is necessary for us, to reach a high initial density in the ODT, as we will discuss later, in chapter 5.

### 3.3.3 Probe beam

The probe laser diode is injection locked directly with light from the probe master laser. It is frequency shifted close to the \( F = 2 \rightarrow F' = 3 \) transition of the \( D_2 \) line by a double-pass AOM setup and is subsequently coupled into a SM fiber. It is used for all probing purposes in the experiment: e.g. the evanescent-wave probing technique that will be discussed in chapter 4, and for absorption imaging [35, 36]. After the SM polarization maintaining optical fiber, the probe beam is expanded to a diameter of about 1 cm. The power of the probe is about 0.2 mW, which corresponds to a saturation parameter of about 0.1. We use a relay telescope to image an absorption image from the vacuum cell on a CCD array (more details...
The relay telescope consists of 2 lenses with different focal length (10 cm and 15 cm), yielding a magnification $\sim 1.5$. At the position of the cold atoms a pixel corresponds to 9 $\mu m$, which seems a good compromise for us between resolution and signal-to-noise ratio.

### 3.3.4 The vacuum cell

Compared to [34] we introduced a few essential changes to the vacuum system: we added a titanium sublimation (TS) pump, and removed the prism from the vacuum chamber. This prism was used to modify the vacuum radiation field as described in chapter 4. In addition we attached the Rb dispenser on a boron nitride block to increase the thermal conductivity between the dispenser and the rest of the world. In this way, the dispenser can be switched on and off faster. This straightforward improvement leads to a drop of the measured base pressure from $3 \times 10^{-10}$ to $4 \times 10^{-11}$ mbar. However, the ion gauge is likely to underestimate the pressure in the glass cell, due to the large distance between cell and gauge head. Ultimately the figure of merit is the observed trap life-time, which is about 8 seconds.

### 3.4 High power lasers

Here we describe the high power lasers with 1 $\mu m$ and 2 $\mu m$ wavelength. In the field of cold atoms high power lasers are commonly used to produce trapping beams. However, the use of 2 $\mu m$ light for an ODT is completely new. Two points are in particular relevant on the experiment: one needs to control the power of the trap beams, i.e. potential depth as a function of time, and to align the trap beams with respect to the MOT and each other.

#### 3.4.1 Diode-pumped Yb: YAG Thin-Disk Laser.

Our Yb:YAG laser produces over 30 W of linearly polarized light (extinction $> 100 : 1$), in a beam characterized by $M^2 = 1.05$. This laser was produced by ELS Elektronik Laser System GmbH. We use the Versadisk-1030-30-SF model. The noise is specified to be $< 0.04 \%$ rms (20 Hz to 100 MHz). However we have measured a substantially higher noise level, typically a few percent of the total power. This is a single frequency laser, with a linewidth $< 2$ GHz, operating at a wavelength of 1030 nm. The characteristic frequency of the laser noise was measured to be in the kHz range, which is close to the typical frequency of an ODT. We changed the laser cavity and removed an etalon from the cavity in order to reduce this noise. As a result we achieve a substantial decrease of the noise in the kHz range. However this led to a multi-mode behavior of this laser. The 1030 nm wavelength is relatively popular, and there is no problem to obtain optical components compatible with this wavelength.
3.4.2 Thulium Fiber Laser

The Thulium Fiber Laser model TLM-50-2000-LP produces about 50 W of linearly polarized light, in a beam characterized by $M^2 < 1.1$. The polarization extinction ratio is higher than 20. This laser was made by IPG Photonics. The wavelength is 2000 nm. The linewdith of the laser is specified as less than 1 nm. The output power instability is less than 1 percent in the range from 1 kHz to 1 MHz. Actually we found that the Thulium Fiber Laser is much more quiet than the VersaDisk laser. The wavelength of 2000 nm is the maximum wavelength where normal optical materials, such as quartz are still transparent.

3.4.3 Setup of an optical dipole trap

A laser beam can create an attractive potential for an atom, if the light field is tuned below the atomic transition frequency (red detuning). Therefore, the focus of a laser beam can be used as a trap for neutral atoms. An important parameter of a such trap is the scattering rate of light by the atoms, because scattering of trap light will heat the atoms in the trap. The trap potential scales as $I/\delta$, whereas the scattering rate scales as $I/\delta^2$. Therefore, one has to use a large detuning and a high power to obtain a conservative but still deep potential. This defines the choice of the laser for our ODTs for cold $^{87}$Rb atoms.

In order to install and use an ODT, one needs: 1) to control the optical power, 2) to focus the beam tightly enough, 3) to align the focus of the beam with the center of the MOT.

A first and surprising problem, which one meets with this type of lasers, is how to extinguish the beam. Normal shutters (Vincent Associates, Uniblitz) are not useful, since these shutters are not able to hold the high power of the beams. One can use the first order reflection from an AOM, but then a substantial fraction of the power will be wasted. Another solution is to use shutters from NM Laser Products, Inc. which are supposed to handle high power beams. A practical disadvantage is that the shutters become hot (more than 100 ºC), and sometimes get stuck (usually for powers above 40 W).

We use an EOM in combination with a Glan Laser Polarizer (PGT8610) to control the optical power. Of course, the EOM can not extinguish a beam completely. We achieve an extinction ratio for the EOM of about 1/400, which is certainly not enough. Thus the use of additional mechanical shutters is necessary. We use a Pockels cell to adjust the 2 µm beam power. The Pockels cell can rotate the polarization only over $\lambda/4$, thus we have to use it in a double pass geometry (see Fig. 3.3).

After power regulation the beams are expanded, collimated and focused into the MOT. The diffraction limit for a Gaussian beam waist is:

$$w_0 = \frac{\lambda f}{\pi w},$$

(3.1)

where $\lambda$ is the wavelength, $f$ the focal distance of the lens and $w$ the diameter of the collimated beam. In order to produce a tight focus, one needs to make the diameter of the beam rather large, and to use a lens with a short focal distance.
We expand both the 1 µm and 2 µm beam up to 10 mm and use lenses with focal lengths in the range 10-20 cm. In practice, we never reached the diffraction limit, but still we are able to focus the beams to waists of 20-30 µm. To approach the diffraction limit, one needs to take care, that a trap beam is going through the center and perpendicular to the lenses. The realized waists allow to have deep optical traps (a few hundred µK) with a realistic beam power (about 10 W or less).

It is crucial for efficient loading of an ODT to spatially overlap the focus of the trap beam with the center of the MOT. There is no way to see a far-off-resonant beam in a vacuum chamber, unless some cold atoms are trapped. Therefore we overlap a resonant probe beam with the trap beam. When we are sure that the probe and the trap beam are accurately co-propagating, we align the resonant probe to the MOT center. Resonant light destroys the MOT easily. With the help of this normally unwanted effect we carefully overlapped the trap beam with respect to the center of the MOT. Unfortunately, it is difficult to find the position of the laser focus with respect to the MOT. An obvious way is to trap MOT atoms. Trapped atoms will naturally move to the potential minimum, i.e. to the beam focus. Thus, it is a good idea to put the last lens, which focuses the trap beams, on a translation stage (which allows for freedom in the direction of the trapping beam).

![Figure 3.3: Schematic overview of the 2 µm laser optical setup. LH - laser head, HPS - high-power shutter, PC - Pockels cell, L - lenses, M - mirrors, HWP - half-wave plate, GLP - Glan-Laser polarizer, MOT - magneto-optical trap (actual target).](image_url)

### 3.4.4 Waist of a trapping beam

One needs to measure the beam waist, since the potential depth of an ODT is inversely proportional to the square of the waist. The most straightforward way is to use a razor blade on a translation stage. We move a blade across the beam and measure the transmitted optical power, as a function of the blade positions. Using
this last method we find a beam waist by fitting with the function

\[ P(x) \sim w \left\{ 1 + \operatorname{Erf} \left[ \frac{\sqrt{2}}{w} (x - \xi_0) \right] \right\}. \]  

\[ (3.2) \]

**Figure 3.4:** Scan of a razor blade through the 2 \( \mu m \) beam (upper picture). We measure the transmitted optical power and fit it in order to get a result for the waist. Points are experimental data; the solid line is fit function (2). The waist is \( w = 43.7 \mu m \). This procedure is repeated for different positions of the blade along the beam and the measured beam waist vs. blade position is obtained (lower picture). The line is a fit of equation (3) through the data points. The Rayleigh length is \( z_R = 1.96 \text{ mm} \).

An example of such a waist measurements is shown in Fig. 3.4. Unfortunately one can not perform these measurements inside the vacuum chamber. Thus we placed a mirror between the last lens and the vacuum chamber and made this measurement outside. In order to simulate a potential influence of the quartz cell on the beam waist, we inserted a spare cell to mimic it. We did not measure any substantial difference with and without the spare cell. A drawback of this procedure
3.4 High power lasers

is that it is not possible to insert the mirror between the vacuum chamber and the last lens, if the last lens is too close to the vacuum chamber. However this position is necessary to achieve a tight focus of the beam. In this case one has to perform trap frequency measurements directly. We discuss this in the last section of this chapter.

We find that the measured waist can substantially differ from the calculated diffraction limit. Since an exact position of the focus is hard to determine, it is convenient to employ an additional translation stage along the beam (in the longitudinal direction), and to measure a set of diameters. Thus one can measure a Rayleigh length $z_R$ using:

$$w(x) = w_0 \sqrt{1 + \left(\frac{x - x_0}{z_R}\right)^2}. \tag{3.3}$$

In Fig. 4 we present measurements of the beam diameters based on the fits, and we show measurements of the fitted waists depending on the longitudinal translation stage position.

The measured Rayleigh length (1.96 mm) is close to the calculated Rayleigh length (1.92 mm) (obtained as: $z_R = \pi w_0^2/\lambda$).

In this way, we measure both the waist and the Rayleigh length of the beam. Both measured parameters can show substantial differences from calculations, if the beam was not aligned properly.

Using the values of a waist $w_0$ and an optical power $P$ of the trapping beam the trap frequencies for a single beam ODT can be calculated as (approximating a harmonic potential near the bottom of the trap):

$$\omega_r = \sqrt{\frac{4\alpha P}{\pi c \varepsilon_0 m w_0^4}}, \tag{3.4}$$

$$\omega_z = \sqrt{\frac{2\alpha P}{\pi c \varepsilon_0 m z_R^2 w_0^2}}, \tag{3.5}$$

where $\omega_z$ and $\omega_r$ are frequencies perpendicular to and along the trap beam propagation, respectively. $\alpha$ is the polarizability for a given wavelength, $c$ and $\varepsilon_0$ are the speed of light and vacuum permittivity.

3.4.5 Measuring trap frequencies of an ODT

Trap frequencies are crucial parameters for calculation of the density, collision rate, and phase-space density of atoms in a trap (we assume that we know the temperature and the number of atoms, since these two are relatively easy to measure). In general, we know two methods to measure trap frequencies directly: 1) observation of trap losses by modulating the trap depth with a varying frequency, 2) observation of the oscillation of an atomic cloud position or shape resulting from a sudden change of trap depth i.e. a change of optical power.
We spent a long time trying to measure trap frequencies by means of the first method: resonant atomic losses. We modulated the EOM/Pockels cell input voltage using a function generator. This leads to modulation of the optical power and therefore of the trap depth and the trap frequencies. If the modulation frequency is close to one of the trap frequencies ($\nu_{\text{trap}}$) multiplied by a factor of two, the trap starts to lose atoms due to parametric heating ($\nu_{\text{res}} = 2\nu_{\text{trap}}$) [66].

This method has two basic problems: anharmonicity and overtones. Due to its finite depth an ODT is not harmonic. The anharmonicity is more pronounced if the atomic temperature is comparable to the trap depth. Anharmonicity shows itself in resonance broadening. Secondly, atomic losses happen at frequencies $\nu_{\text{res}}/n$, where $n$ is any integer. Therefore it is advisable to measure atomic losses in a single beam dipole trap, so that one does not mix the trap frequency with overtones of a higher trap frequency. A typical result of trap losses vs. modulation frequency is shown in Fig. 3.5. Measurements were done using a 2 $\mu$m single beam ODT. The optical power of the beam was about 3.7 W. The measured frequency (0.98 kHz) corresponds to a beam waist of a 37 $\mu$m, which is confirmed by a direct measurements of the waist by the blade method ($\sim 37.5$ $\mu$m).

The second way to measure trap frequencies is to observe directly the oscillations of an atomic cloud. One can either observe an atomic cloud in situ (thus, one hopes to observe space or shape oscillations), or take a picture of an expanded cloud (to observe velocity oscillations). In practice, it is hard to measure either shape, or velocity oscillations for an optically trapped atomic cloud. The atomic cloud is small, on the order of the beam waist. In our case this corresponds to a few pixels. The atomic cloud contains a small number of atoms, so that an expanded cloud barely absorbs any light. This method suffers from poor signal-to-noise ratio.

We succeeded only in measuring the frequencies of the longitudinal oscillations of
3.4 High power lasers

Figure 3.6: Oscillation of the axial cloud size. Measurements were done using a 1 \( \mu \)m single beam ODT. The initial optical power of the beam was 0.6 W. The trap depth was slowly ramped down during 1.5 sec to \( \sim 1/3 \) of the initial trap depth to evaporatively cool the sample. Then (at \( t=0 \)) the trap depth was quickly restored to the initial value. This leads to oscillations of the atomic cloud size. The fitted frequency is 119 Hz, which corresponds to a trap frequency \( \omega_z/2\pi = 60 \) Hz.

Figure 3.7: Center of mass position vs. time under the same conditions as Fig. 7. The fitted frequency is \( \omega_z/2\pi = 59 \) Hz.

a single beam optical trap. The optical power was ramped down near exponentially, so atoms were actually evaporated. After this ramp we immediately switch the power to its initial level. We let the atomic cloud oscillate in the trap for a certain time and finally, switched the trap beam completely off and took an absorption picture of the atomic cloud. We observe shape and center-of-mass oscillations as shown in Fig. 3.6 and Fig. 3.7.

These obtained axial frequencies (60 and 59 Hz, from a breathing and a center of mass oscillations respectively) are reasonably close to the axial frequency of 63 Hz calculated using Eq.(3.5). The opposite calculation is also possible: we can estimate the beam waist \( w_0 \) and transverse trap frequency using the fitted longitudinal
frequency and formula: $z_R = \pi w_0^2/\lambda$. We have to mention, that this method gives correct results only for a Gaussian beam. However as this example shows us, this calculation gives reasonably consistent results.
Observation of modified radiative properties of cold atoms in vacuum near a dielectric surface

We have observed a distance-dependent absorption linewidth of cold $^{87}\text{Rb}$ atoms close to a dielectric-vacuum interface. This is the first observation of modified radiative properties in vacuum near a dielectric surface. A cloud of cold atoms was prepared using a magneto-optical trap (MOT) and optical molasses cooling. Evanescent waves (EW) were used to observe the behavior of the atoms near the surface. We observed an increase of the absorption linewidth with up to 25% with respect to the free-space value. We performed a numerical integration of the optical Bloch equations OBE. This method takes into account the transient internal dynamics and its possible effects on the observed linewidth. Approximately half the broadening can be explained by cavity-quantum electrodynamics (CQED) as an increase of the natural linewidth and inhomogeneous broadening. The remainder we attribute to local Stark shifts near the surface. By varying the characteristic EW length we have observed a distance dependence characteristic for CQED.
4.1 Introduction

An electronically excited atom (or molecule) can decay to the ground state by spontaneous emission. The characteristic rate at which this occurs is not simply an intrinsic property of the atom but also depends on the environment. The spontaneous emission rate is proportional to the density of electromagnetic field modes (DOS, or “density of states”), which is determined by the electromagnetic boundary conditions. The DOS can thus be modified, and with it the spontaneous emission rate. The boundary conditions imposed by the environment not only change the radiative linewidth but also induce energy level shifts and thus change the transition frequencies. These include the electrostatic or Van der Waals shift, the Casimir-Polder shift (modification of the Lamb shift), and resonant radiative shifts. For a review see, e.g. Ref. [67].

Modified spontaneous emission was first observed by Drexhage [24, 25], using dye monolayers separated from an interface by fatty acid layers. Both inhibited and enhanced spontaneous emission have since then been observed by others in a variety of geometries and circumstances [28–31]. Remarkably, the radiative linewidth of an atom in vacuo at a distance of the order of an optical wavelength from a dielectric surface has never been investigated experimentally. Energy level shifts have been studied for atom inside cavities [68, 69] and in vapor cells, using selective reflection spectroscopy [70, 71]. The situation of an atom in front of a distant mirror has recently been investigated using a single trapped ion. Both the broadening of the radiative linewidth and energy level shifts have been reported for this system [26, 27].

4.2 Method: evanescent-wave spectroscopy

The radiative linewidth $\Gamma$ is proportional to the power spectral density of the vacuum field fluctuations at the position of the atom [55], i.e. the local DOS. The proximity of a dielectric surface imposes a boundary condition on the field, changing the DOS. This leads to a modification of $\Gamma$ and to energy level shifts [23, 31, 54, 55, 70–72]. Both the linewidth broadening and the level shifts are significant mainly at distances $z \lesssim \lambda/2\pi$, where $\lambda$ is the wavelength of the dominant electronic transition. In our case this is the $D_2$ resonance line of Rb, and $\lambda/2\pi = 124$ nm.

Therefore we probe the cold atoms near the glass surface using evanescent-wave (EW) spectroscopy [37]. This method is selectively sensitive to atoms very close to the surface. An EW appears when our probe beam undergoes total internal reflection at the glass surface with index of refraction $n = 1.51$, see Fig. 4.1. The optical field on the vacuum side decays exponentially with the distance $z$ to the surface, $E(z) \propto \exp(-z/\xi)$. Atoms can absorb light from the EW, if their distance to the surface is on the order of the decay length $z \lesssim \xi \sim \lambda$. The decay length can be adjusted by changing the angle of incidence $\theta$ according to $\xi(\theta) = (\lambda/2\pi)(n^2 \sin^2 \theta - 1)^{-1/2}$. By adjusting $\theta$ we can thus vary the distance scale at which the atoms interact with the probe light. By increasing $\theta$ further above the critical angle $\theta_c = \arcsin n^{-1}$, the absorption will occur closer to the surface, where $\Gamma$ is more strongly modified.
4.3 Experiment

For atoms in free space, the absorption profile is given by a Lorentzian profile, centered at the (free-space) atomic transition frequency $\omega_{eg} = \frac{2\pi c}{\lambda}$ and with a full width at half maximum (FWHM) equal to the natural linewidth $\Gamma_\infty / 2\pi = 6.07 \text{ MHz}$ for the $D_2$ resonance line of $^{87}\text{Rb}$. Both $\omega_{eg}$ and $\Gamma_\infty$ change in the proximity of the surface. Roughly speaking these $z$-dependent Lorentzians become convoluted with the EW energy density $U(z) \propto \exp(-2z/\xi)$. Therefore, in our experiment $\omega_{eg}(z)$ also contributes to the observed absorption linewidth through inhomogeneous broadening. We expect the width of the resulting absorption profile to increase with the angle of incidence $\theta$. In the experiment we measured this by tuning an EW probe laser across the profile and measuring the absorption.

Figure 4.1: Scheme of the experiment. A weak evanescent-wave probe beam is reflected from the prism surface and collected on a photodiode. The reference beam has equal power but bypasses the vacuum cell. The difference photocurrent of probe and reference beams yields the absorption signal, typically a fraction of $10^{-3} - 10^{-4}$ of the probe. A second reference beam (not shown) was used to monitor variations in the probe power. We also collect the fluorescence from the MOT to normalize for shot to shot variations of the number of atoms.

The major part of our experimental setup has been described previously [73]. We produced clouds of cold $^{87}\text{Rb}$ atoms using magneto-optical trapping inside a ultra high vacuum cell (base pressure $p \approx 10^{-10}\text{ mbar}$). After postcooling in optical molasses we ended up with about $3 \times 10^7$ atoms, at a temperature of 9 $\mu\text{K}$. At this temperature the Doppler width is 90 kHz (FWHM). The cooling lasers were then switched off and the atoms fell down toward the surface of a glass prism, about 3.6 mm below. The center of the cloud reached the prism surface and the EW spot after 27 ms. Just before hitting the surface, the atoms briefly interact with a weak, $p$-polarized, EW probe beam, see Fig. 4.1. The intensity of the probe was kept well
below the saturation intensity to avoid power broadening. Using 0.35 µW and a waist of about 1 mm, the maximum saturation parameter was $s \simeq 0.08$.

The probe beam was derived from a home-built diode laser system, locked to the $F = 2 \rightarrow F' = (1,3)$ crossover resonance in the $D_2$ line of $^{87}\text{Rb}$ (780 nm). We used an acousto-optic modulator (AOM) to shift the probe frequency near resonance with the $F = 2 \rightarrow F' = 3$ transition and to tune it across the resonance. Before sending it into the cell, a fraction of the probe beam was split off and sent to a photodiode as a reference. After total internal reflection on the prism surface the probe beam was focused on a second photodiode. The photocurrents of the two photodiodes were subtracted to obtain our signal, typically a fraction of $10^{-3} - 10^{-4}$ of the probe.

The difference photocurrent was amplified by a low-noise current amplifier (Femto, LCA-100K-50M, 50 MV/A transimpedance) and sent through a low-pass filter ($RC = 1$ ms) to further reduce the noise. All photodiode signals, including a power monitor and a MOT fluorescence monitor were acquired using a digital storage oscilloscope.

In Fig. 4.2 we show a typical EW absorption time trace with the probe beam tuned near resonance, for an angle of incidence $\theta_i = \theta_c + 0.52^\circ$. The signal has been averaged 100× to reduce the noise. Without filtering, the absorption signal has a Gaussian shape due to the velocity distribution of the falling atoms. However, in order to suppress slow drifts in the difference photocurrent we used AC coupling (i.e. a high-pass filter) on the oscilloscope. As a result the Gaussian signal has been distorted. Furthermore it is superposed on an exponentially decaying transient originating from switching off the MOT/molasses beams. Although we shielded the photodiodes from the molasses light as much as possible, some light is still detected. Unfortunately the time between switching off the lasers and the arrival of the atoms at the surface is fixed by gravity.

The position of the peak corresponds to the fall time of the atoms. The width is given by the ratio of the size ($\sim 3$ mm) and velocity ($\sim 0.3$ m/s) of the atom cloud as it reaches the surface. The height of the peak is $\sim 2$ mV, which corresponds to an absorbed power of $\sim 80$ pW. The time-integrated signal amounts to $\sim 3 \times 10^6$ absorbed photons, or $\sim 2$ scattered photons per atom in the center of the EW. If we tune the probe laser away from resonance, or increase the angle of incidence, the signal amplitude decreases and the number of scattered photons drops to much less than one per atom. Eventually the signal disappears in the noise, which is dominated by shot noise.

### 4.4 Data processing

Despite the signal distortion, we can extract the amplitude and width of the original Gaussian, by fitting the filtered time trace to an analytical expression. This expression involves error functions due to the known step response function of the high pass filter. We took the fitted height of the Gaussian as the measure for the amount of absorption. The Gaussian width is essentially constant. Thus, the height
4.4 Data processing

Figure 4.2: A typical time-of-flight signal taken with the evanescent-wave probe beam (100× average). The peak around 27 ms is due to the absorption of evanescent probe light by cold atoms arriving at the surface. As a result of high pass filtering the signal has been distorted and sits on top of an exponentially decaying transient (1/e time 26 ms).

Of the peak is proportional to the absorbed EW power, which depends on the EW detuning and the angle of incidence.

For a given angle of incidence we measured time traces for different detunings of the EW probe. The fitted Gaussian height as a function of probe detuning yields an absorption profile as shown in Fig. 4.3. From this we extracted a Lorentzian line width by fitting a Voigt profile

\[
V(\omega) = \frac{A}{\sqrt{2\pi}\Delta} \text{Re} \left\{ \exp \left[ -\left( \frac{\omega - \omega_{eg} + i\Gamma/2}{\sqrt{2\Delta}} \right)^2 \right] \times \right.
\]

\[
\times \text{erfc} \left( -i \frac{\omega - \omega_{eg} + i\Gamma/2}{\sqrt{2\Delta}} \right) \right\},
\]

where \( A \) is an amplitude and \( \text{erfc}(.) \) is the complementary error function. This Voigt profile is the convolution of a Gaussian with a fixed laser line width \( \Delta/2\pi = 1 \text{ MHz} \) and a Lorentzian with variable width \( \Gamma \).

The Lorentzian linewidth \( \Gamma \) contained in the Voigt profile is a fit parameter. We performed measurements of the absorption profile several times for two different angles of incidence \((\theta - \theta_c = 0.16^\circ \text{ and } 0.52^\circ)\). For each absorption profile we find one value for \( \Gamma \). The results are shown in Fig. 4.7. For larger angles, the absorbed power became too small compared to the noise, due to the decreasing EW volume. The vertical error bars in Fig. 4.7 are entirely determined by the scatter of the datapoints as seen in Fig. 4.3.

In the limit of large EW decay length, or \( \theta \to \theta_c \), we expect \( \Gamma \) to tend to the free-space value \( \Gamma_\infty \). Unfortunately, at angles very close to the critical angle, \( \theta - \theta_c < 0.05^\circ \), the probe beam cannot be treated as a plane wave due to the finite diffraction angle. In order to avoid this complicated situation, we performed an independent check by measuring the linewidth in free space. The same probe laser was used to measure the absorption by the atomic cloud while falling, at a height
of 2 mm above the surface. A short flash of probe light was used to illuminate the atoms. The probe beam containing a "shadow" due to absorption by the atoms was recorded on a CCD camera. In addition to this picture (a), we took two more: a background picture with no probe light (b), and a reference picture with probe light but no atoms (c). The optical density was then obtained as \( \ln[(a - b)/(c - b)] \) and was measured for different values of the probe detuning. The detuning was varied in a similar way as in the EW probe measurements. The linewidth was again determined by fitting a Voigt profile. We plotted this free-space value in Fig. 4.7 as the datapoint for \( \theta - \theta_c = 0 \). Our measured free-space linewidth is in good agreement with the known value. This shows that there were no unknown systematic broadening effects, such as saturation, stray magnetic fields, or residual Doppler broadening. We used for this procedure a normal MOT, and a weak probe \( (I \sim 0.1I_{sat}) \) to avoid broadening.

### 4.5 The linewidth of the probe laser

The linewidth \( \Delta \) of the probe laser plays an important role in the measurements. Essentially it pronounced itself as an additional broadening of the absorption linewidth. The grating stabilized diode laser system has a spectral linewidth comparable to the observed atomic linewidth broadenings \([59, 74]\). We determined the laser linewidth in a separate experiment. We used two independent external cavity diode lasers (ECDL) and observed their beat note on a fast photodiode. The ECDL lasers were independently locked using saturation radio-frequency (RF) spectroscopy. We assume that both lasers have equal linewidth. A single laser can also be superimposed with itself via a fiber delay loop, but such self-homodyne systems are ill-suited to measure laser noise frequencies, since impracticaly long delays are required. The non-zero finite linewidth of the lasers will appear as decoherence of the beat signal,
it gives us a numerical value for the laser linewidth.

We measured an averaged beat signal, which is well described by:

$$V(t) = A \sin(\omega t + \varphi) \exp\left(-\frac{t^2}{2\tau^2}\right),$$

(4.2)

where $\omega$ is the frequency difference between lasers, and $\tau$ is a dephasing time, which gives us the linewidth value.

The decaying envelope $\exp[-t^2/2\tau^2]$ appears after averaging on the oscilloscope. The observed envelope has a Gaussian width $(1/\sqrt{\tau}) \tau \simeq 0.76 \, \mu\text{sec}$. We deduced the laser linewidth of about 1 MHz. This value is typical for grating stabilized diode laser systems and it is used in our further calculation. This is the linewidth at relatively short (ms) time scales. For longer timescales we rely on the feedback loop of the laser locking electronics.

4.6 Theory for $\Gamma(z)$

We will now compare the measured $\theta$-dependence of the linewidth to CQED calculations. When the atom approaches the dielectric surface, both the radiative linewidth $\Gamma$ and the resonance frequency $\omega_{eg}$ change in a $z$-dependent way. Although the latter does not change the excited state lifetime, it appears as inhomogeneous broadening in the experiment, because the evanescent wave performs an integration over $z$.

The absorption from the EW probe beam can be calculated by performing a spatial integration of the photon scattering rate over the vacuum half space $z > 0$. The photons that are scattered out of the EW by atoms are missing from the reflected beam.
Figure 4.5: A typical note beat time-trace. First ECDL was locked to saturated absorption resonance of the $F = 2 \rightarrow F' = 3$ transition, second ECDL was locked to saturated absorption resonance of the $F = 2 \rightarrow F' = (2,3)$ cross-over, and was shifted by double-pass an AOM to make the beat frequency in order of 15 MHz. Beams of the two ECDL’s were combined with a non-polarizing beamsplitter and superimposed onto the fast photodiode ($\omega > 100 MHz$). The beams were carefully aligned and have roughly the same power to achieve a clear signal.

probe beam so that the reflectivity drops below unity. This approach works well if the absorption is small ($\ll 1$, i.e. no probe depletion). In the limit of low saturation the photon scattering rate is $\Gamma(z)s(z)/2$, with the saturation parameter given by

$$s(z) = \frac{cU(z)}{I_0} \frac{\Gamma^2_{\infty}/4}{\delta^2(z) + \Gamma^2(z)/4}, \quad (4.3)$$

where $U(z) \propto \exp(-2z/\xi)$ is the EW energy density and $I_0 = 1.6 \text{mW/cm}^2$ is the (free space) saturation intensity. Note that an increase of $\Gamma(z)$ not only increases the Lorentzian width but also multiplies into the photon scattering rate, thus increasing the on-resonance rate. This effect tends to favor the detection of atoms near the surface.

The modification of the radiative linewidth of an atom near a plane dielectric surface has been described theoretically in terms of dipole damping rates $\Gamma_\perp$ and $\Gamma_\parallel$ [23,55]. The subscripts ($\perp, \parallel$) refer to dipoles oriented perpendicular and parallel to the surface, respectively. The dipole damping rates vary with the distance to the
surface $z$, in the notation of Ref. [23] (note also $^*$):

\[
\frac{\Gamma_\perp(z)}{\Gamma_\infty} = 1 + \frac{3}{2} \text{Re} \int_0^\infty \frac{u^3 r^p(u) du}{\sqrt{1-u^2}} \exp(2ikz\sqrt{1-u^2}),
\]

(4.4)

\[
\frac{\Gamma_\parallel(z)}{\Gamma_\infty} = 1 + \frac{3}{4} \text{Re} \int_0^\infty \frac{u du}{\sqrt{1-u^2}} \left(r^s(u) + (u^2 - 1)r^p(u)\right)
\]

\[
\times \exp(2ikz\sqrt{1-u^2}).
\]

(4.5)

Here $k = 2\pi/\lambda$, and $r^p(u)$ and $r^s(u)$ are the Fresnel reflection coefficients for $p$ and $s$ polarization:

\[
r^p(u) = \frac{n^2\sqrt{1-u^2} - \sqrt{n^2-u^2}}{n^2\sqrt{1-u^2} + \sqrt{n^2-u^2}},
\]

(4.6)

\[
r^s(u) = \frac{\sqrt{1-u^2} - \sqrt{n^2-u^2}}{\sqrt{1-u^2} + \sqrt{n^2-u^2}}.
\]

(4.7)

In our experiment we probe $^{87}\text{Rb}$ atoms on the transition $5S_{1/2}(F = 2) \rightarrow 5P_{3/2}(F' = 3)$. An atom in the excited magnetic hyperfine state $|F', m_F\rangle = |3, m\rangle$ can decay to the ground state $|2, m - q\rangle$ with $q = 0, \pm 1$. Choosing the quantization axis perpendicular to the surface, the $q = 0$ decay channel is governed by $\Gamma_\perp$, the $q = \pm 1$ channels by $\Gamma_\parallel$. The decay rate for a given sublevel $|3, m\rangle$ is then given by

\[
\Gamma_m(z) = c_{m,0}\Gamma_\perp(z) + (c_{m,-1} + c_{m,1})\Gamma_\parallel(z),
\]

(4.8)

where $c_{m,q}$ is shorthand for the square of a Clebsch-Gordan coefficient, $c_{m,q} = (2, m - q, 1, q|3, m)^2$. Note that this implies that close to the surface the $m$ states have different lifetimes [23]. The different $\Gamma_m(z)$ curves are shown in Fig. 4.6, together with $\Gamma_\perp(z)$ and $\Gamma_\parallel(z)$. The curve for $|m| = 3$ is not relevant in our experiment because our $p$-polarized probe does not excite these $m$-states. Our $^{87}\text{Rb}$ atoms are in a random mixture of all five $|2, m\rangle$ states. The probe light is linearly polarized, perpendicular to the surface, thus exciting $q = 0$ transitions.

### 4.7 Modeling the measured linewidth

We calculated the amount of photon scattering by two different methods. The first method is based on the steady-state solution of the optical Bloch equations (OBE) in each point along the atom’s trajectory toward the surface. The second method is based on numerical integration of the OBE. As the falling atoms move through the EW, they encounter an increasing Rabi frequency and a position dependent laser

$^*$Note that there appears a printing error in Eq. (58) of [23]; the coefficients $r^p(u)$ and $r^s(u)$ have been interchanged.
Modified radiative properties near a dielectric surface

Figure 4.6: Distance dependence of line widths (upper) and level shifts (lower) of the relevant hyperfine magnetic sublevels. In the upper graph, the dashed curves show the dipole damping rates, see Eqs. (4.4, 4.5). The curves for the $F = |m_F| = 3$ states coincide with that for $\Gamma_\parallel$ and do not contribute in the experiment. The dotted line is the free-space value. In the lower graph the level shifts of the $F = |m_F| = 3$ states do not contribute and are again shown as dashed.

detuning due to the energy level shifts. The internal dynamics of the atom can therefore be rather complicated.

We start with the first method integrating the photon scattering rate $\Gamma_s/2$ over all $z$ and averaging over the $m$-states, we arrive at the absorption of the probe; expressed as a fraction:

$$-\frac{\Delta P}{P} \propto \int_0^\infty \sum_{m=-2}^2 \Gamma_m(z) c_{m,0} \rho(z) e^{-2z/\xi} \frac{\delta_m^2(z) + \Gamma_m^2(z)/4}{dz}, \quad (4.9)$$

where $\rho(z)$ is the atom density. It is $z$-dependent due to the ground-state level shift that accelerates the atoms to the surface. This is well approximated by the Van der Waals potential $-C_3/z^3$, resulting in a depletion of the density near the surface according to $\rho(z) = \rho_0(1 + (z_W/z)^3)^{-1/2}$ with $z_W = (C_3^g/m gh)^{1/3} \approx 50 \text{ nm}$. Here $C_3^g = 5.6 \times 10^{-49} \text{ J m}^3$ is the Van der Waals coefficient and $m gh$ is the kinetic energy with which the atoms fall onto the surface.

The $z$-dependence of the laser detuning $\delta_m(z)$ in Eq. (4.9) accounts for the energy level shifts by unequal amounts for the ground and excited states. The shift of the ground $5S_{1/2}$ state is dominated by the Van der Waals shift $-C_3^g/z^3$. The shift of the excited $5P_{3/2}$ state is more complicated, containing also a resonant component...
4.7 Modeling the measured linewidth

Figure 4.7: Fitted linewidths for varying angle of incidence of the evanescent-wave (EW) probe. The data points at $\theta - \theta_c = 0$ have been measured in free space, instead of with an EW probe. The thin lines are the calculated widths based on Eq. (4.9), the thick lines based on an integration of the optical Bloch equations (see text). The dashed curves show the result when level shifts are not taken into account, the solid curves take into account both broadening and level shifts.

with oscillatory $z$-dependence. We have used expressions for the shifts of both the ground and the excited states from Ref. [54], using transition line strengths taken from Ref. [56]. We have extended the expressions from [54] to account for hyperfine structure. Furthermore we have multiplied the results by a factor $(n^2 - 1)/(n^2 + 1)$, because our surface is a dielectric instead of a mirror. This is known to be correct in the nonretarded limit [57], which gives the dominant contribution in the experiment.

It is evident from Eq. (4.9) that the absorption profile is a convolution of Lorentzians with different widths, amplitudes, and central frequencies. The resulting absorption profiles are strictly speaking no longer Lorentzian. We have numerically calculated the expected absorption profiles using Eq. (4.9). In practice the deviation from a Lorentzian is sufficiently small that we can fit a Lorentzian to the calculated profiles. The fitted widths are plotted in Fig. 4.7, together with the measured widths. In the same Figure we also show the result of the calculation if we do not take the level shifts into account. Clearly, the effect of the level shifts on the observed linewidth is comparable with the direct broadening effect.

Our second calculation is based on a numerical integration of the optical Bloch equations OBE. This method takes into account the transient internal dynamics and its possible effects on the observed linewidth. Note that such transients would at least qualitatively have the observed signature of a broadening that increases with a decreasing EW decay length (increasing angle of incidence).

To investigate this we numerically integrated the time-dependent optical Bloch equations for an atom moving through the EW field. We made the approximation that the atom is a two-level system. First, the known ground state level shift was used to solve for the accelerated motion $z(t)$ towards the surface. This solution was then used to define a time-dependent Rabi-frequency $\Omega_1(z(t)) = \Omega_{10} \exp(-2z(t)/\xi)$, and similarly for the detuning $\delta(z(t))$, and radiative linewidth $\Gamma(z(t))$. Using these
time dependent parameters we numerically integrated the optical Bloch equations to obtain the time evolution of the Bloch vector \((u(t), v(t), w(t))\). Note that power broadening is naturally included in this method.

The first step is to calculate the perpendicular motion of the atom \(z(t)\) in the potential near the surface. We write the ground state potential as \(V_g(z)\), and the velocity at which the atom enters the EW as \(v_i < 0\). The velocity at a given height \(v(z)\) is fixed by energy conservation, \(mv^2(z)/2 + V_g(z) = mv_i^2/2\). The motion of the atom is then found by solving the differential equation

\[
\dot{z} = v(z) = \sqrt{v_i^2 - 2V_g(z)/m}.
\]  

(4.10)

In Fig. 4.8 we show the numerical solution of this equation of motion, for \(v_i = -0.26\text{ m/s}\) and \(z(0) = 0\).

The internal dynamics in the EW are governed by the optical Bloch equations (OBE). The falling atoms are in a mixture of all five magnetic sublevels \(|F = 2, m_F\rangle\). The EW probe is \(p\)-polarized, i.e. the polarization is approximately linear, perpendicular to the surface. We make the approximation of treating every pair of levels \(|F = 2, m_F\rangle, |F' = 3, m_F\rangle\) as a two-level system, i.e. we neglect effects of optical pumping during the transit through the EW.

The input parameters for the OBE are the Rabi frequency \(\Omega\), the laser detuning \(\delta_L\), and the radiative linewidth \(\Gamma\). All three change in a known \(z\)-dependent way. Since we have already solved the motion \(z(t)\) of the atom, the \(z\)-dependent input parameters of the OBE can now be written in time-dependent form, \(\Omega(t) = \Omega(z(t))\), and similar for \(\Gamma\) and \(\delta_L\). The OBE for the pair of levels \(|F = 2, m_F\rangle, |F' = 3, m_F\rangle\) now take the following form \cite{75}:

\[
\dot{u} = \delta_m(z(t))v - \frac{\Gamma_m(z(t))}{2}u
\]  

(4.11)

\[
\dot{v} = -\delta_m(z(t))u - \Omega_m(z(t))w - \frac{\Gamma_m(z(t))}{2}v
\]  

(4.12)

\[
\dot{w} = \Omega_m(z(t))v - \Gamma_m(z(t))w - \frac{\Gamma_m(z(t))}{2}
\]  

(4.13)
4.8 Discussion

We numerically solve the OBE choosing as the initial condition that the atom is in the ground state at a large negative time: \( u(t_0) = v(t_0) = 0, \ w(t_0) = -1/2 \), where \( \Gamma_\infty t_0 = -1000 \). We then numerically integrated the OBE until \( t = 0 \). A typical example of the solution is shown in Fig. 4.9. From this solution we obtain the number of scattered photons by the atom on its way down to the surface by integrating \( \int \Omega(t)v(t)dt \) [75]. We average this quantity over the different \( m \) states and vary the laser detuning \( \delta_L \) over the absorption resonance. This yields an absorption profile, to which we fit a Lorentzian profile. Finally, the resulting Lorentzian linewidth is determined for different values of the EW decay lengths (angles of incidence). The final result is plotted together with our experimental data points in Fig. 4.7.

4.8 Discussion

Both the optical Bloch analysis and the steady-state analysis predict only about half the observed broadening of up to about 25%. The difference cannot be explained by obvious sources of spurious broadening. These include Doppler broadening (< 2%), Zeeman broadening due to a spurious magnetic field (< 3%), and power broadening (< 0.5%). Furthermore these broadening mechanisms would not show the observed signature of increasing with the angle of incidence. The calculation based on integration of the OBE shows that transit time broadening cannot explain the difference either. Later experiments have shown an importance of the prism surface roughness [34]. However the influence of the surface roughness on the absorption is unclear.

As a tentative explanation we invoke the presence of local Stark shifts caused by charged or polarized particles on the surface. Based on a straightforward model calculation we find that a surface charge density of \( 45e/\lambda^2 \) yields a 10% linewidth increase. Remarkably, such a charge density corresponds to an average distance between the charges of order \( \sim 100 \text{ nm} \), which is just the distance scale to which our experiment is very sensitive. These calculations only weakly reproduce the angular dependence shown by the data. Recently McGuirk et al. have reported that Rb adsorbed on a Si or Ti surface generates local Stark shifts that were measurable as a change in the trapping frequency of their magnetic trap [76]. The authors mention that similar effects on a glass surface like ours are very small. However, their experiment measured only changes upon depositing clouds of Rb atoms on the surface, whereas our experiment is also sensitive to statically present adsorbates. Furthermore, there may be other charged or polarized adsorbates on the surface. For these reasons local Stark shifts do seem to present a likely mechanism to explain our results. Our experiment is also complementary to Ref. [76] in the sense that the latter measures a global effect, whereas our experiment is sensitive only to local variations of the electric fields. Unfortunately we have no detailed information about possible adsorbates to make a more quantitative analysis.

The present method does not let us to distinguish between CQED linewidth broadening and inhomogeneous broadening. However a shift of the resonant frequency is a feature only of level shifts. So the averaged effect of level shifts can be
Figure 4.9: Top: time-dependent detuning, for $\delta_L/\Gamma_\infty = -0.5$, $\xi = 500\,\text{nm}$. The curves for the different magnetic sublevels are indistinguishable on this scale. Middle: time-dependent Rabi frequency for the magnetic sublevels $|m_F| = 0, 1, 2$ (top down). Bottom: Corresponding numerical solution for the $v$ component of the optical Bloch equations. The three curves are for the magnetic sublevels $|m_F| = 0, 1, 2$ (top down).

A straightforward way to improve the experimental data would be to use a colder and denser atomic cloud. Special attention need to be paid to a roughness and to local charge of the surface.
In conclusion, we have used evanescent-wave (EW) spectroscopy to observe a broadening of the absorption linewidth of the $D_2$ resonance line of $^{87}\text{Rb}$, caused by the surface. Part of the broadening can be explained as a combined effect of CQED (homogeneous) linewidth broadening and inhomogeneous broadening caused by level shifts due to the proximity of a dielectric surface. On the basis of models discussed above, we expect that both mechanisms make comparable contributions to the broadening. The present method does not allow distinction between the two.

The total observed broadening of up to 25% was about two times larger than expected from CQED calculations. Although the source of this additional broadening remains essentially unknown, we can eliminate several possible spurious sources of broadening. A telltale feature of the observed broadening is an increase with the angle of incidence, i.e. an increase for shorter EW decay length. This is expected for CQED type effects but not for most spurious effects. A plausible candidate to explain the additional broadening is the presence of local Stark shifts due to charged or polarized adsorbates on the surface. Support for this tentative explanation can be found in the literature [76].
We trap cold $^{87}$Rb atoms in a far off-resonance optical dipole trap (ODT). We employ a 2 $\mu$m wavelength, 50 W Thulium doped fiber laser and a 1 $\mu$m wavelength, 30 W diode-pumped Yb: YAG thin-disk laser. ODTs were created using either one or two focused 1 $\mu$m and for 2 $\mu$m beams. This is the first time that an ODT is created with 2 $\mu$m IR light. We vary the waist of the trapping beams and trap geometry in order to reach good starting conditions for evaporative cooling. Atoms are directly loaded from the MOT. We describe in detail the loading procedure of the cold atoms from the MOT to the ODT, which is based on a dark molasses. We observe an increase of the phase-space density in ODTs of more than three orders of magnitude compared to the phase-space density in a MOT.
5.1 Introduction

The use of focused laser beams as traps for neutral atoms was first proposed in 1978 [40]. In 1986, S. Chu and coworkers succeeded in holding about 500 sodium atoms for several seconds in the tight focus of a red-detuned laser beam [41]. Optical dipole traps (ODT) became a standard tool with the development of sub-Doppler cooling and the invention of the magneto-optical trap (MOT) as a source for cold atomic samples. The far-off resonant ODT gained attention with its demonstration by Miller et al. in 1993 [8], where the authors used a conventional MOT as a source for the cold atoms. In this work $1.3 \times 10^3$ cold atoms of $^{85}$Rb were trapped in far-off-resonant single focused beam at an atomic temperature of 0.4 mK. Spontaneous emission of photons is negligible in such a trap, so a far off resonance ODT provides a virtually conservative potential. In 1996 about $9 \times 10^5$ cold Cs atoms were trapped in quasi-electrostatic trap (QUEST) 20 W CO$_2$ single beam trap [45]. A QUEST is a variant of an ODT, in which the laser frequency $\omega$ is less than half the frequency of the lowest electric dipole transition. In 2001 Barrett et al. demonstrated the all-optical formation of an atomic Bose-Einstein condensate [10] by employing a QUEST, which was created by 2 crossed, focused CO$_2$ laser beams. ODTs open the possibility of forming condensates of nonmagnetic atoms, of molecules, and of several spin states of the same atom, because they do not rely on a magnetic moment of the trapped species. Recently ODTs were used to create a spinor condensate [10, 11] and to trap and evaporate nonmagnetic species [20, 21].

Here we describe our experiments where we use focused 2 $\mu$m and 1 $\mu$m wavelength laser beams to create ODTs for cold $^{87}$Rb atoms. A MOT is used as a source of cold atoms. We employ a dark molasses stage before loading atoms into an ODT. We experimentally investigate the properties of an ODT. We study how the loading of cold atoms into an ODT depends on various parameters, and in particular discuss the role of the volume of an ODT.

The trap depth is limited from above by the available optical power, and from below by the temperature of atoms in the dark molasses. The waist of the trapping beam is a crucial parameter of an ODT. The waist of the beam is essentially the only free parameter that can be varied over a large interval. Cennini et al. [14] provide a simple scaling law for the number of atoms and their collision rate in a single beam ODT. However, we think that these scaling laws are applicable to a fixed trap depth and not to the fixed optical power of the trapping beam, since the ratio of the temperature of atoms and the trap depth should be suitable for further evaporative cooling. Essentially one can reach an impressive density of atoms by increasing the optical power of the trapping beam (before or after loading) of the ODT, however for further evaporation one needs to decrease the optical power, and consequently decrease the density of the atoms back to the initial value due to decompression. Therefore it is useless to produce an ODT as deep as possible, because one can not benefit from the high density and collision rate of the atoms due to the absence of a ”RF-knife” for ODTs.

Several types of ODT’s are discussed in this chapter, namely: 1) a 2 $\mu$m single focused beam, 2) a 1 $\mu$m single focused beam, 3) a crossed-beam trap, where we
cross a 1 µm and a 2 µm beam. We study the difference of the loading efficiency for different colour (1 µm and 2 µm) ODTs. We also discuss the differences between single focused beam and crossed-beam trap ODTs. The temperature, the number of atoms, and the density of the atoms in ODTs have been measured. We have observed a substantial increase of the atomic density, as well as the phase-space density and the collision rate in ODTs, compared to the situation in the MOT.

5.2 2 µm: a new wavelength in the ODT family

Here we introduce the use of a 2 µm laser for an ODT. Like a CO₂ laser (10 µm wavelength) a 2 µm Thulium fiber laser provides a very large detuning. However, unlike light from a CO₂ laser, 2 µm light is still transmitted by most optical materials such as quartz and BK7. It thus has the advantage that no special vacuum windows are needed. The large detuning and a high power of a Thulium fiber laser give us the possibility to create a deep (up to 1 mK) and conservative potential. Light of 2 µm creates a confining potential for both the ground and the excited state of the $5s^2S_{1/2} \rightarrow 5p^2P_{3/2}$ transition, just as 10 µm light does, which has been argued to be a substantial advantage of a CO₂ trap compared to a 1 µm trap [14, 34], and which was one of the motivations to study 2 µm traps for $^{87}$Rb atoms.

We demonstrate the essential differences of 1.03 µm, 2.0 µm and 10.6 µm wavelength on an example of a single beam ODT with a trap depth of 500 µK. All three trap beams are assumed to be focused to a waist of 50 µm, although it is more difficult to focus a 10.6 µm beam to the same waist as a 1.03 µm beam due to a larger diffraction limit.

First of all, $^{87}$Rb atoms have different polarizabilities for the ground ($\alpha_g$) and excited ($\alpha_e$) level [10, 34]. There is a substantial difference in the value and even the sign of the light shift of the excited level. It is positive for 1 µm light and negative for 2 µm and 10.6 µm light [34]. We summarize the values of the polarizabilities in Table 1. For the polarizability of the excite state we took the $F = 3, m_F = 3$ state and linearly polarized light.

<table>
<thead>
<tr>
<th>λ, (µm)</th>
<th>$\alpha_g$, ($\times 10^{-39}$ m²C/V)</th>
<th>$\alpha_e$, ($\times 10^{-39}$ m²C/V)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.03</td>
<td>13.6</td>
<td>-13.6</td>
</tr>
<tr>
<td>2.0</td>
<td>6.2</td>
<td>26.3</td>
</tr>
<tr>
<td>10.6</td>
<td>5.3</td>
<td>9.0</td>
</tr>
</tbody>
</table>

Table 5.1: Polarizabilities for the ground ($\alpha_g$) and the excited ($\alpha_e$) level of $^{87}$Rb for different wavelengths.

The different polarizabilities lead to different powers (10.6 W, 23.3 W, and 27.2 W) that are necessary for a trap depth of 500 µK. The trap frequency perpendicular to the direction of the trap beam propagation (radial frequency $\omega_r$) is the same in all three cases and is equal to $\omega_r = 2\pi \times 1.4$ kHz. The trap frequency along the direction of the trap beam propagation is different due to a different Rayleigh length.
The frequencies along the trap beam direction are $\omega_z = 2\pi \times (6.4; 12.5; 62.4)$ Hz for a wavelength of 1.03, 2 and 10.6 $\mu$m respectively. This is a substantial difference. It seems quite likely that the higher density that was achieved in CO$_2$ ODTs compared to near infrared ODTs results in a stronger axial force, pulling the atoms into the trap minimum.

**Figure 5.1:** Level shift of the ground (solid) and excited (dashed) levels in ODTs formed by a focused 1.03 $\mu$m, 2 $\mu$m, 10.6 $\mu$m beam. We choose equal trap depths and equal beam waists (50 $\mu$m). The level distance far from the center of the trap corresponds to the MOT detuning of -12 MHz.

**Figure 5.2:** Same situation as Fig. 5.1. but using the dark MOT detuning of -90 MHz.

The difference in level shifts of the ground and excited levels are demonstrated in Fig. 5.1 and Fig. 5.2. In Fig. 5.1 one can see that 2 $\mu$m light shifts the atoms resulting in blue detuning of to the MOT light. For dark molasses the situation is less extreme due to the large red detuning of the MOT beams, but still atoms in the 2 $\mu$m ODT feel a different detuning than atoms outside of the ODT. It remains unclear which shift of the excited level is optimal for an efficient loading. If an atom in the ODT is shifted closer to resonance (i.e. 2 $\mu$m and 10.6 $\mu$m), it scatters more MOT photons, which leads to more efficient cooling but restricts the atomic density due to rescattering of the MOT light. If the atom in the ODT is shifted away from the resonance (i.e. 1.03 $\mu$m), they scatter less photons, which leads to less efficient
cooling but a higher atomic density. Extreme cases must be avoided where atoms are tuned into resonance (or even blue detuned) with respect to the MOT light, or detuned very far from the resonance with respect to the MOT light. This prohibits the use of very deep traps, where the level shift is equal to or larger than the dark molasses detuning.

5.3 Depth and frequencies of a far off resonance optical trap

Trap depth and trap frequencies are the key parameters of the dipole trap. For a Gaussian beam propagating in the $z$ direction the trap depth can be calculated as:

$$U(r, z) = \frac{\alpha P}{\pi \varepsilon_0 w(z)^2} \times \exp \left[-\frac{2r^2}{w(z)^2}\right],$$

(5.1)

where $w(z) = w_0 \sqrt{1 + (z/z_R)^2}$, $z_R$ is the Rayleigh length, $r$ is the coordinate perpendicular to the beam axis, $\alpha$ is the polarizability for a given wavelength, $P$ is the optical power and $w_0$ is the waist of the beam (defined as a $1/e^2$ radius). Finally, $c$ and $\varepsilon_0$ are the speed of light and vacuum permittivity.

Practically useful values for the trap depth are set by the temperature of the atoms. During loading the trap should be substantially deeper than the atomic temperature. On the other hand, there is also no reason to make the trap many orders of magnitude deeper than the atomic temperature; since it does not give any advantages and it can tune atoms into resonance with the MOT light. Furthermore the ratio of the trap depth and the atomic temperature (known as the $\eta$ parameter) should be suitable for further evaporation (typically in the range between 6 and 10).

In a harmonic approximation near the bottom of the trap, the trap frequencies for a single beam ODT can be calculated as:

$$\omega_r = \sqrt{\frac{4\alpha P}{\pi \varepsilon_0 m w_0^4}},$$

(5.2)

$$\omega_z = \sqrt{\frac{2\alpha P}{\pi \varepsilon_0 m z_R^2 w_0^2}},$$

(5.3)

where $\omega_r$ and $\omega_z$ are frequencies perpendicular to and along the trap beam propagation, respectively.

Throughout this chapter we calculate the phase-space density $\rho$, collision rate $\gamma$ and peak density $n_0$ as:

$$\rho = \frac{1}{3} \frac{N}{1.2} \left(\frac{\hbar \bar{\omega}}{k_B T}\right)^3,$$

(5.4)

$$\gamma = \frac{m \bar{\omega}^3 \sigma N}{2\pi^2 k_B T}$$

(5.5)
\[ n_0 = \frac{N\bar{\omega}^3}{(2\pi k_B T/m)^{3/2}} \]  

where \( N \) is the number of atoms, \( T \) is the temperature, \( \sigma \) is the scattering cross-section and \( \bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3} \) is the geometric mean of the trap frequencies. In Eq.(4) the factor of 1/3 accounts for the fact that the atoms are equally distributed over the spin states \( (F = 1) \). With this definition a phase-space density equal to 1 corresponds to the phase transition to quantum degeneracy.

Since both the collision rate and the phase-space density are proportional to the cube of the mean trap frequency (assuming a fixed temperature and atom number), it is very useful to create an ODT with high initial frequencies. For a fixed trap depth, the beam waist \( w_0 \) determines the trap frequencies.

For realistic values of the available optical power and for useful trap frequencies the waist has to be substantially smaller than the MOT diameter. The choice of the waist is usually a trade off between a larger trap volume (and thus better overlap with the MOT), and higher frequencies of the ODT. It is always favorable to make a tight focus to have good initial conditions i.e. high collision rate for forced evaporation. Unfortunately, this can lead to unpractically few atoms in the dipole trap.

### 5.3.1 Heating in an ODT

The heating rate of atoms is an important parameter of an ODT. The two main sources of heating are: 1) scattering of light of the ODT, 2) parametric heating caused by instabilities in the power and the alignment of the trap beam. The scattering heating is unavoidable, and is defined by the wavelength and intensity of the ODT. For a chosen trap depth of 500 \( \mu \)K, we calculated a scattering rate of 3.4, 1.7, and 1.5 photons per second, for \( \lambda = 1.03, 2.0, 10.6 \mu m \), respectively. These values for the scattering rate correspond to heating rates of 150, 76, 63 nK/s, respectively. These heating rates are still negligible in our experiment. However we notice that generally 2 \( \mu m \) and 10.6 \( \mu m \) ODTs have lower heating due to the scattering of ODT photons. If long evaporation ramps (\( \sim 1 \) min) are used, heating in an ODT become noticeable. Heating due to laser intensity noise is generally unrelated to the colour of the trap light. In practice the choice of lasers for a given wavelength is restricted. CO\(_2\) lasers have been reported to be stable [22]. For 2 \( \mu m \) we use a fiber laser from IPG. This laser is relatively quiet, and we did not measure any noticeable heating in the 2 \( \mu m \) ODT due to laser noise, although we measure strong heating in the 2 \( \mu m \) crossed beam ODT, which we attribute to an unstable interference pattern in the cross region. For the 1 \( \mu m \) wavelength, we use a Yb:YAG Thin-Disk Laser from ELS, which we find to be quite noisy. We measured a heating rate of about 20 \( \mu K/s \). A substantial improvement of the noise level was achieved by removing the etalon from the laser cavity and by shortening the laser cavity. This lowered heating rate is less than 0.1 \( \mu K/s \). The drawback of this solution is multi-mode behavior of the laser without an etalon, which prohibits e.g. the creation an optical lattice with this laser.
5.4 Loading procedure

We produced clouds of cold $^{87}$Rb atoms using a conventional MOT inside a ultra high vacuum cell (base pressure $p \approx 5 \times 10^{-11}$ mbar). The MOT loading time is about 10 s. We load about $100 \times 10^6$ atoms into the MOT, with a peak density $\sim 1.5 \times 10^{10}$ cm$^{-3}$. Then we switch off the Rb dispenser and switch on the high power laser beams. We wait about 3 s to let the dispenser cool down, which improves the life-time of the trapped atomic cloud. We use dark molasses to load atoms from the MOT into the ODT. The purpose of the dark molasses is to suppress photon rescattering of MOT light, which is known to be the density limiting mechanism in an ODT [10]. This dark molasses resembles the dark MOT procedure, which has been used in [10, 14]. The absence of the magnetic field is the only difference. During the dark molasses stage the atoms spend most of their time in the $F = 1$ dark state. Therefore the atoms scatter less light, and reach higher densities. An efficient loading procedure is a compromise between suppressing the photon rescattering and keeping the atoms in the cloud.

We optimized our loading procedure for the following parameters: 1). the molasses beam detuning, 2). the size of the dark spot in the repumper beam, 3). the duration of the dark molasses phase.

This has led to the following loading procedure. We simultaneously:

1. turn off the normal repumper beam and switch on a repumper beam with a spatial dark region (a square with about 1 mm$^2$ area) in the center of the beam. The optical power of the ”dark” repumper is about 200 µW.

2. increase the MOT beam detuning to -15 $\Gamma$ with respect to the $F = 2 \rightarrow F' = 3$ transition of the $D_2$ line (the MOT transition).

3. switch off the magnetic field of the MOT.

This dark molasses lasts for about 30 ms. Then we turn off the dark repumper beam. Finally, after an additional 3 ms we switch off the molasses beams, so that most of the atoms end up in the $F = 1$ state. We achieve a higher loading efficiency using a repumper with a dark spot than by simply decreasing the repumper intensity, which is described in [65].

We wait about 50 ms until the MOT atoms fall down, and only the atoms into the ODT are left. Then we switch off the dipole beams to release the atoms trapped in the ODT. We shine a resonant probe beam, mixed with repumper light (because most of the atoms are in the $F = 1$ state). Because of the large light shift of the levels in the ODT it is important to image the atoms, only after releasing them from the ODT. Otherwise the atoms are out of resonance with the probe beam. This is illustrated in Fig. 5.3.

Using a time of flight technique, we can measure the number of trapped atoms and their temperature. The dark molasses leads to sub-Doppler cooling of the MOT atoms. We measure a temperature of 25 µK after the dark molasses stage. The temperature of the atoms in the ODT is similar to the dark molasses temperature,
Figure 5.3: Absorption image (1×4 mm²) of an atomic cloud in an ODT, formed by a single 2 µm beam focused to 36 µm. One can see a light line in the middle of the atomic cloud. Due to the large light shift atoms in the center of the trap are transparent for the probe beam. These atoms are detuned for the probe light and thus, and they focus the probe beam. As a consequence we see that the line in the center of the trap is lighter than a background. At the side of the trap the light shift is smaller so that atoms can be observed (the two dark lines above and below).

unless the ODT is either too deep or too shallow compared to the temperature of the dark molasses. We will discuss this point later in this chapter. This temperature is substantially lower than a regular, bright MOT temperature (about 200 µK), but still larger than the low density conventional molasses temperature (about 5 µK for our experimental setup). We measure an increase of the peak density to 3 × 10¹⁰ cm⁻³ during the dark molasses, which is approximately a factor of 2 higher than in a normal MOT. In addition we measure that about 90 percent of the dark molasses atoms are in F = 1, which is believed to be optimal for loading [65]. The dark molasses loading procedure increases the number of atoms in the ODT by a factor of 5, compared to direct loading from optical molasses. Therefore, a proper loading procedure is important to achieve a dense atomic sample into an ODT.

Unfortunately, the loading efficiency of MOT atoms into an ODT is difficult to model. The so-called two box model proposed in Ref. [78], and later developed in Ref. [79], is not applicable to our situation for several reasons. First, atoms in dark molasses continuously interact with the MOT light, which leads to cooling and limits the density. Second, the typical life-time in the dark molasses is about 50 ms [65], whereas the collision rate in the dark molasses is about 1 s⁻¹. Therefore the atoms do not have enough time to reach thermal equilibrium in the system of dark molasses+ODT during 50 ms. Scattering of MOT light probably plays a far more important role than the atomic collisions in the dark molasses, as evidenced by the substantial increase of the number of atoms in the ODT during a dark molasses or a dark MOT time. The model developed in [79] does not apply to loading directly from dark molasses. For example, Cennini et al. [14] reported the loading of 4 × 10⁶ atoms from a MOT of 60 × 10⁶ atoms. Equations (1) of Ref. [79] would predict that almost all MOT atoms (more than 95%) end up in the ODT. The disagreement is even more striking if we compare the prediction with the result of loading of near infrared ODTs. For example, Dumke et al. [48] report about loading 2 × 10⁹ sodium atoms in the ODT from a MOT of 7 × 10⁹ atoms, where the model predicts a loading efficiency of about 50%. We conclude that this thermodynamic approach does not yield a good description of loading of cold atoms into an ODT directly from the
dark molasses, dark MOT or conventional MOT. The description should include the interaction of atoms with the MOT light.

Various groups have reported different loading efficiencies. It seems that spatial overlap of the ODT and the MOT is a key factor in the loading efficiency.

We measured the dependence of the number of atoms loaded into the ODT and the loading efficiency on the number of atoms in the MOT. We define the loading efficiency as the ratio of the atom number loaded into the ODT and the atom number in the MOT. Since the density of the MOT does not change as function of the atom number, larger atomic number in the MOT leads to a larger size of the MOT.

We see in Fig.5.4 that the number of atoms loaded into the ODT grows with the MOT atom number, but slower than linear. We observe a clear decrease of the loading efficiency, when the number of atoms in the MOT increases. This supports the idea that the spatial overlap of the ODT with the MOT plays a crucial role.

Another important process is the redistribution of cold atoms inside the crossed-
beam ODT after switching off the MOT beams. Indeed the cross region of the ODT is the potential minimum for atoms in the wings. Since atoms are not into the equilibrium immediately after the loading, the central dimple collects atoms from the wings. This process occurs on the time-scale of the collision rate of the atoms in the wings. Indeed we observe a loading of the central peak during the first 300 ms after switching off the MOT beam, see Fig. 5.5. After that, the atom number in the central peak slowly decreases due to vacuum background collisions.

5.5 A 2 μm single beam optical dipole trap

We performed trapping of MOT atoms in an ODT formed by a single focused 2 μm beam. The single beam ODT is the simplest type of ODT. It is easy to align such a trap and important features of loading are more clear than in case of a crossed beam ODT. The ODT was aligned to the center of the MOT. We used the loading procedure described earlier in this chapter. A distinctive feature of 2 μm light is the large negative light shift of the excited 5P3/2 state, which is a factor of 5 larger than for the ground 5S1/2 state of the same transition. A large detuning of the cooling laser is therefore needed during the dark molasses phase, in order to prevent the cooling laser to be effectively blue detuned in the center of the ODT.

We show the dependence of the number of atoms in the 2 μm ODT vs. the optical power of the 2 μm trap beam in Fig. 5.6 using a single beam focused to a waist of 36 μm. We can see that the number of atoms loaded into the ODT initially increases linearly with power, reaches a maximum for an optical power of about 10 W, and finally decreases. This counterintuitive result, that a deeper trap leads to smaller number of trapped atoms, can be understood as a consequence of the large light shift of the excited state. The effective detuning of the cooling beams decreases with an increase of the trap beam intensity, which makes the loading less efficient.
We have made several single beam 2 $\mu$m traps. We observed a clear advantage of tightly focused trapping beams. We use a 2 $\mu$m trap beam with a waist of about 22 $\mu$m for a detailed study of a 2 $\mu$m single beam ODT including evaporation. The waist of 22 $\mu$m is smallest waist, which we succeed to create. A smaller waist is hard to make, because it requires to position the focusing lens closer to the MOT, which is not possible due to the MOT beams, MOT coils etc. In this particular case, and generally in this chapter, we measure the trap frequency by observing atomic losses caused by modulation of the trap beam power (this measurement is described in a previous chapter). From this measured frequency we deduce the beam waist and calculate the trap frequency for other optical powers of the trap beam. For a single beam ODT we also measured axial frequency by observing the axial breathing (also described in a previous chapter).

A waist of 22 $\mu$m provides high trap frequencies. For a beam power of 3 W the frequencies are $\omega_r = 2\pi \times 2.6$ kHz and $\omega_z = 2\pi \times 53$ Hz, with a trap depth of $U=330$ $\mu$K. We trapped about $160 \times 10^3$ atoms at a temperature of 15 $\mu$K. The temperature of the trapped atoms was measured using the standard time-of-flight procedure [35] (see Fig. 5.7). This procedure includes: 1. taking an absorption image of the cloud as in Fig. 5.8. after a variable expansion times, 2. for each image we obtain a cloud diameter by fitting in the vertical direction, 3. we fit the cloud diameter vs. expansion time to the function $\sqrt{\sigma_0^2 + (\sigma_v t)^2}$, where $\sigma_0$ is an initial size, $\sigma_v$ is a velocity of thermal expansion. We deduce the temperature from $\sigma_v = \sqrt{k_B T/m}$.

It should be noted that for a 2 $\mu$m single beam ODT the atomic temperature is usually a little bit lower than the dark molasses temperature and also lower than the temperature of a 1 $\mu$m single beam ODT, apparently due to the large light shift of the excited state (atoms experience a stronger dissipative force in the ODT).

For a typical set of measured temperature, atomic number and trap frequencies we calculate a peak atomic density $n_0 = 1.6 \times 10^{13}$ cm$^{-3}$, a collision rate $\gamma = 410$ s$^{-1}$, and phase-space density $\rho = 5.1 \times 10^{-4}$, all before evaporation. The achieved numbers are substantially lower that numbers achieved in CO$_2$ ODTs [10, 14].

Figure 5.6: The number of atoms in the ODT vs. the optical power of the trap beam. The ODT was formed by a single 2 $\mu$m beam focused to 36 $\mu$m.
Figure 5.7: A measured time of flight sequence: the atomic cloud was released from the ODT and expanded for a variable time (0, 1, 2, 3 ms). Directly after the expansion an absorption image of the cloud is taken. The cloud grows because of thermal expansion. Each picture shows an area of 2×3 mm$^2$ (in the vertical×horizontal direction respectively).

![Figure 5.7: A measured time of flight sequence: the atomic cloud was released from the ODT and expanded for a variable time (0, 1, 2, 3 ms). Directly after the expansion an absorption image of the cloud is taken. The cloud grows because of thermal expansion. Each picture shows an area of 2×3 mm$^2$ (in the vertical×horizontal direction respectively).](image1)

Figure 5.8: The atomic cloud was released from the ODT and expanded for a variable time. The vertical size of the atomic cloud was measured by taking an absorption image of the cloud as in Fig. 5.7. For each image we obtain a cloud diameter by fitting in the vertical direction.

![Figure 5.8: The atomic cloud was released from the ODT and expanded for a variable time. The vertical size of the atomic cloud was measured by taking an absorption image of the cloud as in Fig. 5.7. For each image we obtain a cloud diameter by fitting in the vertical direction.](image2)

Despite the fact that the collision rate seems to be large, straightforward evaporation from this trap leads to a decompression and thus, to a fast decrease of the collision rate.

5.6 A 1 µm single beam optical dipole trap

We also performed trapping of MOT atoms in an ODT formed by a single focused 1 µm beam. This allows us to compare the loading of different colour ODTs from the same MOT. Beside this it is easier in practice to tightly focus a 1 µm ODT due to a smaller diffraction limit. Again the ODT was aligned with the center of the MOT. We used the same loading procedure as for the 2 µm single beam ODT. The 1 µm single beam ODT is very similar; it also creates a deep and almost
conservative potential. The difference in wavelength pronounces itself in two ways: 1) the Rayleigh length is twice as long for a 1 µm beam compared to a 2 µm beam, if they are focused to the same waist. 2) the diffraction limit of the waist of a focused 1 µm beam is a factor of two smaller than for 2 µm. This results in higher trap frequencies for a 1 µm ODT, than for a 2 µm trap. 3) the 1 µm light shift of the excited 5P3/2 state is opposite to the shift of the ground 5S1/2 state. The laser cooling beams are effectively more red detuned in the center of the 1 µm ODT with respect to the atoms in the dark molasses.

We used a 1 µm trap beam with a waist of about 13.5 µm for a study of the 1 µm single beam ODT and for evaporation in this trap. This waist is practically achievable and provides high trap frequencies. For a beam power of 0.85 W the calculated frequencies are \( \omega_r = 2\pi \times 5.8 \text{ kHz} \) and \( \omega_z = 2\pi \times 100 \text{ Hz} \), with a trap depth of \( U=550 \mu \text{K} \). Again the temperature of the trapped atoms was measured using a time-of-flight procedure. The number of trapped atoms monotonically increases with the trap depth, in contrast to a 2 µm ODT. This allows us to use a deeper trap to collect more atoms.

For a typical set of measured temperature, atom number and trap frequencies we calculate an atomic peak density \( n_0 = 1.2 \times 10^{14} \text{ cm}^{-3} \), a collision rate \( \gamma = 3.9 \times 10^3 \text{ s}^{-1} \), and a phase-space density \( \rho = 1.7 \times 10^{-3} \). Achieved density and collision rate are comparable with numbers achieved in CO₂ ODTs [10, 14]. However these density and collision rate are obtained by producing a very deep trap. The number of trapped atoms is an order of magnitude lower than in Ref. [10, 14]. This is the result of the smaller waist, and thus smaller trapping volume. The collision rate of 3.5 kHz is very large compared to most cold atom experiments. Nevertheless, we do not gain much in terms of efficient evaporation because of the large trap depth, and therefore the high \( \eta \) parameter. A decrease of the trap depth leads to decompression and thus a reduction of density and of collision rate.

### 5.7 Crossed beam optical dipole traps

Two crossed focused laser beams can create a deep attractive potential and a strong confinement in all dimensions for cold \(^{87}\text{Rb}\) atoms. The possibility to reach high trap frequencies in all dimensions is the main motivation for building a crossed beam optical dipole trap. We use a focused 1 µm and a focused 2 µm beam to produce such a trap. The use of beams with two different wavelengths has several advantages. First, one can modify and optimize the excited level shift in the center of the trap. Second, one automatically avoids problems with interference fringes of the two beams. Nevertheless we have to point out that the use of two colours for an ODT is not necessary. Since the waist of a focused beam is much smaller than the size of the MOT in our experiment, alignment of a crossed beam ODT is more delicate and complicated than alignment of a single beam trap. In practice one has only limited control over the quality of the overlap of the laser beams. We produced several crossed beam optical dipole traps with different values for the beam waists.
Here we discuss two crossed beam ODTs. First, a 'large' volume ODT (the waist of the 1 µm beam is $w_0=75$ µm, and the waist of the 2 µm beam is $w_0=60$ µm), and a 'small' volume ODT (the waist of the 1 µm beam is $w_0=15$ µm, and the waist of the 2 µm beam is $w_0=37$ µm). We compare the atom number, the density and the collision rate of these two crossed beam ODTs, then we discuss their potential as a starting point for forced evaporation. The beams in both crossed beam ODTs were crossed at near right angle, which provides a tight confinement in all directions. We used the same loading procedure as for the single beam traps. In both ODTs the beams were crossed in the horizontal plane.

We start with the large volume ODT. For beam powers of 17.5 W and 4.9 W this ODT has trap frequencies of about 820 Hz and 420 Hz respectively (associated with 2 µm and 1 µm single beam transverse frequencies respectively) in the intersection region. The total trap depth is about 360 µK.

![Figure 5.9: An absorption image of the cloud immediately after release from the large volume crossed beam ODT. We observe wings, and a dense central part at the position where the trap beams cross. The probe beam lies in the plane of the crossed beams. The 2 µm trap beam is almost perpendicular to the probe beam and to the 1 µm trap beam. The 1 µm trap beam is almost parallel to the probe beam. Thus, the wings should be associated with atoms trapped only in the 2 µm trap beam.](image)

We observe about $4 \times 10^6$ atoms in the large volume crossed ODT. Approximately $1.1 \times 10^6$ of these atoms are located in the intersection region. We measure a temperature of the trapped atoms of about 28 µK. The atoms that are trapped in the intersection region, have a phase-space density of $4.3 \times 10^{-4}$, a peak density $3.6 \times 10^{13}$ cm$^{-3}$ and a collision rate $1.2 \times 10^3$ s$^{-1}$.

The small volume ODT ($w_0=15$ µm and $w_0=37$ µm) captures about $1.1 \times 10^6$ atoms, with $300 \times 10^3$ atoms in the intersection region. The trap frequencies are 1.2 kHz and 3.0 kHz, associated with 2 µm and 1 µm single beam transverse frequencies. The total trap depth is about 405 µK. The temperature of the atoms is about 38 µK, which is higher than the temperature in the large volume trap. The atoms in the intersection region have a phase-space density of $1.8 \times 10^{-3}$, a peak density $2.3 \times 10^{14}$ cm$^{-3}$ and a collision rate $9.2 \times 10^3$ s$^{-1}$.

We observe a substantial increase of the collision rate in the small trap at the cost of a smaller atom number. Therefore it confirms the statement that the choice of the waist of the trapping beams is a trade-off between a large atom number (large waist) and a high density (small waist) in the trap.
A pure 1 µm cross-beam trap has been made as well. Two 1 µm focused beams with orthogonal polarization were crossed in the center of the MOT. The beam waists were 18 µm and 50 µm, and the trap frequencies were 0.9 and 2.5 kHz. We trapped about 1.4 × 10⁶ of atoms (300 × 10³ in the intersection region), and we measured a surprisingly high temperature of about 60 µK. The atoms in the intersection region have a phase-space density of 2.6 × 10⁻⁴, a peak density 6.6 × 10¹³ cm⁻³ and a collision rate 3.3 × 10⁻³ s⁻¹. We did not succeed to decrease this temperature by changing the loading procedure. The temperature can be decreased by decreasing the trap depth, although this leads to a smaller number of atoms in the trap. So, we see that the density and collision rate in a pure 1 µm cross-beam trap are well comparable with the density and collision rate, which have been achieved in a multi-colour ODT. We conclude that the sign of the level shift of the excited level does not play a crucial role.

Also a pure 2 µm crossed beam trap has been made. Unfortunately we observe a strong heating of the atoms in the cross region. We believe that this is the result of an unstable interference pattern, which causes parametric heating. The heating rate depends on the relative polarization of the crossed beams. The use of orthogonal polarizations decreases the heating substantially. This supports the idea, that we deal with a moving fringes pattern, caused by interference. This was unexpected, since the specified linewidth of 1 nm should prohibit the interference in this trap. The high heating rate prohibited further use of a pure 2 µm crossed beam trap.

### 5.8 Summary and discussion

We summarize this chapter with an overview, where we present the type of the investigated traps, and their main characteristics. For the crossed beam traps, the quoted atom number \( N \) is the number in the intersection region.

<table>
<thead>
<tr>
<th>Type of ODT</th>
<th>( U_0 ), µK</th>
<th>( \bar{\omega}/2\pi ), Hz</th>
<th>( N )</th>
<th>( T ), µK</th>
<th>( n_0 ), cm⁻³</th>
<th>( \gamma ), s⁻¹</th>
<th>( \rho/\rho_c )</th>
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<tr>
<td>Single, 2 µm</td>
<td>330</td>
<td>700</td>
<td>1.6×10⁵</td>
<td>15</td>
<td>1.6×10¹³</td>
<td>410</td>
<td>5.1×10⁻⁴</td>
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<tr>
<td>Single, 1 µm</td>
<td>550</td>
<td>1.9×10³</td>
<td>1.3×10⁵</td>
<td>25</td>
<td>1.2×10¹⁴</td>
<td>3.9×10³</td>
<td>1.7×10⁻³</td>
</tr>
<tr>
<td>Cross, 1&amp;2 µm</td>
<td>360</td>
<td>650</td>
<td>1.1×10⁶</td>
<td>28</td>
<td>3.6×10¹³</td>
<td>1.2×10³</td>
<td>4.3×10⁻⁴</td>
</tr>
<tr>
<td>Cross, 1&amp;2 µm</td>
<td>405</td>
<td>2.2×10³</td>
<td>3×10⁵</td>
<td>38</td>
<td>2.3×10¹⁴</td>
<td>9.2×10³</td>
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<td>6.6×10¹³</td>
<td>3.3×10³</td>
<td>2.6×10⁻⁴</td>
</tr>
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</table>

There is a clear correspondence between a high trap frequency \( \bar{\omega} \) and a high spatial density \( n_0 \), and therefore also to a high collision rate \( \gamma \) and a high phase-space density \( \rho \). We do not see any strong advantages of 2 µm light, compare to 1 µm. Contrary to the argument given in [14] the sign of the shift of the excited state does not seem to play a crucial role. However, this large excited state shift pronounced itself by the maximum in the number of atoms loaded into a 2 µm single beam ODT as a function of trap depth. We do measure a lower temperature in a 2 µm single beam ODT and a crossed beam ODT with 1 µm and 2 µm beams than a 1 µm single beam ODT and a crossed beam ODT with two 1 µm beams. However
this temperature difference does not yield a crucial advantage for forced evaporative cooling.

We conclude that the small waist of trapping beams is a condition for achieving a high density of the atoms into an ODT. Indeed following [14] we can assume that the number of trapped atoms \( N \sim \omega_0^2 \), where \( \omega_0 \) is the beam waist. Assuming that the temperature does not depend on trap frequencies, for fixed trap depth we obtain a density \( n \sim \lambda/\omega_0^2 \). The collision rate and the phase-space density of the atoms is also proportional to \( \lambda/\omega_0^2 \). These scaling laws show an advantage for ODTs with a tight focus and a long wavelength. Although the presented scaling laws are an obvious simplification, they serve as a good first assumption and give the qualitative explanation of our experimental results.

We succeeded to reach an atomic density comparable to the atomic density in CO\(_2\) traps [10], [14]. However, we had to produce a trap with a substantial higher mean frequency (2.4 kHz compared to 1.5 kHz) to reach a comparable density \( 2 \times 10^{14} \text{ cm}^{-3} \) as reported in [10]. This statement qualitatively holds for both 1 \( \mu \text{m} \) and 2 \( \mu \text{m} \) single beam ODT compared to the result of [14] as well. Indeed the measured axial trap frequency in our 1 \( \mu \text{m} \) single beam ODT was measured to be 135 Hz in the axial direction. This is comparable to the axial frequency of 350 Hz, which was measured by Cennini [14]. The axial oscillation period is comparable with the dark molasses life-time (50 ms). Therefore we believe that a strong axial confinement is one of the crucial requirements to reach a high density in an ODT. This explains why CO\(_2\) traps [10] are more successful than near infrared traps with similar trap frequencies. The scaling laws presented above should be modified to take into account the role of the axial frequency. It makes the importance of a long wavelength and a small waist more pronounced. The high density can potentially cause a high three-body collisions rate. In this way it can lead to high losses of atoms. However, up to now we do not know of any experiment, where these high densities were actually obtained. Therefore, a small waist of the trap beam either in a single beam

or a crossed beam configuration is highly preferable. Still we have to stress that a small waist is unfavorable for the loading efficiency in terms of the fraction of MOT atoms loaded into the ODT. The use of a long wavelength (i.e. 10.6 \( \mu \text{m} \)), has a natural advantage of creating a strong longitudinal confinement with a relatively large trap volume. Another solution could be the use of near infrared beam with a large value of \( M^2 \gg 1 \), so a beam divergence is larger than diffraction limited.

The achievement of a tight focus can be complicated technical problem. The diffraction limit for a waist is expressed as: \( w_0 = (\lambda/\pi)(f/d) \), where \( f \) is the focal distance of the lens and \( d \) is the diameter of the beam on the lens. Therefore it is necessary to use a large lens with a short focal distance in order to obtain a small waist. However, this is difficult in reality, because the region around the MOT is crowded due to the presence of a vacuum system and MOT beams etc. Therefore a tight focusing of the trap beam is a key task, either for CO\(_2\) or for an infrared ODT.

In summary in near infrared ODTs one can reach the same numbers for the density, collision rate and phase-space density as in CO\(_2\) ODT, but the number of atoms will be substantially lower. There are a few ways to overcome this problem.
For example one can first evaporate in a magnetic trap, and then load the atoms into an ODT. This will be more efficient, because the size of the atomic cloud will be reduced. Another possibility is to use a large ODT to trap atoms from a MOT, and then to load atoms into a small volume ODT [17]. Finally, a very promising technique is to use a compressible optical dipole trap [15].
6 Evaporative cooling of atoms in optical dipole traps

We evaporate cold \(^{87}\text{Rb}\) atoms from various optical dipole traps (ODTs) with wavelengths of 1 and 2 \(\mu\text{m}\). Evaporative cooling was achieved by reducing the power of the trapping beams on various time scales. We investigate evaporative cooling by varying the trap geometry and the shape of the evaporation ramp. We cool atoms in ODTs to the nK range. We observe an increase of phase-space density by a few orders of magnitude. For a crossed-beam ODT, crossed in the vertical plane we achieve a phase-space density just above the BEC phase transition. We estimate a fraction of condensed atoms about 10 \%. We also observe a drop of the atomic collision rate during evaporation from the ODT due to decompression of the atomic cloud, which seems to be a serious drawback of direct evaporation from an ODT. We present a compressible ODT, where the trap volume can be changed during the evaporative cooling.
6.1 Introduction

The first and still most common method to produce a BEC is by forced evaporation from a magnetic trap. The main advantage of this method is that it allows to ramp down the trap depth at a constant trap frequency, which makes it possible to reach the run-away regime. The needed ratio of "good" to "bad" collision rates is typically about 300. In addition it is usually easier to load a large number of atoms into a magnetic trap, due to its large volume.

Nevertheless an all-optical method to produce a BEC remains attractive. This method allows trapping of species with a weak magnetic moment or with multiple spin states. Evaporative cooling can be made more efficient by tuning the atomic scattering properties using Feshbach resonances in a magnetic field. For example a Cs BEC was achieved by tuning the atomic scattering properties by a magnetic field [17]. This is not possible in a magnetic trap, and thus an ODT needs to be used.

The main drawback of evaporative cooling from an ODT is a decrease of the trap frequency due to a decrease of the trap depth. Therefore it is not possible to reach the run-away regime during evaporation from an ODT. This demands a much higher initial ratio of good-to-bad collision rates and also a high initial phase-space density. Nevertheless all-optical BECs were achieved by several groups.

The first evaporative cooling in an optical dipole trap was demonstrated in 1995 [9]. The authors trapped about $5 \times 10^3$ sodium atoms at a temperature of 140 $\mu$K in an optical dipole trap formed by the intersection of two 1.06 $\mu$m laser beams. They evaporatively cooled the atoms to a final temperature of 4 $\mu$K by reducing the trap depth during 2 seconds. The final number of atoms was 500. A 28 fold increase in atomic phase-space density was achieved in this work.

In 2001 Barrett et al. demonstrated the all-optical formation of an atomic Bose-Einstein condensate [10]. They employed a quasi-electrostatic trap (QUEST), which was created by two crossed, focused CO$_2$ laser beams. The authors trapped $2 \times 10^6$ $^{87}$Rb atoms at a temperature of 75 $\mu$K. Evaporative cooling through the BEC transition was achieved by lowering the power of the trapping beams over 2 s. The resulting condensates were $F = 1$ spinors with $35 \times 10^3$ atoms distributed between the $m_F = (-1, 0, 1)$ states.

Later, in 2003, Cennini et al. [14] produced an all optical BEC of $^{87}$Rb atoms by a simpler approach using only one focused CO$_2$ laser beam, which was oriented horizontally, and was focused to a waist of 27 $\mu$m. They started with forced evaporation of $4 \times 10^6$ atoms with a temperature of 140 $\mu$K. After evaporative cooling during 7 seconds, a spinor condensate was produced with $12 \times 10^3$ atoms distributed among the $m_F = (-1, 0, 1)$ states.

Recently, condensation of cold ytterbium atoms in an ODT was demonstrated [21] in a crossed-beam ODT. In this experiment the ODT beams were generated from a cw diode-pumped solid state laser with a wavelength of 532 nm and a power of 10 W. The horizontal and vertical beam powers were 5.0 and 0.2 W, and the beam waists were 14 and 11 $\mu$m respectively. The potential depths induced by the horizontal and the vertical beam (both were calculated from the beam waists and
the peak intensities) were 630 µK for the horizontal beam and 53 µK for the vertical beam. Evaporative cooling was carried out by gradually decreasing the power of the horizontal beam in 2 s while keeping the vertical beam power constant. Finally, a BEC of \(5 \times 10^3\) atoms of \(^{174}\text{Yb}\) was produced. An essential difference of this experiment is that the trap beams were not ramped down proportionally as it was done in \([9,10]\); here the optical power of the vertical trap beam was unchanged.

Up to now mostly \(^{12}\text{CO}_2\) lasers are used for the all-optical formation of a BEC of \(^{87}\text{Rb}\) atoms. The only exception is the result of Kinoshita et al. \([15]\). However, the authors had to change the volume of their ODT during evaporation in order to achieve a BEC. We think that this is a consequence of a higher loading efficiency of an 10 µm ODT compared to a near infrared ODT, as mentioned in the previous chapter. Nevertheless there is no fundamental reason that would forbid achievement of a BEC in a near infrared static ODT. The all-optical formation of a BEC of \(^{87}\text{Rb}\) in a near infrared ODT still seems to be an unresolved question.

In this chapter we describe our experiments on evaporative cooling in various ODT geometries. We start with a single beam ODT as simple example, and discuss the basic properties of evaporative cooling in ODTs. We continue with a few examples of crossed beam ODTs. Then we discuss a compressible ODT. We finish this chapter with a simulation of evaporation from an ODT and with a summary.

The main purpose of this chapter is to study the possibility of reaching BEC by evaporative cooling from an ODT. We will give simple recommendations for the geometry of ODTs and for evaporation from them. We determine a rough parameter, which predicts the final phase-space density achievable by evaporative cooling from the ODT. For more accurate calculations, a numerical model was developed. Finally we suggest few ways to improve evaporative cooling from an ODT.

### 6.2 Evaporation from a single beam optical dipole trap

We start this chapter by describing evaporation from a single beam ODT as a first simple example. Many important features of forced evaporation from ODTs are here more clear than in the case of a crossed beam ODT.

As described in the previous chapter, a 2 µm trap beam with a waist of about 22 µm and a beam power of 3 W provides trap frequencies \(\omega_r = 2\pi \times 2.5\) kHz, \(\omega_z = 2\pi \times 74\) Hz, with a trap depth of \(U/k_B=300\) µK. We loaded \(140 \times 10^3\) atoms into this trap, at a temperature of 13 µK. We calculated a collision rate (already defined in previous chapters as \(\gamma = (m\bar{\omega}^3\sigma/2\pi^2k_B)N/T\)) of \(\gamma = 510\) s\(^{-1}\). These temperature and density are reasonable parameters to start forced evaporation. Indeed, the ratio of good-to-bad collision rates is more than 2000.

 Forced evaporative cooling was performed by a gradual decrease of the optical power of the trap beam, and thus of the trap depth. It is obvious that the speed of the evaporation ramp should decrease during evaporation, in view of the decreasing collision rate. An optimal function of the trap depth vs. time is a subject of discussion. Barrett et al. \([10]\) used three linear ramps in the following way: immediately
after loading, the power in both trap beams is ramped to 1 W in 1 s (from an initial power of 12 W), then the power in both beams is ramped to the final value (typically about 260 mW) in 1 s and is maintained at this low level for 0.5 s.

O’Hara et al. [47] proposed scaling laws for the number of atoms, the collision rate, and the phase-space density as a function of the trap depth for evaporative cooling in an adiabatically lowered optical trap. This solution is valid for a specific type of function of trap depth vs. time and neglects atom losses due to background gas collisions. Unfortunately, we find that neglecting atom losses due to background gas collisions is unrealistic for our experimental parameters. Besides this there is no reason to consider a constant $\eta$ parameter as the most efficient way of evaporation from an ODT.

In practice we optimize our evaporation ramp experimentally. We tried to reach maximum evaporation efficiency, i.e. maximum final phase-space density.

In the ramp, which we present here, we chose a time-trace consisting of 2 decaying exponents with different $1/e$ decay times $\tau_1$ and $\tau_2$. We vary the ratio of amplitudes of these 2 exponents together with $\tau_1$ and $\tau_2$ to optimize our time-trace. This voltage time-trace is amplified and sent to a Pockels cell. Therefore, the optical power of the trap beam is proportional to $\sin^2\left(\frac{\pi}{4}V(t)/V_{\pi/4}\right)$, where $V(t) = V_{\pi/4}(a_1 \exp[-t/\tau_1] + a_2 \exp[-t/\tau_2])$. The sum $a_1 + a_2$ should be equal to 1. The $\pi/4$ voltage of the Pockels cell is 3.1 kV. The Pockels cell is used in double pass, thus the $\pi/4$ voltage corresponds to a $\pi/2$ rotation of the polarization. In practice the trap beams cannot be extinguished completely due to the finite polarization extinction, but the final trap depth is much smaller than the initial one and even smaller than the minimum depth needed to overcome gravity (thus atoms would finally fall out of the trap).

Figure 6.1: A sequence of absorption images of atomic clouds, each after 4 ms expansion. The size of each image is $3 \times 3$ mm$^2$. Due to expansion the vertical size of the clouds is defined by the temperature of the cloud. Each image corresponds to a different evaporation time (0, 1, 2, 5 and 9 seconds respectively). The ODT was formed by a single 2 $\mu$m beam, focused to a waist of 22 $\mu$m.

We demonstrate one evaporation ramp as an example. For this specific evaporation ramp we choose $V(t) = (0.25 \exp[-t/\tau_1] + 0.7 \exp[-t/\tau_2] + 0.05) \times 3.1$ kV, with $\tau_1 = 1$ s, and $\tau_2 = 3$ s; the small offset of 0.05 compensates for gravitational sag. In Fig. 6.1 we present this evaporation sequence of the expanded cloud of cold atoms. Atoms were evaporated from the ODT for various times. After evaporation the trap beam was turned off. The cloud was released and expanded for 4 ms. The number
6.2 Evaporation from a single beam optical dipole trap

Figure 6.2: Various parameters of the atomic cloud during evaporation from a 2 \( \mu \text{m} \) single beam ODT: a) temperature of the atomic cloud (filled triangles) and the number of atoms (open squares) vs. evaporation time. b) the ratio of the trap depth to the temperature: the \( \eta \) parameter vs. evaporation time. c) the collision rate (open squares) and the phase-space density (filled triangles) vs. evaporation time. d) the atom density vs. the temperature. The line indicates the BEC phase transition. We observe a slow decrease of number of atoms in the ODT from \( 140 \times 10^3 \) to \( 2 \times 10^3 \) and a fast cooling during the first 2 seconds to temperature of about 1 \( \mu \text{K} \). Further cooling is decelerated due to a drop of the atomic collision rate (down to few collisions per second).

of atoms was calculated from each absorption image. After 4 ms of free expansion the vertical size of the cloud is much larger than initially, so it is defined by the temperature of the cloud. The temperature was calculated from the fitted vertical diameter of the absorption image of the expanded cloud neglecting the initial size of the cloud and assuming a Boltzmann distribution. We observe a substantial cooling together with a loss of atoms. As one can see in Fig. 6.2 the temperature falls from 13 \( \mu \text{K} \) initially to a remarkable value of about 100 nK. Unfortunately the phase-space density increases only from \( 10^{-5} \) to \( 30 \times 10^{-5} \) (a phase-space density of 1 corresponds to the phase transition). We assume a Boltzmann distribution for calculation of the atomic densities in Fig. 6.2. For the sake of simplicity, we continue to use a Boltzmann distribution for the calculation of the atomic density in an ODT during this chapter unless otherwise indicated. The final atomic density
is $4.6 \times 10^{11}$ cm$^{-3}$, and the final trap frequencies are 235 and 4.8 Hz. Evaporative cooling became inefficient due to a low collision rate: we observe a striking drop of the atomic collision rate from $500 \text{ s}^{-1}$ to $1 \text{ s}^{-1}$. This drop of atomic collision rate is very typical for forced evaporation from an ODT, and was also reported by other groups [10, 14, 20, 48]. Another typical feature of evaporative cooling in an ODT is the decrease of efficiency of the evaporation during the evaporation ramp. Indeed, in the present example we observe an increase of the phase-space density from the $0.76 \times 10^{-3}$ to $5.7 \times 10^{-3}$ during the first second. At the same time the number of atoms decreases from $166 \times 10^3$ to $119 \times 10^3$. This corresponds to an evaporation efficiency $\ln[(\rho/\rho_0)/(N_0/N)]$ equal to 6. This is an impressive number, since the typical evaporation efficiency from a magnetic trap is about 2. This high initial evaporation efficiency can be explained by the high initial collision rate, and by "natural" evaporation. However, during the next four seconds the phase-space density increases only to $14.0 \times 10^{-3}$ (less than a factor of 3), while the number of atoms falls from $119 \times 10^3$ to $18 \times 10^3$ (more than a factor of 6). This corresponds to an evaporation efficiency of 0.5. Such a drop of efficiency during evaporation was observed in all ODTs that are discussed in this chapter. So it is highly favorable to have atoms in an ODT with a high initial phase-space density, so that one can reach the phase transition while evaporative cooling is still efficient.

This single beam experiment shows the importance of maintaining high collision rates and thus high trap frequencies. Therefore a crossed beam ODT seems to be more favorable for efficient evaporation compared to a single beam ODT. So, the next step is to perform evaporation from a crossed beam ODT.

### 6.3 Evaporation from a crossed beam optical dipole trap

Two crossed focused laser beams can create a strong confinement in all dimensions for cold $^{87}$Rb atoms. The high confinement in all dimensions is the main motivation for evaporation from a crossed beam optical dipole trap. We use a few configurations of focused 1 $\mu$m and a focused 2 $\mu$m beams to produce such a trap. In all cases the two trapping beams lie in a horizontal plane. We discuss evaporation from two 1 $\mu$m $\times$ 2 $\mu$m crossed beam ODTs and from a 1 $\mu$m $\times$ 1 $\mu$m crossed beam ODT.

#### 6.3.1 A large volume ODT, by crossing a 1 $\mu$m and a 2 $\mu$m beam

We start with a large volume crossed-beam trap, as described in the previous chapter. As was mentioned, approximately $10^6$ atoms are located in the intersection region, with an atomic temperature of about 28 $\mu$K. The trap frequencies in the cross region are 1.35 kHz and 0.53 kHz for beam powers of 19.6 W and 8.9 W for the 2 $\mu$m and 1 $\mu$m beams, respectively. The atoms in the intersection region have a phase-space density of $1.2 \times 10^{-3}$, a density of $9.2 \times 10^{13}$ cm$^{-3}$ and a collision rate of $3.1 \times 10^9$ s$^{-1}$. Our evaporation voltage ramp consists now of three exponents
6.3 Evaporation from a crossed beam optical dipole trap

plus a small offset to compensate for gravitational sag at the end of the evaporation. We keep the beam powers constant during the first 300 ms, in order to load atoms from the wings into the cross region. Then we start to evaporate with $\exp[-t/\tau]$ with $\tau=0.75$ s. After 1 second we switch to $\tau=1$ s, and finally after 1 more second of evaporation we switch to $\tau=1.5$ s, which is maintained till the end of the evaporation ramp. This signal is amplified and then sent to the Pockels cell. In Fig. 6.3 we present an evaporation sequence with subsequent expansion of clouds of cold atoms. The atoms were evaporated from the crossed-beam ODT for various times. Then the trap beam was turned off, and the cloud was released and expanded for 8 ms. Therefore the size of the atomic cloud is determined by the temperature of the cloud. Initially the cloud is horizontally elongated because of the width of the trap. Gradually the atoms leave the wings, and the expanded atomic cloud became round. After 3 seconds of evaporative cooling we end up with $7 \times 10^3$ atoms at a temperature of 500 nK, and a phase-space density of about $23 \times 10^{-3}$. The final trap frequencies are 305 and 123 Hz. The final density is $8.7 \times 10^{11}$ cm$^{-3}$, and the collision rate is then 5 s$^{-1}$. Further evaporation is useless due to the low trap frequencies and subsequently low collision rates. We present the dependence of various parameters of the cloud vs. evaporation time in Fig. 6.4.

We believe that the use of tighter focused beams can help to maintain the trap frequencies at a useful level. A distinctive feature of the crossed beam ODT is the existence of the "wings" of the trap. The wings of the ODT collect atoms during the MOT and dark molasses phase. These atoms are gradually transferred to the intersection region. However, during forced evaporation, atoms that leave the intersection region are still trapped in the wings of the ODT. Potentially these highly energetic atoms can come back to the intersection region and make the evaporative cooling less efficient. This could be the reason why we achieve a somewhat higher phase-space density in a single beam ODT, despite the fact that the mean frequencies of these two traps are similar.

Figure 6.3: A sequence of absorption images of atomic clouds each taken after 8 ms expansion. The size of each picture is $4 \times 4$ mm$^2$. The diameter of the clouds is determined by the temperature of the cloud. Each image corresponds to a different evaporation time (0, 1, 2 and 3.5 s respectively). The ODT was formed by crossing two beams (the waist of the 1 $\mu$m beam was 75 $\mu$m, and the waist of the 2 $\mu$m beam was 60 $\mu$m). One can see an obvious reduction of the cloud size, which is evidence for cooling.
Evaporative cooling of atoms in optical dipole traps

6.3.2 A small volume ODT, by crossing a 1 µm and a 2 µm beam

We continue with describing evaporation from a crossed-beam trap consisting of a 1 µm beam focused to a waist of 15 µm crossed with 2 µm beam focused to a waist of 37 µm (a small volume ODT). The trap frequencies in the cross region are $2\pi \times 1.2$ kHz and $2\pi \times 3.0$ kHz for beam powers of 5 W and 0.4 W for the 2 µm and 1 µm beams, respectively. The total trap depth is about 400 µK. Our evaporation ramp again consists of three exponents. We keep the beam powers constant during the first 300 ms. Then we start to evaporate with $\exp[-t/\tau]$ with $\tau=0.7$ s, after 1 second of evaporation we switch to $\tau=1$ s, and finally after 1 more second of evaporation we switch to $\tau=2$ s, and maintain this till the end of the evaporation ramp ($\sim 5$ seconds after the start of the evaporation). We present the dependence of various parameters of the cloud vs. evaporation time in Fig. 6.5. In this particular sequence
6.3 Evaporation from a crossed beam optical dipole trap

**Figure 6.5:** Various parameters of the atomic cloud during evaporation from a small volume crossed-beam trap: a) temperature of the atomic cloud (filled triangles) and the number of atoms (open squares) vs. evaporation time. b) the \( \eta_1 \) parameter (open squares) and the \( \eta_2 \) parameter (filled triangles) vs. evaporation time, c) collision rate (open squares) and phase-space density (filled triangles) vs. evaporation time. d) the atom density vs. the temperature. The line indicates again the phase transition.

We observe an increase of the phase-space density from \( 2.0 \times 10^{-3} \) to \( 72.0 \times 10^{-3} \). Just as in the previous section the evaporation became inefficient due to a low collision rate at the end of the evaporation ramp (\( \gamma = 7 \) s\(^{-1} \)). At the end of evaporation ramp we observe \( 2 \times 10^3 \) atoms at a temperature of 300 nK and a density of \( 1.9 \times 10^{12} \) cm\(^{-3} \). The final trap frequencies are \( 2\pi \times 320 \) and \( 2\pi \times 94 \) Hz.

Comparing the results of evaporation from large and small volume ODTs, we learn that evaporative cooling in a small crossed beam ODT is substantially more efficient than evaporation from a large crossed beams ODT. Indeed we achieve a higher phase-space density (0.072 vs. 0.023) and a lower temperature (300 nK vs. 500 nK).

### 6.3.3 ODT created by crossing two 1 \( \mu \)m beam

We also studied evaporation from a pure 1 \( \mu \)m crossed-beams ODT. Two 1 \( \mu \)m focused beams with orthogonal polarization were crossed in the center of the MOT.
Evaporative cooling of atoms in optical dipole traps

Figure 6.6: Various parameters of the atomic cloud during evaporation from a pure 1 μm crossed-beams ODT: a) the temperature of the atomic cloud (filled triangles) and the number of atoms (open squares) vs. evaporation time. b) the $\eta_1$ parameter (open squares) and the $\eta_2$ parameter (filled triangles) vs. evaporation time. c) collision rate (open squares) and phase-space density (filled triangles) vs. evaporation time. d) the atom density vs. the temperature. The line indicates the BEC phase transition.

The beam waists were 18 μm and 50 μm, and the trap frequencies were $2\pi \times 2.5$ and $2\pi \times 0.9$ kHz for beam powers of 0.6 and 4.7 W, respectively. The total trap depth was 360 μK. We trapped about $300 \times 10^3$ atoms in the intersection region, at a temperature of about 55 μK. The atoms had an initial phase-space density of $2.6 \times 10^{-4}$. As in the previous subsection we start the evaporation ramp with a voltage ramp of $\exp[-t/\tau]$ with $\tau=0.7$ s, after 1 second we switch to $\tau=1$ s, and finally after 1 more second of evaporation we switch to $\tau=2$ s, and this we maintain until the end of the evaporation ramp. The dependence of the various measured parameters of the cloud vs. evaporation time are shown in Fig. 6.6. After 4 s of evaporative cooling we observe about $10 \times 10^3$ atoms at a temperature of 500 nK. The final density and the final collision rate are $4.4 \times 10^{12}$ cm$^{-3}$ and 21 s$^{-1}$ respectively. The final trap frequencies are 312 and 93 Hz. The maximum achieved phase-space density equals $37 \times 10^{-3}$. Thus, we reach a phase-space density that is comparable to evaporative cooling in a 1 μm × 2μm ODT. This suggests that the
wavelength of the trap beam does not play an important role during the evaporation and that rescattering of trap light is still negligible, even for a 1 µm beam.

Despite the improvement compared to evaporation from a single beam ODT, the drop of collision rate during evaporation remains a crucial problem for the crossed beam ODTs. With the small volume ODT we achieve better results, compared to the large volume ODT, but a drawback is that the number of trapped atoms is smaller. In the following, it is crucial to solve the drop of the evaporation efficiency at the end of the evaporation ramp.

6.4 Evaporation from an ODT crossed in the vertical plane

As was already stated, decompression of the atomic cloud is the main drawback of evaporation from an ODT. In principle one can reduce the power of only one trap beam, as was done in [21]. This allows to keep two trap frequencies at the initial value and therefore to maintain a relatively high collision rate until the end of the evaporation ramp. However, if both beams lie in the horizontal plane, such a reduction will not have much effect, because the evaporated hot atoms will still be trapped in the wings of the unreduced beam. Therefore, to perform evaporation we aligned the unreduced beam vertically. In this case gravity acts along the weak axis of the evaporated trap and atoms can leave the trap completely [21]. Two trap frequencies will be maintained, so the drop of collision rate will be reduced. Theoretically, it is possible to reach the run-away regime in this geometry. Indeed for two fixed frequencies, we modify the scaling laws written in a chapter 2 and obtain $\gamma \sim N/\sqrt{T}$. So, a decrease of the temperature can compensate a decrease of the number of the trapped atoms, as it happens in a magnetic traps. An essential drawback is that the $\eta$ parameters in two directions will increase and suppress evaporation in two dimensions. Essentially we deal with 1-D evaporation in this case.

One can also decrease the optical power of the trap beams at the same speed. Still the vertically crossed beams ODT has the advantage compared to a horizontally crossed beams ODT. During forced evaporation, due to gravity, more atoms that leave the intersection region leave the wings of the ODT. We describe first an example where both the horizontal and the vertical beam are reduced with the same speed. Then we describe evaporation from a vertically crossed beams ODT, where the power of the vertical beam is reduced much more slowly than the horizontal beam.
Evaporative cooling of atoms in optical dipole traps

Figure 6.7: An absorption picture of atoms just released from the ODT 100 ms after loading (no forced evaporative cooling). Most of the atoms have already been collected into the deeper potential of the intersection region due to a high trap frequencies.

6.4.1 Evaporation from an ODT with beams crossed in the vertical plane: identical ramps for the power of both beams

The next crossed beam ODT consists of a horizontal 1 µm trap beam, crossed with a near vertical (∼ 20° from vertical) 2 µm beam. The 1 µm beam is focused to a waist of 14 µm and the 2 µm beam is focused to a waist of 15 µm. We load about $110 \times 10^3$ atoms into this ODT. This trap has the advantage of higher initial trap frequencies compared to the previous one, but the drawback of a smaller volume. The initial frequencies are 5.6 and 4.5 kHz in the horizontal and vertical direction, respectively.

We reduce both the horizontal and vertical beam exponentially with a $1/e$ time $\tau = 2$ s. In Fig. 6.8 the evolution of the cloud parameters during evaporation is shown. After 5.5 s of evaporative cooling we have about $1.0 \times 10^3$ remaining atoms. The final horizontal trap frequency is 650 Hz and the final vertical trap frequency is 560 Hz. On basis of the final frequencies and the number of atoms we estimate a critical temperature of 170 nK ($T_c = \frac{\sqrt{N}}{3 \times 1.2} \hbar \bar{\omega}/k_B$). After 3 ms of expansion the cloud demonstrates a clear anisotropy. We estimate an average kinetic energy of 185 nK from the vertical expansion, and 250 nK from the horizontal expansion. A calculation of the mean-field energy gives 120 nK (assuming 100% BEC). The chemical potential is calculated as: $\mu = (15 \hbar^2 \sqrt{mN \bar{\omega}^3} a)^{2/5}/2$, where $a$ is the scattering length, $m$ is the mass and $\bar{\omega}$ is the mean trap frequency. We clearly approach the phase-transition very closely in this evaporation ramp. The calculated final phase-space density is 0.42.
6.4 Evaporation from an ODT crossed in the vertical plane

Figure 6.8: Various parameters of the atomic cloud during evaporation from the ODT, which consists of a horizontal 1 \( \mu m \) trap beam, crossed with a near vertical 2 \( \mu m \) beam: a) the temperature of the atomic cloud (filled triangles) and the number of atoms (open squares) vs. evaporation time. b) the \( \eta_1 \) parameter (open squares) and the \( \eta_2 \) parameter (filled triangles) vs. evaporation time. c) collision rate (open squares) and phase-space density (filled triangles) vs. evaporation time. d) the atom density vs. the temperature. The line indicates the phase transition.

6.4.2 Evaporation from an ODT with beams crossed in the vertical plane: power reduction of the horizontal beam much faster than of the vertical beam

Our final ODT consists of a horizontal 2 \( \mu m \) trap beam, crossed with a near vertical 1 \( \mu m \) beam. The 1 \( \mu m \) trap beam is aligned about 20° from the vertical direction. Both beams are focused to a waist of 22 \( \mu m \). We load about 150 \( \times 10^3 \) atoms into this ODT. Initially almost all atoms are trapped in the horizontal beam and only a small fraction is located in the intersection region. The atoms gradually fill the region of the intersection of the beams during the ramp down of the horizontal beam power. The loading of the central intersection region (see Fig. 6.9) increases both the phase-space density and the collision rate. Unfortunately, this process continues only as long as atoms in the wings are present. This loading process is demonstrated in Fig. 6.10 in terms of the number of atoms in the wings and in the intersection part
(the "dimple"). This closely resembles the loading of a so called "dimple" [78], where the atoms gain phase-space density by loading from a large trap into a small volume with a lower potential energy. Stamper-Kurn et al. [78] achieved a large increase of the phase-space density in the "dimple" due to a large ratio of the volume of the magnetic trap and the volume of the "dimple". In our case the ratio between the "dimple" volume to the trap volume is much smaller, and therefore the potential gain in a phase-space density is smaller.

Figure 6.9: Absorption pictures of atoms just released from the ODT, without any expansion. The size of each image is 2×2 mm². Each image corresponds to a different evaporation time (0, 1, 2, 3, and 4 s, respectively). The power of the horizontal beam is ramped down exponentially, with a 1/e time \( \tau \) of 1 second. The power of the vertical beam is left unchanged.

![Absorption pictures](image)

Figure 6.10: Number of atoms in the wings (open squares) and in the intersection part (filled triangles) of the ODT. The power of the horizontal beam was reduced exponentially, with a 1/e time \( \tau \) of 1 second. The power of the vertical beam was left unchanged. Despite a decrease in the number of atoms in the wings, the number of atoms in the intersection region initially grows (due to the loading from the wings), but starts to decrease when most of the atoms in the wings are gone.

![Number of atoms](image)

We achieve the best results by reducing the horizontal beam power exponentially with a 1/e time \( \tau \) of 1 second and by slowly reducing the vertical beam exponentially with a 1/e time \( \tau \) of 5 seconds. In Fig. 6.11 a decrease of the temperature of the trapped atoms and an increase of the phase-space density is visible. A few things are worth to note in these results. A distinctive feature of panel c of this plot
6.4 Evaporation from an ODT crossed in the vertical plane

![Graphs showing various parameters vs. evaporation time]

**Figure 6.11:** Various parameters of the atomic cloud during evaporation from the ODT, which consists of a horizontal 2 µm trap beam, crossed with a near vertical 1 µm beam: a) temperature of the atomic cloud (filled triangles) and the number of atoms (open squares) vs. evaporation time. b) the $\eta_1$ parameter (open squares) and the $\eta_2$ parameter (filled triangles) vs. evaporation time. c) collision rate (filled triangles) and phase-space density (open squares) vs. evaporation time. d) the atom density vs. the temperature (filled squares). For the points above the phase-transition line we applied the model for extraction of the temperature (open triangles), which takes into account the quantum statistics nature of the density distribution, and which is described below. The line indicates the BEC phase transition.

is that the collision rate is maintained at a level close to the initial one. This is achieved by a relatively slow decrease of the near vertical 1 µm trap beam, which leads to an increase of the $\eta_1$ parameter. The phase-space densities of the two last data points are above the transition line, which implies the existence of a condensed fraction in the atomic cloud. We also observe a clear anisotropy for the last two images of the evaporation sequence, which is another sign for condensation of the cloud. For example we observe about 700 atoms in the last absorption image. We estimate an average kinetic energy of 130 nK on basis of the vertical cloud size, and an average kinetic energy of 190 nK on basis of the horizontal cloud size. These temperatures are smaller than the calculated critical temperature ($T \approx 260 \text{nK}$) for the last absorption image.

However, the extraction of the temperature directly from the Gaussian width
leads to an error if the cloud is partially condensed. Indeed in this case a condensed fraction of the cloud (which expands slowly) is mixed with a thermal fraction (which expands with the normal thermal velocity), so the temperature is underestimated and the phase-space density is overestimated. Therefore, for a more detailed study we developed a model, which numerically calculates the expected size of an atomic cloud consisting of a thermal and a condensed part. For a given temperature and number of atoms the model calculates the fraction of condensed atoms. We assume a Thomas-Fermi profile for the condensed part. For both the Thomas-Fermi and the thermal component we calculate the expanded density distribution. We calculated the position of the minimum of the trap potential including gravity and we calculate the trap frequency at this point. These frequencies are used for calculation of the density profiles. We add the calculated density profiles of the thermal and condensed fractions. The result is then sampled in order to incorporate the finite pixel size of the camera. Finally a Gaussian is fitted to this calculated density distribution (just as it was done with experimental data). In this way using the number of the atoms, the trap frequencies, the temperature and the expansion time as input parameters we obtain a calculated expected Gaussian width in both the horizontal and vertical direction after a certain expansion time. We show an example in Fig. 6.12, where calculated points are fitted with a Gaussian. This procedure allows us to estimate a small condensed fraction, which is difficult to do directly from an absorption image due to the small number of atoms.

![Figure 6.12: The calculated column density (dots) of a partially condensed atomic cloud is fitted by a Gaussian profile. The values of column density are calculated using a numerical model, described in the text, which assumes a bimodal distribution.](image)

We find the number of atoms from an absorption image and mimic the fitted size of the cloud using this model. We use this model for the last two absorption images of the evaporation ramp (Fig. 6.11). We choose the temperature such that the calculated aspect ratio of the cloud coincides with the observed aspect ratio. This way, we extract the temperature of the cloud for the last image of 252 nK. The calculated aspect ratio for this temperature is about 1.23, consistent with the observed aspect ratio of 1.22. The fraction of condensed atoms is about 10 % for this temperature.
6.4 Evaporation from an ODT crossed in the vertical plane

6.4.3 Anisotropic expansions: time of flight series and discussion of the results

For a more detailed study we perform additional time of flight measurements (see Fig. 6.13 and Fig. 6.14). We use the same evaporation ramps as in the previous subsection. After 5 s of evaporative cooling we end up with about $1.5 \times 10^5$ atoms on average which demonstrate an anisotropic expansion. We observe an expansion in the vertical direction of 3.8 mm/s (which corresponds to a kinetic energy of 150 nK) and an expansion in the horizontal direction of 6.4 mm/s (which corresponds to a kinetic energy of 430 nK); both are shown in Fig. 6.13. The final horizontal trap frequency is 1.31 kHz and the final vertical trap frequency is 0.18 kHz. We estimate the final collision rate after 5 s of evaporation to be about 200 s$^{-1}$. Therefore, this trap does not suffer from a dramatic decrease of the collision rate due to decompression. For these frequencies and number of atoms ($1.5 \times 10^5$) we estimate a critical temperature of 340 nK, which is higher than a critical temperature in a previous example due to a higher final number of the atoms in the trap.

![Figure 6.13](image_url)

**Figure 6.13:** After 5 seconds of evaporation the atomic cloud was released from the ODT and expanded for a variable time. The size of the atomic cloud was measured by taking an absorption image of the cloud. For each image we obtain the cloud diameter by fitting in the vertical (open squares) and horizontal (filled triangles) directions. From a fit to the cloud diameter vs. expansion time we deduce an apparent temperature. We observe a difference in vertical and horizontal expansion. From the vertical expansion we deduce the kinetic energy of 150 nK, and from the horizontal expansion we deduce the kinetic energy of 430 nK.

We took the last four images of this time-of-flight sequence and we use the model (described in the previous subsection) to extract the temperature and number of atoms in the condensed fraction. We choose the temperature such that the calculated aspect ratio of the cloud coincides with the observed aspect ratio. Then we compare the calculated and observed width in the vertical direction. For all four absorption images a temperature of 325 nK gives a reasonable prediction for the aspect ratio. For this temperature the condensed part consists of 140 atoms, and 1360 atoms in
the thermal part. At the same time the observed vertical diameter (as well as the horizontal) is slightly larger than calculated for this temperature.

The probable reason is that the observed clouds have a size comparable with the resolution of our probe system. Indeed the resolution of our probe system is restricted by the pixel size (9 \( \mu m \)) and a diffraction limit. The calculation of the resolution for a diffraction-limited lens with a Rayleigh criterion \([81]\) gives \( r = 0.61 \lambda / N A = 3.5 \, \mu m \), which defines an upper limit of the resolution of our probe system. This leads to an overestimate of the cloud size. A table with calculated (calc.) and observed (obs.) values is shown in Fig. 6.15. We again observe that the atomic cloud is partially condensed. This time of flight series supports conclusion of the previous subsection, we observe the atomic cloud just above the BEC phase transition with a fraction of condensed atoms about 10 %. The total number of involved atoms is small, which explains why we do not observe a bimodal distribution directly.

<table>
<thead>
<tr>
<th>expansion time, (ms)</th>
<th>( \sigma_h/\sigma_v ), obs.</th>
<th>( \sigma_h/\sigma_v ), calc.</th>
<th>( \sigma_v ), obs. (( \mu m ))</th>
<th>( \sigma_h ), calc. (( \mu m ))</th>
</tr>
</thead>
<tbody>
<tr>
<td>3</td>
<td>1.36</td>
<td>1.18</td>
<td>16.0</td>
<td>13.9</td>
</tr>
<tr>
<td>4</td>
<td>1.18</td>
<td>1.25</td>
<td>20.3</td>
<td>16.8</td>
</tr>
<tr>
<td>5</td>
<td>1.30</td>
<td>1.30</td>
<td>26.5</td>
<td>19.6</td>
</tr>
<tr>
<td>6</td>
<td>1.45</td>
<td>1.32</td>
<td>28.0</td>
<td>22.4</td>
</tr>
</tbody>
</table>

As a final remark of this section we note that evaporative cooling in a crossed beam ODT, with beams crossed in the vertical plane, has the advantage of keeping the collision rate at the useful level. This allows us to perform efficient evaporation till the end of the ramp. Using the vertical crossed beam trap we finally reached phase-space density just above critical. However it does not solve the problem of the small number of atoms at the end of the evaporation ramp. This makes it difficult to accurately extract final temperatures and to make solid conclusions.
6.5 Simulation of forced evaporative cooling from an ODT

As was mentioned in the second chapter, the change of the number and the temperature of atoms in the trap can be written as [44]:

\[ \dot{N} = -N \Gamma_{ev} - N \Gamma_{vac} \]  
\[ \dot{T} = -\Gamma_{ev} \frac{\eta}{3} T + \frac{\dot{\omega}}{\omega} T, \]

(6.1)

(6.2)

where \( \Gamma_{ev} \) is the evaporation rate, \( \Gamma_{vac} \) is the background vacuum collision rate and \( \bar{\omega} = (\omega_x, \omega_y, \omega_z)^{1/3} \) is the mean trap frequency. Note that the last term in Eq.(2) describes adiabatic cooling due to decompression. This term is usually absent when evaporation is performed from a magnetic trap. However, this does not lead to an increase of phase-space density, since adiabatic cooling is combined with a decompression.

The evaporation rate can be written as

\[ \Gamma_{ev} = \gamma \eta e^{-\eta}, \]

(6.3)

where \( \gamma \) is the atomic collision rate, and \( \eta \) is the ratio of the trap depth and the atomic temperature.

It is not possible to find an analytical solution for the number of atoms and the atomic temperature during forced evaporative cooling from an ODT based on these equations (unless one is willing to neglect the vacuum background collision term). However a numerical solution of this problem is rather straightforward, and it can include terms describing three-body recombination and parametric heating. Including the influence of the wings of the crossed-beam ODT is the main problem which we face during this calculation. Another delicate problem is the role of gravity, which we will discuss now.

6.5.1 The role of gravity during the evaporation from an ODT

In general, gravity leads to a decrease of the trap depth. For low beam power it can open the trap. Because of this, the majority of the groups tries to align the trap beams (or at least one of them) strictly horizontally.

It is easy to calculate the offset due to gravity in the center of a single horizontal beam ODT. The potential along the vertical axis is \( U(z) = -U_0 \exp[-2z^2/w^2] - gmz \), where \( U_0 \) is the trap depth of the ODT in the absence of gravity, \( w \) the waist of the beam, and \( m \) the mass of a \(^{87}\)Rb atom. The first zero of the first derivative corresponds to the trap minimum. The derivative in the inflection point must be positive for a non-zero trap depth, yielding \( U_0 > \sqrt{\varepsilon m g w/2} \). Therefore the critical trap depth (which defines the critical power) is \( U_{0cr} = 0.82 gmw \). Note that this
is calculated strictly in the center of the trap. If an atom moves away horizontally from the center the offset becomes larger due to a larger waist of the trap beam.

![Graph showing potential vs. z.](image)

**Figure 6.16:** The potential of an optical trap in combination with gravity. There is a visible shift from the center of the ODT (which corresponds to $z = 0$). The numbers refer to an ODT created by a single 2 $\mu$m beam with a waist of 22 $\mu$m and a power of 50 mW.

As we can see, gravity decreases both the trap depth and trap frequency of an ODT. However it is easy to calculate that the change of trap depth is large compared to the change of frequency. For our example of a 2 $\mu$m beam of 50 mW focused to a waist of 22 $\mu$m we calculated that the trap depth decreases due to gravity from 5.6 $\mu$K to 3.7 $\mu$K; the trap frequency decreases negligibly from 331 Hz to 328. Although one can ask how valid the assumption of a harmonic trap is for a large gravitational sag, it seems that gravity helps to maintain the trap frequency in combination with a decreased trap depth, similar to the RF-knife for a magnetic trap. Therefore we believe that gravity has a positive (albeit uncontrolled) effect on evaporative cooling.

### 6.5.2 Numerical results

We present numerical results for evaporation from a single beam ODT and we compare these with experimental results presented in the second section of this chapter.

The numerical curves are shown in Fig. 6.17. These numerical calculations give us reasonable predictions. However, we observe two differences. First, it seems that in reality we have a slightly stronger evaporative cooling than we obtain in the calculations. This leads to a faster reduction of temperature and number of atoms. It looks like the trap frequencies were underestimated by about 20 %. We measure the trap frequencies by modulating the trap power and observing the trap losses (described in chapter 3). In practice the atoms are distributed along the trap beam and thus, they radially oscillate with a lower local frequency due to a larger beam waists. This statement is true both for evaporation and for the trap modulation experiment. However, in the trap modulation experiment, the atoms are hotter. Consequently, they are more distributed along the beam and a larger fraction sees a
larger waist of the trap beam. This might explain the deviation from the numerical model.

Second, in reality we reach substantially lower final temperatures than the model predicts. We believe that this is a result of the influence of gravity. The gravitational offset is calculated in the center of the trap. Further from the center the role of gravity increases due to a larger beam waist. If an atom has enough energy, it can reach positions where gravity is stronger than the trap confinement. Therefore, we believe that we underestimate role of the gravity during the calculations. Another point is that the results can be sensitive to a small tilt of the trap with respect to the horizontal axis.


6.6 Compressible optical dipole trap and evaporation from it

During the loading stage we prefer to have a large volume ODT. Because of the larger overlap with the MOT more atoms are collected into the trap. However, during the evaporation stage the situation changes to the opposite: a small volume ODT is preferable, because it naturally leads to high trap frequencies and an efficient evaporation. This dilemma can potentially be solved by using a compressible ODT, i.e. an ODT where the waists of the trap beams are changing during the experiment. Thus the decrease of the trap frequencies due to reduction of the beam power is compensated by a rise of the trap frequencies due to a decrease of the trap volume. Recently, this strategy led to the achievement of a large BEC [15].

6.6.1 Compression of the cold atoms in optical dipole trap

In Fig. 6.18, we present our setup, that can perform this task. Indeed, as we know, the diffraction limit of a focused Gaussian beam varies inversely proportional to the beam diameter on the last lens. This is directly obvious from the formula for the diffraction limit of a waist: \( w_0 = \lambda \pi (f / d) \), where \( \lambda \) is the wavelength, \( f \) the focal length, \( d \) the diameter of the beam, and finally, \( w_0 \) the waist.

\[ \text{Figure 6.18: Scheme of a single beam compressible ODT. The following abbreviations are used: } L_1 \text{ and } L_2 \text{ are lenses of the telescope, } L_3 \text{ is the final lens, which focuses the beam into the MOT, the actual target of the focused trap beam. Compression is achieved by moving one of the lenses of the telescope (in our case the second lens } L_2, \text{ which was mounted on a motorized translation stage). This leads to a change of the beam diameter on the last lens, and a subsequent change of the minimum waist of the trap beam, i.e. change of the ODT volume.} \]

The obvious drawback of this method is that a movement of the lens of the telescope causes a movement of the focus of the trap beam and thus, a movement of the ODT itself. This phenomenon is unavoidable, but the obvious way to reduce this problem is to make the distance between the telescope and last lens much larger than the length of the telescope. However, a large distance between the telescope and the last lens makes alignment much more delicate. Since the telescope was not aligned ideally, a displacement of the telescope lens causes a small change of the
6.6 Compressible optical dipole trap and evaporation from it

beam direction, and thus a displacement of the beam position on the last lens. In practice we find it difficult to work with distances between telescope and last lens larger than 1.5 m.

Generally speaking the movement of the ODT is a complicated function of the change of the telescope lens. It can be either larger or smaller than the actual movement of the telescope lens. Based on a thin lens approximation we get the following expression:

\[ p(x) = \frac{f_3}{1 - \frac{f_3(x - f_1 - f_2)}{L(x - f_1 - f_2) - f_2(x - f_1)}} \] (6.4)

where \( p(x) \) is the position of the ODT with respect to the last lens, \( f_1 \) and \( f_2 \) are focal lengths of the lenses \( L_1 \) and \( L_2 \) of the telescope, \( x \) is the distance between \( L_1 \) and \( L_2 \), \( f_3 \) is the focal length of the last lens \( L_3 \), and \( L \) is the distance between lenses \( L_2 \) and \( L_3 \). One can see that if \( x = f_1 + f_2 \), i.e. when the beam is collimated between \( L_2 \) and \( L_3 \), \( p = f_3 \).

![Image of absorption images](image)

**Figure 6.19:** Absorption images (1 × 5 mm²) of the atomic cloud during compression. The first picture is an image of the uncompressed cloud. The following 4 pictures correspond to a movement of the telescope lens \( L_2 \) by 0, 1, 2, 3 mm. It is noticeable that initially the movement of the focus of the last lens is large enough to compensate movement of the last lens, but gradually compensation became insufficient, and the cloud moves from its initial position. During this sequence we observe an axial compression of the cloud from 49.5 to 9.0 pixels.

The movement of the ODT is a serious problem. We observe a movement of
about 10 mm for a change of the beam waist of a factor of two (from ~ 50 µm to ~ 25 µm) for a telescope with a length of 15 cm, and a distance between telescope and last lens $L_3$ of about 1.5 meters (in agreement with Eq. (6.4)). Such a movement moves the atomic cloud out of view for the camera, and subsequently to the wall of the vacuum chamber. Besides this, the acceleration of the ODT spills atoms out of the trap. In order to solve this problem we mechanically attach a cable between the mounts of the lenses $L_2$ and $L_3$. Thus the last lens follows the movement of the telescope lens and compensates the movement of the ODT.

We reduced the length of the telescope in order to perform faster compression. The first lens has a focal distance of 5.0 cm and the second lens has a focal distance of 2.5 cm. For this choice we can find a region of the telescope lens position where the movement of the ODT can be roughly compensated by a movement of the last lens. In this particular sequence we change the beam waist from an initial value of 50 µm to a final value of 16 µm. The trap volume $(w_0^2 \times z_R)$ decreases by a factor of almost 100, by moving the 2nd lens $L_2$ of the telescope by 4.5 mm. Due to the finite speed of the motorized translation stage it takes about 2 seconds to perform this compression. The initial cloud consists of about $900 \times 10^3$ atoms at a temperature of about 20 µK. We combine the compression with a reduction of the trap beam power in order to avoid adiabatic heating and to keep $\eta$ at a useful value. We reduced the power of the trap beam from an initial value of 8 W to 780 mW. The trap depth after compression is equal to 161 µK, which is roughly equal to the trap depth before compression (170 µK). We finish the compression with about $173 \times 10^3$ atoms at a temperature of 7.9 µK. Thus, we loose about 80 % of the atoms during the compression stage. However the phase-space density during the evaporation increases from roughly $20 \times 10^{-6}$ to $2.2 \times 10^{-3}$. Therefore, we observe an efficient evaporation during the compression stage. The efficiency of evaporative cooling $\alpha$ is $2.8 (\alpha = \ln[(\rho/\rho_0)/(N_0/N)])$. This is a high efficiency compared to the typical value of 2 for evaporation from a magnetic trap. The atomic density changes from $8.3 \times 10^{11}$ cm$^{-3}$ to $2.2 \times 10^{13}$ cm$^{-3}$. This is a rather moderate increase of the density, despite the substantial drop of the trap volume. The increase of the density is suppressed by a drop of the number of atoms and decrease of the power of the trap beams. In addition the collision rate increase from 24 s$^{-1}$ to 370 s$^{-1}$.

### 6.6.2 Evaporation from the compressible optical dipole trap

Evaporative cooling from the compressible ODT after the compression is performed by reducing the optical power of the trap beams as described in previous sections. We observe a gradual cooling of the atomic sample.

We chose a time-trace consisting of 2 decaying exponents with different 1/e decay times $\tau_1$ and $\tau_2$. The optical power of the trap beam is proportional to $\sin^2(a_1 \exp[-t/\tau_1] + a_2 \exp[-t/\tau_2])$. For this specific evaporation ramp we choose $V(t) = (0.25 \exp[-t/\tau_1] + 0.75 \exp[-t/\tau_2]) \times 3.1$ kV, with $\tau_1 = 1$ s, and $\tau_2 = 3$ s. The evaporative cooling during 14 s (including 2 s of compression) leads to a sample of $12 \times 10^3$ atoms at a temperature of 240 nK and phase-space density of $64 \times 10^{-3}$. The final trap frequencies equal $2\pi \times 435$ and $2\pi \times 12$ Hz, the final density equals
6.7 Summary and discussion

We have studied evaporative cooling from ODTs of various geometries. Evaporative cooling from an ODT suffers from a drop of the collision rate due to decompression

4.1 \times 10^{12} \text{ cm}^{-3}\), and the collision rate equals 12 s$^{-1}$.

We conclude that, the evaporative cooling from a compressible ODT can be efficient during the compression. After compression an evaporative cooling suffers from a decrease of the collision rate, as in a normal ODT. Generally evaporative cooling from a compressible ODT seems to be a promising way to efficiently increase the phase-space density in an ODT. In our case further compression requires a further decrease of the waist of the trap beam, and thus further increase of the trap beam diameter on the last lens. This is not possible due to the finite size of the last lens and mirrors. Another solution might be to make a compressible crossed beams ODT, as was done in Ref. [15].

Figure 6.20: The various parameters of the atomic cloud during evaporation from the compressible ODT: a) the temperature of the atomic cloud (filled triangles) and the number of atoms (open squares) vs. evaporation time. b) ratio of trap depth to the temperature the $\eta$ parameter (open squares) vs. evaporation time. c) collision rate (open squares) and phase-space density (filled triangles) vs. evaporation time. d) the atom density vs. the temperature, the line is the BEC transition.
during the evaporation, which is a substantial drawback of evaporative cooling in an ODT. Evaporative cooling from a tight ODT is more efficient than evaporative cooling from a shallow ODT, because this trap has higher frequencies and thus can longer hold the collision rate on a useful level. Therefore we conclude that a high initial collision rate in combination with a high initial phase-space density is a crucial condition for successful evaporation from an ODT. A way to avoid a fast decompression is to cross the trap beams in the vertical plane and to keep the power of the vertical beam high. This allows to maintain the collision rate at a useful level for a longer time.

Straightforward evaporation from an ODT aiming to achieve degeneracy suffers from ”an illness of smallness”. In order to reach good initial parameters, such as a high collision rate and a high phase-space density, one has to focus the trap beam (or beams) tightly. This leads to a small initial number of atoms in the trap. This spoils the signal-to-noise ratio of the absorption images. It does not prohibit the achievement of quantum degeneracy using evaporative cooling, but does lead to a small final number of atoms.

We did not find in literature any parameter, that can tell if degeneracy is reachable by straightforward forced evaporation from an ODT or not. As a rule of thumb, we propose the following. We seem to win about one order of magnitude of phase-space density from evaporation from an ODT, by giving up about one order of magnitude in the collision rate. At some critical value of the collision rate, forced evaporation can no longer be efficient, since the loss of phase-space density due to vacuum background collisions is larger than the gain in phase-space density due to evaporation. This critical value of the collision rate can be numerically calculated using Eqs. (6.1), (6.2) and (6.3). We get a critical collision rate of about $3 \text{s}^{-1}$ for a vacuum life-time of 10 seconds. It should be pointed out that this critical collision rate has no relation with the criteria for run-away evaporation, because the latter tells us if one can keep the collision rate constant; our criterion determines if there is any gain of phase-space density possible, which is a weaker statement. Thus, our rule of thumb is that the initial collision rate multiplied by the initial phase-space density should be larger than the critical collision rate calculated for a given vacuum lifetime. It is easy to see that the initial conditions of Barrett et al. [10], and of Cennini et al. [14] fulfill this criteria. Unfortunately, this criteria does not say anything about the final number of atoms. We face this problem at evaporative cooling in the crossed beam ODT, where one beam is vertical. Despite the fact that the collision rate after 5 s is still large enough for efficient evaporation, further evaporation is not useful due to the unpractically small number of atoms. Although the experiments were performed using $^{87}\text{Rb}$, our conclusions are also valid for others species.

The best results were obtained with a crossed beam ODT, with beams crossed in the vertical plane. This geometry of an ODT is closely resemble the geometry used by Takasu et al. [21]. It is a first use an ODT with beams crossed in the vertical plane for evaporative cooling of $^{87}\text{Rb}$. The maximum phase-space density was achieved by reducing the horizontal beam exponentially with a $1/e$ time $\tau$ of 1 s and slowly reducing the vertical beam exponentially with a $1/e$ time $\tau$ of 5 s. These evaporative ramps in a combination with our geometry of ODT allows us to
maintain a relatively high collision rate until the end of the evaporative cooling. This scheme of evaporative cooling is potentially attractive for heavy species with a weak magnetic moment.

We reached a phase-space density just above the BEC phase transition. For the final trap frequencies we estimate a critical temperature of 340 nK. We observe an anisotropic expansion of an atomic cloud. We extract the temperature (325 nK) below the critical temperature on basis of the observed aspect ratio. The number of condensed atoms is about 140, which is roughly 10% from a total number of the atoms in the trap. The small number of atoms leads to a signal to noise ratio that does not allow us to observe a bimodal distribution directly. However we develop the numerical model which allows us to estimate a small condensed fraction on base of a fitted Gaussian widths of an expanded cloud. This is a first all optical BEC of $^{87}\text{Rb}$, reached by evaporative cooling into a near infrared stationary ODT.

A compressible ODT is described. We report on an efficient evaporative cooling during the compassion. After compression an evaporative cooling form this ODT suffers from a decompression, as in a stationary ODT.

Finally we briefly discuss possible improvements. The use of a compressible ODT is a realistic way to solve the problem of decompression during evaporation, since a drop of the trap frequencies due to decrease of the optical power can be compensated by a decrease of the trap volume. A crossed beam compressible ODT is a straightforward improvement of the present setup, this will allow to reach smaller final volume of the ODT. More efficient loading into an ODT could be achieved if the atoms are evaporatively pre-cooled in a magnetic trap. This will decrease the size of the atomic cloud and improve the overlap between the cloud and the ODT.
References


REFERENCES


Summary

This thesis can be divided in two parts. The first part describes the study of scattering properties of an atom near a dielectric surface. The second part is devoted to trapping and cooling of atoms in a far off-resonance optical dipole trap. We start this summary with the first one.

As is known, an electronically excited atom decays to the ground state by spontaneous emission. The characteristic spontaneous emission rate at which this occurs is not an intrinsic property of the atom but has strict connection with the environment. Basically, the natural spontaneous emission rate is a consequence of vacuum field fluctuations of an electric field. The non-zero field fluctuations are always present because of Heisenberg’s uncertainty principle $\Delta E \times \Delta t \geq \hbar$, where energy can be treated as energy of an electro-magnetic field. Therefore any atom has a finite natural spontaneous emission rate. The electromagnetic boundary conditions modify the intensity of the vacuum field fluctuations. In the case of a simple dielectric surface it leads to decaying contributions of the field fluctuations from the dielectric material to the vacuum, and thus to an increase of the spectral intensity of field fluctuations in vicinity of an dielectric surface. Therefore the spontaneous emission rate of an atom increases in the vicinity of a dielectric surface. The spontaneous emission rate presents itself as an absorption linewidth, which can be measured in an experiment.

![Figure 6.21: A glass prism and a probe beam. A glass prism is located in a vacuum chamber. A resonant probe beam enters the prism and totally reflects from the glass-vacuum interface. The bright spot on the top of the prism is visible in the position of the reflection of the probe beam. The probe beam is visible due to scattering by Rb vapor.](image)

A cloud of cold atoms is a good probe for studying such quantum electrodynamics (QED) phenomena because of its low temperature and low density. We prepare a cloud of cold $^{87}$Rb atoms in a conventional magneto-optical trap (MOT) and cool the atoms to a temperature of $\sim 9 \, \mu K$ using optical molasses. We release the cloud of
cold atoms from the MOT and then it drops on the dielectric surface due to gravity. We measure the absorption of the evanescent waves (EW) light. An EW appears when a beam undergoes total internal reflection at a surface. The optical field at the vacuum side decays exponentially with the distance to the surface. Atoms can absorb light from the EW if their distance to the surface is on the order of the decay length. Therefore this method is selectively sensitive to atoms very close to the surface, at distances in order of the wavelength of the atomic transition (780 nm).

The falling atoms are located in vacuum, and thus the atoms do not interact with any medium (only with the dielectric surface). We experimentally observe the distance-dependent absorption linewidth of cold $^{87}$Rb atoms near a dielectric surface. We observe a broadening of absorption linewidth in order of 20% compared to free space value.

In the second part of this thesis we study trapping and cooling of cold atoms in a far off-resonance optical dipole trap. The focus of a red detuned laser beam can be used as a conservative trap for neutral atoms, so called optical dipole trap (ODT). Nowadays lasers, which can provide enough optical power for substantially deep ODT, are commercially available. Optical power in the order of a few Watts is typically needed. The use 2 $\mu$m light to produce an ODT for cold atoms is completely new. We study the loading of ODTs directly from dark molasses. Dark molasses is a modification of a conventional MOT, which allows a temporary suppression of the photon rescattering of the light.

![Figure 6.22: A high power laser beam and a blade.](image)

The idea of evaporative cooling is based on the preferential removal of atoms with an energy higher than the average energy per atom (i.e. "hot atoms"). A thermal energy distribution establishes with a lower temperature than before the removal of the hot atoms. The evaporative cooling leads to a decrease of temperature, this is also a standard technique to increase the phase space density. We study an
evaporative cooling from an ODT. In a magnetic trap forced evaporative cooling is usually performed by using a so-called radio-frequency (RF) knife, which "cuts" the thermal distribution in a certain point. The RF-knife technique does not work for an ODT, since an ODT traps atoms in all magnetic states. Evaporative cooling in an ODT is performed by a decrease of the optical power of the trapping beam and thus of the trap depth. Therefore evaporative cooling from an ODT suffers from decompression due to a decrease of the trap depth, and therefore of the trap frequencies. In contrast, in magnetic traps the density actually increases during evaporative cooling. Although we reach higher initial values for the density, the collision rate, and the phase space density after loading than achieved with typical magnetic traps, evaporative cooling quickly becomes inefficient (i.e. does not lead to an increase of the phase-space density) due to a drop of the collision rate. There are no known criteria defining an efficient evaporation strategy from an ODT.

The problem of unwanted decompression can be partially solved using a crossed-beam ODT by making power reduction of the horizontal beam faster than of the vertical beam during evaporative cooling. We achieved a phase-space density just above critical after evaporative cooling from this trap. Another way to solve this problem is a compressible ODT, i.e. an ODT where the waist of the trap beam (beams) is reducing during the evaporation. Thus the decrease of the trap frequencies due to reduction of the beam power is compensated by a rise of the trap frequencies due to a decrease of the trap volume. We demonstrated trial example of such trap, where we demonstrate a compression and evaporative cooling from such trap.
De titel van dit proefschrift is "Koude atomen: veranderde stralingseigenschappen en verdampingskoeling vanuit optische vallen". Deze titel doet al een tweedeling van het onderzoek vermoeden, die er inderdaad is. De koude atomen zijn hierbij eerst middel en later doel. Het feit dat de levensduur van veel aangeslagen toestanden van atomen op encyclopedische wijze is vastgelegd, doet vermoeden dat het een intrinsieke eigenschap van het atoom betreft. Dat is niet het geval. Spontaan verval wordt, ondanks de naam, ge¨ınduceerd door vacu¨umfluctuaties van het licht. Hoewel de quantummechanische vacu¨umfluctuaties niet volgen uit de klassieke Maxwell-vergelijkingen, voldoen ze wel aan de randvoorwaarden, die uit de Maxwell-vergelijkingen voortkomen. Het gevolg hiervan is dat stralingsdichtheid in de buurt van een overgang van een materiaal naar een ander materiaal, anders is dan in het homogene geval. In dit proefschrift is de overgang vacu¨um-glas gebruikt om locale veranderingen in de vacu¨umfluctuaties te cre¨eren. Ruimtelijk strekken deze veranderingen zich uit over afstanden van de orde van grootte van de golflengte van het licht. Deze bijna letterlijk microscopisch kleine afstanden maken dat alleen atomen in de buurt van de glas-vacu¨um overgang een veranderd spectrum hebben. En dat, ondanks de naam, veroorzaakt spontaan verval van de atomen in de publieke ruimte van het licht. De tweede eis is dat selectief gebruik kan worden gemaakt van evanescent licht.

Het tweede deel van het proefschrift gaat over de mogelijkheid om atomen gevan- gen te houden en te koelen in een lichtbundel. De overgrote meerderheid van exper- imenteren met koude atomen begint in een zogenaamde magneto-optische val (MOT). Een MOT is een onovertroffen startpunt voor een experiment, maar voor een vol- gende stap is vaak de heftige interactie tussen licht en atomen een probleem. Daarom worden de atomen overgeladen in een andere val die deze beperking niet kent. De bekendste aanpak is omladen naar een magnetische val. Een alternatief is omladen in een optische val, waarvan de kleur zo anders is dan van het licht waarmee de atomen wisselwerken, dat de atomen wel kunnen worden gevangen, maar niet worden opge- warmd. Historisch is deze aanpak veel geprobeerd met licht met een golflengte van 1 µm; vaak met tegenvallende resultaten. Het optisch gevangen houden en vooral het
comprimeren van een wolk atomen werd pas een succes toen infra-rood licht met een golflengte van 10 µm gebruikt werd. Tot voor kort was onbekend waarom het succes met 10 µm licht zo groot was en met 1 micron licht zo klein. Verklaringen werden verdeeld in twee categorieën: het essentiële verschil is de spectroscopie (de energieniveaus van atomen worden door 10 µm licht de ene kant op geduwd en door 1 µm licht de andere kant). De tweede categorie verklaringen is gebaseerd op een verschil in geometrie. Door de kortere golflengte van het licht zijn 1 µm traps veel kleiner dan 10 µm traps. In de praktijk is er een scherp onderscheid tussen succesvolle vallen en tegenvallende vallen: het onderscheid wordt bepaald door de aan- of afwezigheid van de mogelijkheid om een Bose-Einsteincondensaat te maken in een val.

In het kader van dit proefschrift is voor het eerst gewerkt met optische vallen op basis van licht met een golflengte van 2 µm. Dat is een interessante golflengte: als de spectroscopische verklaringen voor het succes van 10 µm vallen waar zijn, dan is de 2 µm val waarschijnlijk ook heel succesvol, wat de licht-geïnduceerde spectroscopische verschuivingen klappen van teken om bij 1.6 µm. In dit opzicht is 2 µm licht ‘rood’. Veel materialen die in het zichtbare gebied doorzichtig zijn, zijn dat bij 2 µm ook. Maar bij 10 µm zijn totaal andere technieken nodig. Als 2 µm met zo goed zou werken als 10, dan is het om praktische redenen een stap vooruit. Als de op geometrie gebaseerde verklaringen juist zijn, dan is 2 µm een keus die maar een beetje minder slecht is dan 1 µm. Experimenten in dit proefschrift hebben laten zien dat de geometrische verklaring de goede is: het succes bij 10 µm is er niet bij 2 µm. De beste resultaten zijn niet behaald met 2 µm licht maar met een combinatie van 1 en 2 µm licht, waarbij ook de zwaartekracht een gunstige rol bij het koelen speelde. Het feit dat 2 µm een factor 2 langgolviger is dan 1 µm heeft ervoor gezorgd dat het net aan mogelijk bleek om Bose-Einsteincondensatie te bereiken. Met 1 µm licht is dat niet mogelijk, althans niet met vallen die een vaste vorm hebben. Collega’s zijn wel in staat geweest om gedurende het koelen de vorm van de val zo te veranderen dat ook bij 1 µm condensatie mogelijk is.

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