

5. Physikalisches Institut – Universität Stuttgart

Masterarbeit

Dipolar and universal scattering of ultracold dysprosium atoms – Demagnetization cooling and broad Feshbach resonances

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3. November 2015

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Stuttgart, den 3. November 2015

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Zusammenfassung

Inhalt dieser Arbeit ist die Verwendung dipolarer Streuung zur Entmagnetisierungskühlung von ultrakalten Dysprosiumatomen, sowie die Erforschung von dipolarer und universeller Streuung anhand zweier breiter Feshbachresonanzen in ¹⁶⁴Dy. In Gasen ultrakalter Dysprosiumatome kann dipolare Streuung aufgrund des hohen magnetischen Moments beobachtet werden. Dieses hohe magnetische Moment hat zur Folge, dass die langreichweitige magnetische Dipol-Dipol-Wechselwirkung stärker ist als die von Alkaliatomen bekannte kurzreichweitige Kontaktwechselwirkung.

Bei der Entmagnetisierungskühlung handelt es sich um eine Kühlmethode, die, im Gegensatz zur herkömmlichen Verdampfungskühlung, theoretisch verlustfrei ist. Hierbei werden inelastische Stöße, die mit dipolaren Relaxationen verbunden sind, verwendet um einen Teil der Bewegungsenergie der Atome in potentielle Energie eines in einem externen Magnetfeld energetisch höheren Spinzustandes umzuwandeln. Durch optisches Pumpen zurück in den vollständig spinpolarisierten und dunklen Grundzustand, kann die entsprechende Energie mit dem Lichtfeld der Atomwolke entzogen werden [1–3]. In Chrom konnte so bereits eine Kühleffizienz von $\chi_{Cr} > 17$ erreicht werden [4]. Im Rahmen dieser Arbeit wurde Entmagnetisierungskühlung von Dysprosiumatomen mit einer Effizienz von $\chi_{demag} = 7.95 \pm 1.58$ demonstriert. Diese Effizienz übersteigt die für Verdampfungskühlung übliche Effizienz von $\chi_{evap} \approx 4$, ist jedoch von der selben Größenordnung wie die Effizienz des von uns an dieser Stelle zusätzlich verwendeten Nachkühlens ($\chi_{626} \approx 6.5$ [5]).

In Dysprosium führt die langreichweitige anisotrope magnetische Dipol-Dipol-Wechselwirkung und die ebenfalls anisotrope Kontaktwechselwirkung zu einem chaotischen Hintergrund schmaler Feshbachresonanzen [6]. Trotzdem konnten wir zwei gebundene Zustände, die mit breiten Feshbachresonanzen einhergehen, beobachten. Diese Zustände bestehen über die Kreuzung mit vielen schwachen Zuständen hinweg und sind vom chaotischen Hintergrund entkoppelt. Wir konnten mithilfe von Atomverlustspektroskopie, Magnetfeldmodulationsspektroskopie und der Analyse unserer Daten mit gekoppelten Kanal-Berechnungen (engl. coupled channel calculations) und den universellen Formeln die Streulänge als Funktion des Magnetfelds und eine Hintergrundstreulänge gewinnen. Wir fanden Anhaltspunkte, dass die Atomverluste in der Nähe der Resonanzen mit der universellen Verlustdynamik vereinbar sind. Durch den Vergleich unserer Daten mit der universellen Verlustdynamik konnten wir so die Fallentiefe abschätzen. Außerdem konnten wir anhand der Temperaturabhängigkeit der Dreikörperverluste im unitären Regime den Elastizitätsparameter ermitteln.

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CHAPTER 1

Introduction

In 1982 Feynman proposed the simulation of physics by universal computers [7]. Within the last years, it could be demonstrated that many-body quantum systems, which are usually used to describe condensed matter physics, may be simulated by ultracold atomic gases [8, 9]. Compared to condensed matter, ultracold atomic samples offer an exceptional possibility to control with high precision both, the external parameters, such as the trapping potential, and the internal degrees of freedom, like the contact interaction [10].

Laser cooling and trapping laid the groundwork for the study of quantum behaviour of dilute atomic vapours at very low temperatures and was awarded in 1997 with the Nobel Prize in Physics [11–13]. By combining laser cooling and evaporative cooling methods, quantum degeneracy could be reached in 1995 [14–16]. Since then quantum gases of alkali atoms with isotropic short-range contact interaction have been studied extensively [8].

Recently, laser cooling and trapping of more exotic atomic species such as thulium [17] and holmium [18] have been reported. In erbium and dysprosium even Bose-Einstein condensation [19, 20] and Fermi degeneracy [21, 22] could be observed.

The complex electronic structure of those species gives rise to enriched interatomic interactions. First, the magnetic dipole-dipole interaction adds an anisotropic long-range interaction and furthermore, the short-range interaction becomes anisotropic. Due to the multiplicity of interactions, even molecular potentials are hard to compute and calculations generally attain a very high intricacy.

The wealth of interactions combined with the restrictions for calculations opens the door to the observation of a bunch of interesting phenomena. For instance, in the presence of anisotropic interactions the Fermi surface is no longer spherical. This deformation of the Fermi surface due to the magnetic dipole-dipole interaction could be observed in erbium [23]. Moreover, identical fermions do not collide at very low temperatures. This hampers evaporative cooling in Fermi gases. However, Fermi degeneracy could be reached in a dipolar gas due to universal dipolar scattering [21] Besides, the observation of discrete, stable, and self-ordered droplets in a quantum ferrofluid in close analogy to the Rosensweig instability [24], could not be predicted by mean field calculations [25, 26] and poses a new riddle for theoreticians. Conceivably, stabilizing quantum fluctuations leading to beyond mean-field effects [27] could be the solution to the puzzle. In addition, erbium and dysprosium show an extremely dense spectrum of narrow Feshbach resonances, which is characterized by a chaotic distribution [28, 29].

This raises the question whether precise control of interactions is also possible in a lanthanide such as dysprosium and the verification of simple universal expressions becomes essential.

About this thesis

In chapter 2 we will have an insight into the intriguing properties of dysprosium.

We use dipolar scattering for demagnetization cooling and address this topic in chapter 3. Here, we will at first concentrate on the magnetic dipole-dipole interaction and subsequent dipolar relaxations which are the prerequisit for demagnetization cooling. In the following, we will direct the attention to actual demagnetization cooling excluding heating mechanisms due to optical pumping in section 3.3 and in the following section 3.4, including the resulting experimental restrictions. In section 3.5, we will discuss the experimental realization including the experimental setup and the analysis of demagnetization cooling of ultracold dysprosium atoms.

In chapter 4, we show that two broad magnetic Feshbach resonances persist across multiple narrow resonances. At first, we will concentrate on the common description of magnetic Feshbach resonances. Second, we will address universal loss dynamics and present rate equations for atom number and temperature. In section 4.3, we present the characterization of the two broad Feshbach resonances by atom loss spectroscopy and magnetic field modulation spectroscopy. Subsequently, we see that the universal expressions are in good agreement with corresponding coupled-channel calculations and the data. In the following, we use this knowledge in order to analyse the two broad resonances in terms of universal loss dynamics in section 4.4.

CHAPTER 2

Properties of dysprosium

Dysprosium is a lanthanide with atomic number 66. Naturally occurring dysprosium is composed of 7 stable isotopes. The by nature most abundant stable bosonic isotopes are ¹⁶⁴Dy (28%) and ¹⁶²Dy (26%) and the naturally most abundant fermionic isotopes are ¹⁶³Dy (25%) and ¹⁶¹Dy (19%) [30]. Dysprosium in its groundstate ⁵I₈ has the electronic configuration [Xe] 4f¹⁰ 6s², which means that four open *f*-states lie beneath a filled 6*s*-shell. This gives rise to an angular momentum L = 6 and a spin S = 2 and thus a total angular momentum J = 8.

In figure 2.1 a part of the energy level scheme of dysprosium is shown. It contains the 684 nm optical pumping transition, which was used in this thesis for continuous demagnetization cooling. Further, the transitions used for Zeeman slowing and transverse cooling of the atomic dysprosium beam (421 nm) [35] and atom trapping in a magneto optical trap (626 nm) are displayed. The 421 nm transition has a linewidth of $\Gamma_{421}/2\pi = 32.2$ MHz [36], which leads to a Doppler temperature of $T_{D,421} = \hbar\Gamma/2k_B = 773 \,\mu\text{K}$ [32]. Here, \hbar is the reduced Planck constant and k_B the Boltzmann constant. The 626 nm transition is closed and due to its narrow linewidth of $\Gamma_{626}/2\pi = 136 \,\text{kHz}$, the Doppler temperature is $T_{D,626} = 3.3 \,\mu\text{K}$ [32]. The 684 nm optical pumping transition is a transition from the ground state [Xe] $4f^{10} \,\text{6s}^2 \,{}^5\text{I}_8$ to a $4f^9(^6\text{H}^0)5d \,\text{6s}^2 \,{}^5\text{I}_8$ excited state, where one of the ten 4f electrons is excited to a 5d electron state. The linewidth of the 684 nm transition $\Gamma_{684}/2\pi = 95 \,\text{kHz}$ leads to a Doppler temperature of $T_{D,684} = 2.3 \,\mu\text{K}$.

In contrast to the bosonic isotopes, the fermionic isotopes additionally have a finite nuclear spin I = 5/2. Therefore, they exhibit a hyperfine splitting with quantum numbers ranging from F = 11/2 to F = 21/2 due to the coupling to the total angular momentum J [37].

Besides, the total angular momentum J = 8 of dysprosium in its groundstate together with the Landé-factor of $g_J = 1.24125$ lead to a magnetic moment of $\mu = 9.93 \,\mu_B$. Therefore, dysprosium has (besides Terbium) the highest magnetic moment of all stable elements naturally abundant on earth. From this high magnetic moment arises one key property of dysprosium concerning the interatomic interaction in a dilute gas. In comparison to alkali atoms, which have only one unpaired electron, the magnetic dipole-dipole interaction is 100 times stronger.



Figure 2.1: Part of the level scheme of Dysprosium [31]. The 421 nm transition is used for Zeeman slowing and the 626 nm transition is used for magneto-optical trapping [32], whereas the 684 nm transition is used for optical pumping. The 953 nm transition could operate like a decay channel for the optical pumping transition, but is very narrow [33, 34]. The gray line indicates the wavelength of 1064 nm which we use for optical trapping.

Apart from this long-range dipole-dipole interaction, the high total atomic angular momentum J = 8 leads to a multitude of non-degenerate molecular potentials and a correspondingly large manifold of collision channels. Furthermore, the complex electronic structure leads to an anisotropic van der Waals interaction. Due to the anisotropy of the molecular potentials, collision channels with large relative angular momentum l are coupled even for ultracold collisions with an initial *s*-wave channel (l = 0) [29]. The corresponding spectrum of Feshbach resonances is presented in chapter 4 and [6].

CHAPTER 3

Demagnetization cooling

In this chapter, I will present the theoretical principle of and a first experimental approach to demagnetization cooling of ultracold dysprosium atoms. At first, I will give a short introduction on the magnetic dipole-dipole interaction, which is essential for demagnetization cooling, as it is the prerequisite for dipolar relaxation collisions. In the following, I will explain the most important theoretical aspects of dipolar relaxation and optical pumping to give a first understanding of the nature of demagnetization cooling. Subsequently, I will briefly review the derivation of dipolar relaxation cross-sections [38, 39] and the resulting relaxation rate coefficients. Further, I will focus on the cooling rate. Next, I will explain optical pumping, which includes experimental restrictions. In order to understand the nature of light-assisted losses, we will consider molecular potentials and extend the typical isolated atom optical pumping picture to atom-atom interactions. Further, I will explain a mechanism to suppress light assisted collisions. In the following, I will present our experimental realization. Therefore, I will briefly explain how we prepare the atomic cloud and will focus on the laser system we use for optical pumping. Finally, I will analyse the performance of demagnetization cooling and show how we optimized it in this first dysprosium prove of principle experiment.

3.1 Magnetic dipole-dipole interaction

In chapter 2, we have seen that dysprosium has a very high magnetic moment. Therefore, the magnetic dipole-dipole interaction is important in scattering processes of dysprosium atoms (see subsec. 3.4.1).

The interaction potential U_{dd} of two magnetic dipoles with magnetic moments μ_1 and μ_2 is given by

$$U_{\rm dd}(\boldsymbol{r}) = \frac{\mu_0}{4\pi} \frac{\boldsymbol{\mu}_1 \cdot \boldsymbol{\mu}_2 - 3\left(\boldsymbol{\mu}_1 \cdot \hat{\boldsymbol{r}}\right)(\boldsymbol{\mu}_2 \cdot \hat{\boldsymbol{r}})}{r^3}, \qquad (3.1)$$

where $\mu_0 = 4 \pi 10^{-7} \text{ N/A}^2$ is the vacuum permeability, \boldsymbol{r} is the interatomic separation vector, r is the distance, and $\hat{\boldsymbol{r}} = \frac{\boldsymbol{r}}{r}$ is the normalized separation vector. When the two dipoles have the same strength $|\boldsymbol{\mu}_1| = |\boldsymbol{\mu}_2| = \mu$ and are aligned parallel to each other $\boldsymbol{\mu}_1 || \boldsymbol{\mu}_2$ (see fig. 3.1) in an external magnetic field, the interaction potential simplifies as follows

$$U_{\rm dd}(\mathbf{r}) = \frac{\mu_0 \,\mu^2}{4\pi} \frac{1 - 3\,\cos\theta}{r^3} \,. \tag{3.2}$$

Here $\theta = \measuredangle(\mu_i, r)$ is the angle between the interatomic separation vector and the orientation of the two dipoles (see fig. 3.1). The dependence on this relative angle θ



Figure 3.1: a) Sketch of two magnetic dipoles, which are polarized by a magnetic field B and separated by a distance r. The angle between the interatomic separation vector and the orientation of the polarized dipoles is called θ . b) Dipole-dipole interaction potential. Yellow represent a positive sign and repulsive interaction and blue a negative sign and attractive interaction. The anisotropic interaction potential depends on the relative angle θ . In head-to-tail configuration it is attractive and if the atoms are oriented parallel to each other it is repulsive. Besides, the interaction is proportional to r^{-3} .

shows that the potential is anisotropic. If the two dipoles are oriented in head-to tail configuration ($\theta = 0^{\circ}$), they attract each other and the interaction potential is negative. If they are oriented parallel to each other ($\theta = 90^{\circ}$), they repel each other and the interaction potential is positive. Moreover, the interaction potential is proportional to r^{-3} . Potentials which scale with r^{-m} , $m \leq 3$ cannot be approxiamted by a contact-like pseudo interaction potential (see subsec. 3.4.1) [40]. Atoms colliding in such a potential are no longer reflected by the centrifugal barrier and thus all partial waves $l \geq 0$ contribute even at zero temperature [41]. The dipole-dipole interaction potential is thus a long-range potential.

Collisions due to dipolar interactions are characterized by the dipolar length [40]

$$D \equiv \frac{\mu_0 \,\mu^2 \,m}{8 \,\pi \,\hbar^2} \stackrel{\text{Dy}}{=} 196 \,a_0 \,, \tag{3.3}$$

which in turn characterizes the magnetic dipole-dipole interaction strength.

3.2 Dipolar relaxations

Despite the fact that ultracold gases are very dilute, collisions are very important in understanding them. The most prominent example is that collisions ensure thermalization in trapped atomic gases, which is essential for evaporative cooling. The high magnetic moment of dysprosium (see chapter 2 and subsec. 3.4.1) gives rise to enriched collisional physics, as for example dipolar relaxation collisions, which are an integral part of demagnetization cooling. Literature on collision physics is abundant [41–44] and goes beyond the scope of this Master's thesis. Here, I will show the connection between dipole-dipole interaction and inelastic relaxation collisions and briefly review the derivation of dipolar relaxation cross-sections [38, 39].

Due to the high magnetic moment of dysprosium, the atoms experience magnetic dipole-dipole interactions. In eq. 3.1 the classical description of the interaction of two magnetic dipoles is given. The corresponding quantum mechanical description of the magnetic dipole-dipole interaction is obtained, if the classical magnetic moments $\boldsymbol{\mu}_i$ in eq. 3.1 are substituted by the quantum mechanical magnetic dipole moment operators $\hat{\boldsymbol{\mu}}_i = g_J \, \mu_{\rm B} \, \boldsymbol{J}_i$

$$\hat{U}_{\rm dd}(\mathbf{r}) = \mu_0 \left(g_J \,\mu_{\rm B}\right)^2 \frac{(\boldsymbol{J}_1 \cdot \boldsymbol{J}_2) - 3 \,(\boldsymbol{J}_1 \cdot \hat{\boldsymbol{r}})(\boldsymbol{J}_2 \cdot \hat{\boldsymbol{r}})}{4 \,\pi \,r^3} \,, \tag{3.4}$$

where, g_J is the Landé-factor, μ_B the Bohr magneton and $J_i = (J_{ix}, J_{iy}, J_{iz})$ is the operator of total angular momentum of the two magnetic dipoles which consists of the spin matrices J_{ix} , J_{iy} and J_{iz} . By defining normalized directions and introducing creation and annihilation operators for orbital and total angular momentum as well as their projection on the magnetic field axis z,

$$\hat{r}_{+} = \frac{x + iy}{r}, \quad \hat{r}_{-} = \frac{x - iy}{r}, \quad \hat{z} = \frac{z}{r},$$
(3.5)

$$J_{+} = J_{x} + iJ_{y}, \quad J_{-} = J_{x} - iJ_{y}, \quad J_{z}, \qquad (3.6)$$

the dipole-dipole interaction (3.4) can be expanded to

$$\hat{U}_{dd} = \frac{\mu_0 \left(g_J \,\mu_B\right)^2}{4 \,\pi \,r^3} \left[J_{1z} \cdot J_{2z} + \frac{1}{2} \left(J_{1+} \cdot J_{2-} + J_{1-} \cdot J_{2+}\right) - \frac{3}{4} \left(2 \,\hat{z} \,J_{1z} + \hat{r}_- \,J_{1+} + \hat{r}_+ \,J_{1-}\right) \left(2 \,\hat{z} \,J_{2z} + \hat{r}_- \,J_{2+} + \hat{r}_+ \,J_{2-}\right)\right].$$
(3.7)

The first term is an elastic scattering term, the second term describes spin-exchange collisions and the third term involves the conversion of orbital angular momentum into spin angular momentum and vice versa, so called dipolar relaxation collisions.

In first-order Born approximation [45] the total cross-section for two identical colliding

bosonic (fermionic) atoms in the same internal state can be expressed as [38, 39, 46]

$$\sigma = \left(\frac{m}{4\pi\hbar^2}\right)^2 \frac{1}{k_{\rm i} k_{\rm f}} \left[\int |\tilde{U}_{\rm dd}(\boldsymbol{k} - \boldsymbol{k'})|^2 \,\delta(|\boldsymbol{k'}| - k_{\rm f}) \mathrm{d}\boldsymbol{k'} \\ \pm \int \mathrm{Re}\left(\tilde{U}_{\rm dd}(\boldsymbol{k} - \boldsymbol{k'}) \,\tilde{U}_{\rm dd}^*(\boldsymbol{k} + \boldsymbol{k'})\right) \delta(|\boldsymbol{k'}| - k_{\rm f}) \mathrm{d}\boldsymbol{k'} \right].$$
(3.8)

Here, $\tilde{U}_{\rm dd}(\boldsymbol{q}) = \langle M'_1, M'_2 | \mathcal{F}(\hat{U}_{\rm dd}) | M_1, M_2 \rangle$ is the Fourier transform of the dipole-dipole interaction (3.4) that has already been contracted between the initial $|M_1, M_2\rangle$ and final internal states $|M'_1, M'_2\rangle^{-1}$, \boldsymbol{k} and $\boldsymbol{k'}$ are the initial and final wavevectors and k_i and k_f are the absolute values of the initial and final wave vectors.

The cross-section σ (3.8) averaged over all possible orientations of **k** reads in the case of elastic relaxations [38, 39, 46]

$$\sigma_{0} = \frac{\pi}{45} \left(\frac{\mu_{0}(g_{J}\mu_{\rm B})^{2}m}{4\pi\hbar^{2}} \right)^{2} \left[C_{S}^{2} \cdot J_{0}^{S}(M_{1}, M_{2}, M_{1}', M_{2}') \left(1 + h(1)\right) + C_{A}^{2} \cdot J_{0}^{A}(M_{1}, M_{2}, M_{1}', M_{2}') \left(1 - h(1)\right) \right]$$
(3.9)

and in the case of inelastic relaxations with $\Delta M = \sum_i M'_i - M_i$ spin-flips

$$\sigma_{\Delta M} = \frac{2\pi}{15} \left(\frac{\mu_0 (g_J \mu_B)^2 m}{4\pi \hbar^2} \right)^2 \left[C_S^2 \cdot J_{\Delta M}^S (M_1, M_2, M_1', M_2') \left(1 + h(k_f/k_i) \right) + C_A^2 \cdot J_{\Delta M}^A (M_1, M_2, M_1', M_2') \left(1 - h(k_f/k_i) \right) \right] \frac{k_f}{k_i}.$$
(3.10)

Here $C_S = \frac{1}{\sqrt{2-\delta_{M_1,M_2}}}$ and $C_A = 1 - \frac{\delta_{M_1,M_2}}{\sqrt{2}}$ are the normalization coefficients of the symmetric (antisymmetric) part of the total wave function of two bosonic atoms [46], $J_{\Delta M}^{S(A)}(M_1,M_2,M'_1,M'_2)$ describes the transition probabilities between two properly symmetrized (anti-symmetrized) two-particle spin states which change their magnetic

1 If an internal (spin) state of atom i = 1,2 is denoted by $|M,J\rangle_i$, the symmetric internal states of the atoms are defined as follows:

$$\begin{split} |M_1,M_2\rangle_{\rm S} &= \frac{1}{\sqrt{2}}(|M_1,J\rangle_1 \, |M_2,J\rangle_2 + |M_2,J\rangle_1 \, |M_1,J\rangle_2) \\ \text{The antisymmetric internal states are defined in close analogy:} \\ |M_1,M_2\rangle_{\rm A} &= \frac{1}{\sqrt{2}}(|M_1,J\rangle_1 \, |M_2,J\rangle_2 - |M_2,J\rangle_1 \, |M_1,J\rangle_2) \\ \text{Here } M_i \text{ is the projection of } J \text{ on the quantization axis.} \\ \text{Having this in mind, the initial total wave function of two bosonic atoms including not only the internal but also the spatial part reads} \\ |\psi\rangle_{\rm in} &= \sqrt{\frac{2}{V}}(C_{\rm S} \, \cos\left(\mathbf{kr}\right) |M_1;M_2\rangle_{\rm S} + C_{\rm A} \, \sin\left(\mathbf{kr}\right) |M_1;M_2\rangle_{\rm A}) \\ \text{Here } C_S &= \frac{1}{\sqrt{2-\delta_{M_1,M_2}}} \text{ and } C_A = 1 - \frac{\delta_{M_1,M_2}}{\sqrt{2}} \text{ are the normalization coefficients.} \end{split}$$

quantum numbers from M_1 and M_2 to M'_1 and M'_2 ¹, $\Delta M = \sum_i M'_i - M_i$ is the number of spin-flips in one collision, and h(x) is defined as

$$h(x) = \begin{cases} -\frac{1}{2} & \text{if } x = 1\\ -\frac{1}{2} - \frac{3}{8} \frac{(1-x^2)^2}{x(1+x^2)} \ln \frac{(1-x)^2}{(1+x)^2} & \text{if } x > 1 \,. \end{cases}$$
(3.11)

To the alert eye it will become apparent that the scattering cross-section σ of dipolar relaxations depends on the magnetic quantum numbers before M_1 , M_2 and after the collision M'_1 , M'_2 and thus a lot of different collision channels arise. However, in experiments we use optical pumping to ensure that the atoms are always fully spin-polarized $M_1 = M_2 = -J$ and thus the population of other states becomes negligible. This reduces the number of collision channels dramatically. The dipole-dipole interaction can mediate direct spin-flips with maximal $\Delta M = 2$. Here we will therefore first focus on spin-flips from the fully spin polarized state to the $|M'_1 = -J + 1, M'_2 = -J \rangle = |M'_1 = -J, M'_2 = -J + 1 \rangle$ and $|-J + 1, -J + 1 \rangle$ states. For a fully polarized initial state with magnetic quantum numbers $M_1 = M_2 = \pm J$ the elastic scattering cross-section reduces to

$$\sigma_0 = \frac{16\pi}{45} J^4 \left(\frac{\mu_0 (g_J \mu_B)^2 m}{4\pi\hbar^2}\right)^2 (1+h(1)) = \frac{64\pi}{45} D^2 (1+h(1))$$
(3.12)

and the inelastic scattering cross-section to states with $\Delta M = 1$ becomes

$$\sigma_1 = \frac{8\pi}{15} J^3 \left(\frac{\mu_0 (g_J \mu_{\rm B})^2 m}{4\pi\hbar^2}\right)^2 \left(1 + h(k_{\rm f}/k_{\rm i})\right) \frac{k_{\rm f}}{k_{\rm i}}$$
(3.13)

and to states with $\Delta M = 2$ respectively

$$\sigma_2 = \frac{8\pi}{15} J^2 \left(\frac{\mu_0 (g_J \mu_B)^2 m}{4\pi\hbar^2}\right)^2 \left(1 + h(k_f/k_i)\right) \frac{k_f}{k_i}$$
(3.14)

If the sample is fully polarized in the $M_1 = M_2 = -J$ state, the above mentioned inelastic scattering cross-sections will lead to a cooling of the motional degree of freedom (see fig. 3.2). Though not that probable, because of constant repumping of the atoms back to the fully spin polarized state, also a heating due to dipolar relaxations from energetically higher, e.g $M_1, M_2 = -J + 1$ states, back to the fully spin polarized state is possible. In the case of a single spin-flip ($\Delta M = -1$) from $|-J, -J + 1\rangle$ to

 $\begin{array}{l} \hline \\ \hline 1 & \text{The transition probabilities } J^{S(A)}_{\Delta M} \text{ for elastic } (\Delta M=0), \text{ single spin-flip } (\Delta M=\pm 1) \text{ and double spin-flip } (\Delta M=\pm 2) \text{ transitions read } [38, 46]: \\ J^{S(A)}_0(M_1,M_2,M'_1,M'_2) = |_{S(A)} \langle M'_1,M'_2| \left(4 J_{1z} J_{2z} + J_{1+} J_{2-} + J_{1-} J_{2+}\right) |M_1,M_2\rangle_{S(A)} |^2 \\ J^{S(A)}_{\pm 1}(M_1,M_2,M'_1,M'_2) = |_{S(A)} \langle M'_1,M'_2| \left(J_{1,\pm} J_{2z} + J_{1z} J_{2\pm}\right) |M_1,M_2\rangle_{S(A)} |^2 \\ J^{S(A)}_{\pm 2}(M_1,M_2,M'_1,M'_2) = |_{S(A)} \langle M'_1,M'_2| \left(J_{1,\pm} J_{2z} + J_{1z} J_{2\pm}\right) |M_1,M_2\rangle_{S(A)} |^2 \end{array}$



Figure 3.2: Spin-flips from the fully spin-polarized state to the neighbouring one lead to a cooling of the motional degree of freedom, as in inelastic dipolar relaxation collisions the Zeeman energy $\Delta E_{\rm Z}$ per spin-flip is transferred from the kinetic energy to the spin energy reservoir. Image taken from [3].

 $|-J, -J\rangle$ the corresponding scattering cross-section is

$$\sigma_{-1} = \frac{4\pi}{15} J^3 \left(\frac{\mu_0 (g_J \mu_B)^2 m}{4\pi \hbar^2} \right)^2 \left(1 + h(k_f/k_i) \right) \frac{k_f}{k_i}$$
(3.15)

It is reduced by a factor $\frac{1}{2}$ compared to σ_1 due to the different normalization factor C_S . This represents the fact that the initial state is not a fully spin-polarized state, which means that there exist for one initial state no longer two relaxation possibilites (here: $|-J, -J\rangle$ to $|-J, -J+1\rangle$ and $|-J+1, -J\rangle$) but only one possibility (here: $|-J, -J+1\rangle$ to $|-J, -J\rangle$ or $|-J+1, -J\rangle$ to $|-J, -J\rangle$).

3.3 Demagnetization cooling and optical pumping

In the previous section 3.2 dipolar relaxations, which couple motional and spin degree of freedom, have been considered. During such relaxations, the sum of the quantum numbers M_1 and M_2 of two colliding atoms is not conserved and the demagnetization of the sample is allowed. If the particles are initially fully spin-polarized in the Zeeman state M = -J, a part of the kinetic energy of the sample is transferred into Zeeman energy during demagnetization. This involves a net cooling effect [3].



Figure 3.3: Illustration of demagnetization cooling [3]. a) Initially, the atoms are fully spin polarized in M = -J at a high magnetic field B. The Zeeman energy separation $\Delta E = g_J \mu_B B$ is much higher than the mean energy of the system $k_B T$. b) When the magnetic field is decreased such that $k_B T \approx g_J \mu_B B$, higher Zeeman sublevels become energetically accessible and dipolar relaxations lead to a demagnetization of the system. c) Optical pumping offers the possibility to polarize the atoms back to M = -J. This cooling cycle may then be repeated several times.

In figure 3.3 demagnetization cooling is illustrated. Initially the particles are fully spin-polarized in the Zeeman state M = -J at high magnetic field B. In this regime, the mean energy $k_{\rm B}T$ is well below the Zeeman energy separation $\Delta E = g_J \mu_{\rm B} B$. Reducing the external magnetic field until $k_{\rm B}T \approx g_J \mu_{\rm B} B$, higher Zeeman-substates become energetically accessible. Hence, dipolar relaxations lead to a demagnetization of the system. During demagnetization a part of the energy of the sample is transferred into Zeeman energy, which involves a heating of the spin energy reservoir and a cooling of the kinetic energy reservoir. The big advantage of cooling an atomic cloud compared to cooling solid state systems by demagnetization cooling, is the possibility to make use of the transparency of an atomic cloud, which allows to cool not only the motional degree of freedom, but also the spin degree of freedom by optical pumping into a dark state. Therefore, the cooling cycle can be repeated several times ¹. In fact, it is possible to use optical pumping to polarize the atoms back to the fully spin-polarized

¹ However, the cooling mechanism is much more efficient in a single circle in solid state systems than in atomic physics. [3, 38, 47]

Zeeman state M = -J, leaving the temperature nearly unchanged. This offers the possibility to constantly recycle the atoms that were thermally excited to higher spin states. Consequently the system does not reach thermal equilibrium in which back and forth relaxations cancel out and hence cooling is possible over a wide temperature range. This cooling scheme has been proposed in 1950 by A. Kastler [1].

In the following subsections 3.3.1 to 3.3.2 the dipolar relaxation rate and the cooling effect of demagnetization are discussed assuming optical pumping which leaves the temperature and the atom number unchanged.

3.3.1 Relaxation rate coefficients

The time-scale on which dipolar relaxation collisions happen is given by the dipolar relaxation rate $\Gamma = \beta n$ with atomic density n. In this subsection we will therefore focus on the relaxation rate coefficient β .

The dipolar relaxation cross-sections which have been discussed in section 3.2 still contain the relation $k_{\rm f}/k_{\rm i}$ of the initial and final wave vectors. If we apply a magnetic field B, the 2J + 1 substates of the atoms are no longer degenerate, but split by the Zeeman energy. The energy difference between two neighboured Zeeman substates is

$$\Delta E = g_J \mu_{\rm B} B \,. \tag{3.16}$$

The energy difference between the final and initial states $E_{\rm f, i}$ is given by $\Delta E_{\rm z} = E_{\rm f} - E_{\rm i} = \Delta E \cdot \Delta M$. The final wavevector then reads

$$k_{\rm f} = \sqrt{k_{\rm i}^2 - m\Delta E_{\rm z}/\hbar^2} \,. \tag{3.17}$$

In order to obtain the relaxation rate coefficients the above mentioned scattering cross-sections have to be averaged thermally.

If $\Delta M > 0$ and the magnetic fields are finite, only atoms with a minimal initial energy corresponding to a velocity $v_{\Delta M, \min} = \sqrt{4\Delta E_z/m} = \sqrt{4\Delta M g_J \mu_{\rm B} B/m}$ have enough energy to reach the energetically higher Zeeman substate and thus can participate in dipolar relaxations. Therefore, the relaxation rate coefficient that describes the transition of atoms from M = -J to M' = -J + 1 reads

$$\beta_{\mathrm{dr}}^{+} = \left\langle \left(\sigma_{1} + 2\,\sigma_{2}\right)v_{\mathrm{rel}}\right\rangle_{\mathrm{thermal}} = \int_{v_{1,\,\mathrm{min}}}^{\infty} \sigma_{1}\,v\,f_{v}\,\mathrm{d}v + 2\int_{v_{2,\,\mathrm{min}}}^{\infty} \sigma_{2}\,v\,f_{v}\,\mathrm{d}v\,,\qquad(3.18)$$

where $v_{\rm rel}$ is the initial relative collision velocity and the Boltzmann distribution

$$f_v = 4\pi \left(\frac{\mu_{\rm red}}{2\pi k_{\rm B}T}\right)^{\frac{3}{2}} v^2 \exp \frac{-\mu_{\rm red}v^2}{2k_{\rm B}T}$$
(3.19)

weights the initial relative collision velocities in the integral with their probability to occur. The initial absolute wavevector k_i reduces with the relative momentum $p_{\rm rel} = \hbar k_{\rm rel} = \mu_{\rm red} v_{\rm rel}$ and the reduced mass $\mu_{\rm red} = \frac{m}{2}$ to $k_{\rm i} = \frac{mv}{2\hbar}$. The factor of 2 in the case of double spin-flips accounts for the double loss of atoms with M = -J compared to a single spin-flip.



Figure 3.4: Dipolar relaxation rate coefficient of a) ¹⁶⁴Dy and b) ⁵²Cr dependent on temperature and for different magnetic field values. B = 400 mG and B = 200 mG are typical magnetic field values used in the experiment. β_{dr}^+ accounts for single and double spin-flips from the fully spin-polarized to the adjacent energetically higher Zeeman-states. β_{dr}^- accounts for single spin-flips to the fully spin-polarized state. The dipolar relaxation rates of Dysprosium are enhanced by approximately one order of magnitude compared to Chromium due to the higher mass and magnetic dipole moment of Dysprosium.

Clearly, the dipolar relaxation rate constant β_{dr} depends on temperature as well as on the magnetic field. In figure 3.4 it is shown a) for Dysprosium and b) for Chromium for different magnetic fields in a temperature range from 0 to 50 µK. The higher mass and magnetic moment of Dysprosium enhance its dipolar relaxation constant by about one order of magnitude.

The dipolar relaxation rate constant β_{dr}^+ which promotes the atoms to energetically higher spin states and thus cools the motional degree of freedom is suppressed when lowering the temperature, since lowering the temperature lowers the velocity of the atoms and hence the frequency of scattering events. The higher the magnetic field, the earlier it is suppressed, as the Zeeman energy increases with the magnetic field and thus less atoms have enough energy to be promoted to M = -J + 1 states. Therefore a reduction of the magnetic field with decreasing temperature seems to be reasonable. However, a reduction of the magnetic field results in a reduction of the Zeeman energy and therefore in a reduction of the cooling effect per scattered atom. Consequently, a theoretical and experimental investigation on optimal magnetic fields for certain temperatures is necessary. A first step in order to investigate it theoretically, is to evaluate the cooling rate, which includes not only the dipolar relaxation rate constant, but also the Zeeman energy, the shape of the trap and the density of atoms in the state M = -J. This will be done in the next subsection 3.3.2.

The reverse process of dipolar relaxations to energetically lower states leads to a heating of the motional degree of freedom and corresponds to the dipolar relaxation rate constant $\beta_{\rm dr}^-$, which is also displayed in figure 3.4 (dashed line). In order to calculate it, only single spin-flips with $\Delta M = -1$ have been taken into account. Furthermore, the integral for thermal averaging takes all velocities from 0 to ∞ into account, since spin-flips to energetically lower states are allowed for all initial kinetic energies.

3.3.2 Cooling rate

Considering the relaxation rate coefficients, the two-body rate equations describing the change of population in the different Zeeman substates due to dipolar relaxation can be expressed as the following coupled differential equations [46]

$$\frac{\mathrm{d}n_{-J}}{\mathrm{d}t} = -\beta_{\mathrm{dr}}^+ n_{-J}^2 + \beta_{\mathrm{dr}}^- n_{-J} n_{-J+1}$$
(3.20)

$$\frac{\mathrm{d}n_{-J+1}}{\mathrm{d}t} = +\beta_{\mathrm{dr}}^+ n_{-J}^2 - \beta_{\mathrm{dr}}^- n_{-J} n_{-J+1} , \qquad (3.21)$$

where $n_M = \frac{N_M}{V}$ denotes the number density of atoms in the Zeeman substate M. As expected β_{dr}^+ is a source for atoms in the M = -J + 1 = -7 state, but constant optical repumping of the atoms back into the fully spin polarized state with M = -J results in number densities of fully spin-polarized atoms n_{-J} that are much higher than the number densities of atoms in the M = -J + 1 state n_{-J+1} . Taking optical pumping into account and assuming that the optical pumping rate $\Gamma_{op} \gg \beta_{dr}^- n_{-J+1}$, in these rate equations 3.20 and 3.21 the term $\beta_{dr}^+ n_{-J}^2$ stays bigger than the term $\beta_{dr}^- n_{-J} n_{-J+1}$ even if eventually $\beta_{dr}^+ < \beta_{dr}^-$. Therefore, the β_{dr}^- term can be neglected in order to calculate the spin-flip rate due to dipolar relaxations:

$$\dot{N}\Big|_{\rm dr} = \beta_{\rm dr}^+ n_{-J}^2 V \,.$$
 (3.22)

For each $\Delta M = 1$ relaxation the corresponding Zeeman energy is transferred from the external to the internal degree of freedom. Consequently, the energy of the external degree of freedom changes as following

$$\dot{E} = -E_{\rm Z} \left. \dot{N} \right|_{\rm dr} \,. \tag{3.23}$$

In order to calculate the cooling effect on a cloud of trapped atoms, the energy of the gas in the trap is needed. A general trap potential which describes most traps is of the form $U(x,y,z) = a_1 x^{n_1} + a_2 y^{n_2} + a_3 z^{n_3}$. With $\alpha = \sum_i \frac{1}{n_i}$, the energy of the trapped gas results in $E = \left(\frac{3}{2} + \alpha\right) N k_{\rm B} T$ [38]. Neglecting atom loss ($\dot{N} = 0$), the temporal

evolution of the temperature arises [46]

$$\frac{\mathrm{d}T}{\mathrm{d}t} = \frac{\dot{E}}{\left(\frac{3}{2} + \alpha\right) N k_{\mathrm{B}}} = \frac{-E_{\mathrm{Z}} \beta_{\mathrm{dr}}^{+} n_{-J}^{2} V}{\left(\frac{3}{2} + \alpha\right) N k_{\mathrm{B}}},\qquad(3.24)$$

where the number density in the trap is a space-dependent Gaussian distribution $n(\mathbf{r}) = n_0 \exp\left(-\frac{U(\mathbf{r})}{k_{\rm B}T}\right)$ and the trap volume V is temperature dependent [38]. In order to simplify calculations, the mean density

$$\bar{n} = \frac{\int n^2(\mathbf{r}) \mathrm{d}^3 r}{\int n(\mathbf{r}) \mathrm{d}^3 r} = 2^{-\alpha} n_0$$
(3.25)

and the mean trapping volume

$$\bar{V} = \frac{N}{\bar{n}} = 2^{\alpha} V \tag{3.26}$$

are used for further calculations. Here $n_0 = \frac{N}{V}$ denotes the peak density.

Considering a harmonic trap ($\alpha = 3/2$) and the constant repumping of the atoms from M = -7 to M = -8, which results in $N_{-J}/N \approx 1$ and $\bar{n}_{-J}^2 \bar{V} \approx \bar{n}_{-J} N_{-J}$, equation 3.24 reduces to

$$\frac{\mathrm{d}T}{\mathrm{d}t} = \frac{-E_Z \,\beta_{dr}^+ \,\bar{n}_{-J}}{3 \,k_B} \,. \tag{3.27}$$

In figure 3.5 the cooling rate $\frac{dT}{dt}$ divided by the mean density \bar{n}_{-J} is displayed for different magnetic fields in a temperature range from 0 to 50 µK. As the density n = N/V depends on temperature as well as the volume V of the cloud, the cooling rate divided by the mean density \bar{n}_{-J} has been calculated here. This way the temperature dependency of the density has not to be included in the calculation, but still the calculation can be compared with experimental data. Obviously, there is no cooling effect for a magnetic field close to zero (e.g. $B = 1 \mu G$), as no Zeeman energy can be transferred from one energy reservoir to the other. In agreement with the dipolar relaxation rate, the cooling effect is suppressed when lowering the temperature. Furthermore, the higher the magnetic field, the higher the temperature at which the cooling is already suppressed. Besides, the higher the magnetic field, the higher the absolute value of the cooling rate can grow due to the Zeeman energy. However, the temperature must be high enough in order to ensure that the cooling rate is still not suppressed. Consequently, the magnetic field value at which demagnetization cooling performs best changes with temperature.

In figure 3.6 the optimal magnetic field and the corresponding cooling rate divided by the density \bar{n}_{-J} over temperature are shown. Both curves have been calculated numerically comparing the value of $\frac{1}{\bar{n}_{-J}} \frac{\mathrm{d}T}{\mathrm{d}t}$ for different magnetic fields. Interestingly, the optimal magnetic field B_{opt} seems to increase linearly with temperature T by



Figure 3.5: The cooling rate $\frac{dT}{dt}$ divided by the mean density \bar{n}_{-J} of fully spin-polarized atoms for different magnetic fields in a temperature range from 0 to 50 µK. A typical experimental mean density is around $\bar{n} = 5 \cdot 10^{17} \text{ m}^{-3}$, thus a theoretical optimal cooling rate ranges from 0 to $40 \frac{\mu K}{s}$ depending on temperature. Moreover, it depends linearly on the density of the cloud. Therefore, a higher density can also improve the cooling rate as long as three-body losses are negligible.

approximately $18 \frac{\text{mG}}{\mu\text{K}}$.

Experimentally, we can realize minimal magnetic fields of around 20 mG. The cooling rate is theoretically optimal for this magnetic field value at a temperature of $T = 1.\bar{1}\,\mu\text{K}$, but already quite small with $0.03 \cdot 10^{-17} \frac{\text{m}^3 \mu\text{K}}{s} \cdot 10^{18} \frac{1}{\text{m}^3} = 0.3 \frac{\mu\text{K}}{s}$ at a density of around $10^{18} \frac{1}{\text{m}^3}$.

Atoms which are trapped in harmonic potentials undergo the phase transition to a Bose-Einstein condensate (BEC) at a critical temperature T_c which is characterized by [48]

$$k_{\rm B}T_{\rm c} \approx 0.94 \,\hbar\bar{\omega} N_{\rm tot}^{1/3} \tag{3.28}$$

Our maximal mean trapping frequency is $\bar{\omega} = 2\pi (132 \cdot 783 \cdot 905)^{1/3} \text{Hz} = 2\pi \cdot 454 \text{Hz}$ (in the crossed optical trap, see sec. 3.5.1) and total initial atom numbers of $N_{\text{tot}} = 0.5 \cdot 10^6$ are realistic in this trap. Therefore, at least a temperature of $T_c = 1.6 \,\mu\text{K}$ should be reached in order to observe the phase transition to a BEC.

In figure 3.7 the optimal temperature evolution for an optimal magnetic field value is displayed for a realistic initial temperature of $T_0 = 27 \,\mu\text{K}$ and a constant mean density of $\bar{n} = 10^{18} \,\text{m}^{-3}$. The critical temperature is also shown as an orange line. Under these conditions, the time it takes to reach quantum degeneracy can hence be estimated to $t_{\text{BEC}} \geq 4.8 \,\text{s}$. However, this model is very simplified and neglects any heating mechansim. In the following section we will therefore address this kind of limitations.



Figure 3.6: Numerically calculated a) optimal magnetic field and b) cooling rate over density (red curve), both dependent on temperature. Interestingly, the optimal magnetic field increases linearly with temperature: $B_{\rm opt} \approx T \cdot 18 \, \frac{\rm mG}{\mu \rm K}$.



Figure 3.7: a) Temperature evolution (blue line) for b) optimal magnetic field values and an initial temperature of $T_0 = 27 \,\mu\text{K}$ neglecting heating effects and assuming perfect optical repumping to the fully spin-polarized state. Furthermore, a comparison with $T_c = 1.6 \,\mu\text{K}$ (orange line), reveals a time of $t_{\text{BEC}} \ge 4.8 \,\text{s}$ to reach the phase transition to a BEC.

3.4 Optical pumping: experimental restrictions

In the previous section 3.2 demagnetization cooling has been discussed assuming optical pumping which leaves the temperature and the atom number unchanged. However, every absorbed photon causes a kick to the atoms with a total angular momentum $\hbar k_{\rm A}$, where $k_{\rm A} = 2\pi/\lambda_{\rm A}$, due to the spontaneous emission of a single photon with wavelength $\lambda_{\rm A}$. Therefore, in every cooling cycle the atoms gain at least the respective recoil energy.

In this section, I will briefly introduce molecular potentials and the contact interaction before discussing optical pumping at low densities and the corresponding optical pumping rate Γ_{op} . This rate ensures the possibility to constantly recycle the atoms that were thermally excited to higher spin states and is also proportional to the loss rate due to single atom recoil. Further, I will address additional heating and loss mechanisms due to light assisted collisions (LAC) and the possibility to suppress LACs.

3.4.1 Molecular potentials and contact interaction

The interaction of two approaching neutral atoms can be described by a molecular potential. There are bound molecular and unbound quasi molecular states and respective attractive and repulsive molecular potentials. Attractive molecular potentials can be described by a Lennard-Jones potential. This molecular potential consists of a repulsive r^{-12} potential, due to the electrostatic repulsion, and of an attractive van der Waals potential, which scales with r^{-6} and describes the interaction between two induced dipole moments. However, in the low energy limit $(T \to 0)$, for partial waves (l > 0)the molecular potential is superimposed with the centrifugal barrier $\hbar^2 l(l+1)/(mr^2)$, where m is the atomic mass and r the interatomic distance [44] sec. 10.5 and [41]. The atoms are reflected by the centrifugal barrier and do not "experience" the full potential. Therefore, in this regime the scattered wavefunction is a spherical wavefunction that does not depend on the actual potential and only *s*-wave scattering with zero relative orbital angular momentum l = 0 has to be taken into account.

Consequently, at low temperatures the effect of this short-range potential is the same as a hard-sphere potential, where only the scattering length a_s characterizes the scattering. The corresponding contact interaction is the simplest interaction between two particles and given by the potential [44] sec. 10.5 and [41]

$$U_{\rm c}(\boldsymbol{r}) \equiv g\,\delta(\boldsymbol{r})\,,\tag{3.29}$$

where $\delta(\mathbf{r})$ is the Dirac delta distribution and g the coupling strength,

$$g \equiv \frac{4\pi\hbar^2}{m} a_{\rm s} \,, \tag{3.30}$$

which is proportional to the scattering length $a_{\rm s}$. Typical values for the scattering

length in 164 Dy and 162 Dy are [6, 49]

$$a_{\rm s} \stackrel{^{162}\rm{Dy}}{=} 112(10) a_0 \tag{3.31}$$

$$a_{\rm s} \stackrel{^{164}\rm Dy}{=} 92(8) a_0,$$
 (3.32)

with the Bohr radius a_0 .

We can compare the contact interaction with the dipolar interaction by comparing the background scattering length (eq. 3.31 and 3.32) with the dipolar length $(D \stackrel{Dy}{=} 196 a_0, eq. 3.3, [40])$. In fact, the dipolar length is larger than the background scattering length for both isotopes. Hence, the importance of the dipolar interaction in Dysprosium becomes evident.

3.4.2 Low densities: optical pumping rate



Figure 3.8: Visualization of the optical pumping scheme between the lowest Zeemansublevels of the ¹⁶⁴Dy groundstate and excited state 4f⁹ (⁶H⁰) 5d 6s² ⁵I₈. The atoms which have been thermally promoted to $|g, M = -7\rangle$ by inelastic dipolar relaxations are excited by σ^- -polarized light to $|e, M = -8\rangle$ with the scattering rate $\Gamma_{\rm sc}$ and decay to the fully spin-polarized ground-state by spontanous emission. Altogether, they are transferred from $|g, M = -7\rangle$ to $|g, M = -8\rangle$ with the optical pumping rate $\Gamma_{\rm op}$. The numbers on the two groundstate levels give the squares of the Clebsch-Gordan coefficients and thus the probability for the respective transitions to $|e, M = -8\rangle$ and vice versa depicted as a dotted line.

The fully spin-polarized ground state is a dark state for σ^- -polarized light. σ^- polarized light induces a change of the magnetic quantum number $\Delta M = -1$. Furthermore, the total angular momentum does not change in the optical pumping transition $(\Delta J = 0)$. Therefore, a state to which an atom in the fully-spin polarized state could be excited to by a σ^- polarized photon does not exist.

Applying σ^{-} -polarized light atoms in the ground state $|g, M = -J + 1\rangle$ are excited to the fully spin polarized excited state $|e, M = -J\rangle$ before they decay back either into the fully spin-polarized ground state $|g, M = -J\rangle$ or into the state they were excited from $|g, M = -J + 1\rangle$. The probability of both events is given by the square of the Clebsch-Gordan coefficients. As the atoms are not excited to another state from the fully spin-polarized dark state, all the atoms will end up in this state after some time. This way energy is transferred from the spin degree of freedom to the light field, as the spontanously emitted photons have more energy than the optical pumping light.

Since atoms which are thermally excited to higher magnetic sublevels by inelastic dipolar relaxations shall directly be optically repumped to the fully spin-polarized state in order to avoid heating due to dipolar relaxations from energetically higher to lower spin states, the optical pumping rate $\Gamma_{\rm op}$ should exceed the $\beta_{\rm dr}^-$ back relaxation process and will be discussed in this subsection. As there exists a bunch of literature on light-matter interaction [44, 50, 51], I will only briefly discuss the optical pumping rate $\Gamma_{\rm op}$, or rather scattering rate $\Gamma_{\rm sc}$, and its dependence on the pumping power P.

As we are discussing the regime of low densities and other decay channels of the excited state used for optical pumping are rare, we can restrict the atom-photon interaction to a two-level system. The rate at which the two-level system scatters photons reads

$$\tilde{\Gamma}_{\rm sc} = \frac{\Gamma}{2} \frac{I/I_{\rm s}}{4\,\Delta^2/\Gamma^2 + I/I_{\rm s} + 1}\,. \tag{3.33}$$

Here, $\Gamma = 2\pi \gamma$ is the linewidth of the transition to the excited state, $\Delta = \omega - \omega_A$ is the detuning of the laser frequency ω from the atomic resonance frequency ω_A , Iis the intensity of the pumping light and $I_s = \frac{\pi}{3} \frac{h c \Gamma}{\lambda^3}$ is the saturation intensity with wavelength λ of the atomic transition. As actually the system is not a simple two-level system, the rate at which a two-level system scatters photons has to be multiplied with the relevant square of the Clebsch-Gordon coefficient to obtain the rate at which the atoms are promoted from the $|g, M = -J + 1\rangle$ state to the $|e, M = -J\rangle$ state:

$$\Gamma_{\rm sc} = {}^J_J C G^{-1}_{-J+1} \cdot \tilde{\Gamma}_{\rm sc} = 0.11 \cdot \tilde{\Gamma}_{\rm sc}$$
(3.34)

Here the square of the Clebsch-Gordon coefficient ${}_{J}^{J'}CG_{M}^{p}$ designs a transition from $|J,M\rangle$ to $|J',M+p\rangle$, where p = -1,0, +1 is the polarization of the photon. This notation is similar to [46, 52]. The value of the Clebsch-Gordon coefficients have been calculated in [53]. For large detunings, $\Delta \gg \Gamma$, eq. 3.33 and 3.34 simplify to

$$\Gamma_{\rm sc} \stackrel{\Delta \gg \Gamma}{\to} 0.11 \cdot \frac{\Gamma}{2} \frac{I/I_{\rm s}}{4 \, \Delta^2 / \Gamma^2}$$

$$(3.35)$$

In order to obtain the actual rate at which the atoms are pumped back to the fully spin-polarized groundstate, also the Clebsch-Gordon coefficient for the transition from $|e, M = -J\rangle$ to $|g, M = -J\rangle$ has to be taken into account:

$$\Gamma_{\rm op} = {}^J_J C G^0_{-J} \cdot \Gamma_{\rm sc} = 0.89 \cdot \Gamma_{\rm sc}$$
(3.36)

Together with the experimentally measurable power $P = \frac{1}{2}\pi w_x w_y I$ and beam waists of a gaussian beam w_x , w_y in x and y direction [54] we can calculate the optical pumping rates on the experiment.

$$\Gamma_{\rm op} = \frac{0.11 \cdot 0.89}{4\pi} \frac{P \,\Gamma^3}{w_x w_y I_{\rm s} \,\Delta^2} \tag{3.37}$$

As already explained, the optical pumping rate should exceed the rate of dipolar backrelaxation collisions, which cause heating, in order to provide a steadily fully spin-polarized sample. On the other hand, there should not be too much light in the system as it causes loss processes by giving the atoms enough energy to escape the trap [55]. Hence, an optimal optical pumping rate has to be found.

3.4.3 High densities: light assisted collisions



Figure 3.9: a) Schematic diagram of molecular potential energy curves for a quasimolecule which comprises of two colliding atoms. With a red arrow the resonant excitation by blue detuned light with frequency ω at the Condon point R_C to a repulsive molecular state and the atomic transition with frequency ω_A at large distances are shown. A resonant excitation to bound molecular states for red detuning is also possible at sharply peaked distinct detunings. Both cases lead to heating as explained in fig. 3.10. Suppression of light assisted collisions is possible when the detuning is chosen such that the respective Condon point coincides with the position of the first node R_N of the groundstate wave function Ψ_g . b) There are different molecular potentials for the different Zeeman-sublevels. Here, optical pumping with σ^- polarized light to the energetically lowest Zeeman-sublevel is shown.

In optical pumping at low densities, we can assume a picture of separated atoms. This has been done in the previous subsection 3.4.2 and is depicted in fig. 3.9b. Whereas, at higher densities the atoms start to interact (see subsec. 3.4.1). Hence, a distance dependent potential has to be integrated into the model (fig. 3.9a). There are ground state molecular potentials and excited state molecular potentials, as well as bound molecular and unbound quasi molecular potentials. Light assisted collisions (LAC) become possible, when photons can excite atoms to bound molecular or unbound quasi



Figure 3.10: Schematic diagram of molecular potential energy curves for a quasimolecule which comprises of two colliding atoms to explain heating by light assisted collisions. a) Two atoms approach each other with a certain velocity (i) until they reach the distance R_C at which the blue detuned laser light is resonant to the excitation to an unbound quasimolecular state (ii). b) After excitation to this molecular state the atoms move apart and retain the energy of the detuning as kinetic energy (i). During spontanous emission at a distance $R > R_C$, the atoms keep the gained kinetic energy (ii).

molecular states with one atom in the ground and one atom in the excited state (fig. 3.9) [55–61].

Two approaching atoms are resonantly excited by either a blue or a red detuned laser with frequency ω at a certain distance R_C , which is called Condon point, to a (quasi-)molecular state. In the case of red detuning only at sharply peaked distinct frequencies this is possible, whereas in the blue detuned case resonant excitation is possible over a large frequency range.

In both cases the atoms gain kinetic energy as depicted in fig. 3.10 for the blue detuned case. After the excitation to a repulsive quasimolecular state, the atoms slow down until they reach a turning point and accelerate into opposite directions gaining kinetic energy. When they pass the Condon point before spontanous emission back to the groundstate, their final kinetic energy has increased [62]. Applying red detuned light, the atoms can be excited to discrete bound states of the attractive potential, if the laser frequency is tuned near the position of a particular bound level. Next they may decay mostly by spontanous emission to rovibrational states which are no longer trapped or hot atoms are formed [55, 58].

3.4.4 Suppression of light assisted collisions

This heating and atom loss mechanism due to light assisted collisions can be suppressed, when the first node of the ground state wavefunction coincides with the Condon point. Since, at the detuning which corresponds to the node of the ground state wavefunction the Franck-Condon factor diminishs. The Franck-Condon factor $f_{\rm C}$ is the overlap between the ground ψ_g and excited vibrational wavefunction ψ_e . The probability of a transition between the corresponding states in the case of weak excitations and *s*-wave scattering is the following¹ [46, 63]

$$K_{\rm loss} = \frac{2\pi\hbar}{\mu_{\rm red}\,k} [\Omega_{\rm R}(R_{\rm C})]^2 \cdot f_{\rm C}\,. \tag{3.38}$$

It is proportional to the Franck-Condon factor. Hence the probability to excite the molecular state depends also on the Franck-Condon factor [46, 63]. Here, k is the relative collision wave vector of the atoms and $\Omega_{\rm R}(R)$ is the Rabi-frequency. The Rabi-frequency can be rewritten as $\Omega_{\rm R}(R_{\rm C})$ when $\Omega_{\rm R}(R)$ varies slowly [63]. Thus the binary loss rate per atom

$$\Gamma_{\rm bin} = n \cdot K_{\rm loss} \tag{3.39}$$

with density n of the atomic cloud diminishs also when the Franck-Condon factor diminishs.

Here, we are focusing at an intermediate regime $R > R_B$. In this regime the van der Waals dominated groundstate molecular potential is almost constant in distance. It scales with $V_g(R) = -C_6/R^6$. However, the excited state interaction potential, a dipole-dipole potential scaling with $V_e(R) = C_3/R^3$, is still relevant. The characteristic van der Waals length scale is $R_B = \left(\frac{\mu_{\rm red}C_6}{10\hbar^2}\right)^{1/4}$. Hence the ground state wave function varies much more slowly with R than the excited state wave function and in a reflection approximation the excited state wavefunction can be seen as a Dirac delta function $\delta(R - R_C)$. Therefore, the Franck-Condon factor $f_C = |\langle \psi_g(E) | \psi_e(E + \hbar \Delta_A) \rangle|^2$ for blue detuning with frequency Δ_A and $f_C = (\langle \psi_g(E) | \psi_\nu \rangle)^2$ for red detuning with excitation to a state with vibrational quantum number ν and the vibration frequency ν_{ν} reduces to

$$f_{\rm C} = \frac{1}{D_C} |\psi_g(R_{\rm C}, E)|^2 \tag{3.40}$$

red:

blue:

$$f_{\rm C} = \frac{h \,\nu_{\nu}}{D_{\rm C}} |\psi_g(R_{\rm C}, E)|^2. \tag{3.41}$$

Here, $D_{\rm C} = \left| \frac{\mathrm{d}(V_e - V_g)}{\mathrm{d}R} \right| \stackrel{R > R_{\rm B}}{\approx} \left| \frac{\mathrm{d}V_e}{\mathrm{d}R} \right| = \frac{3C_3}{R^4}|_{R=R_{\rm C}}$ is the slope of the difference potential

¹ Here, this probability is given for the case of zero BEC fraction. In order to get this probability in a condensate with fraction x, K_{loss} has to be multiplied by a factor of $\frac{2-x}{2}$: $K_{\text{loss}}^x = \frac{2-x}{2}K_{\text{loss}}$.

$V_e - V_g$ at $R = R_C$ [63].

An approximation of the ground state wave function is [63]

$$\psi_g = \left(\frac{2\,\mu_{\rm red}}{\pi\,\hbar^2\,k}\right)^{1/2}\,a(R)\,\sin(k\,\rho(R))\tag{3.42}$$

where

$$a(R) = 1 - \left(\frac{R_{\rm B}}{R}\right)^4$$
, (3.43)

$$\rho(R) = R - a_{\rm s} - \frac{2}{3} \left(\frac{R_{\rm B}}{R}\right)^4 R, \qquad (3.44)$$

with scattering length a_s of the ground state potential. Combining eq. 3.38, 3.39, 3.40 and 3.42 the binary loss rate per atom for blue detuning reads:

$$\Gamma_{bin} = \frac{4n}{\hbar k^2} \left[\Omega_{\rm R}(R_{\rm C}) \right]^2 \cdot \frac{a^2(R_{\rm C}) \,\sin^2(k\,\rho(R_{\rm C}))}{D_{\rm C}} \,. \tag{3.45}$$

As we are only interested in the area around the zero crossing of the wave function, the small angle approximation can be applied leading to

$$\Gamma_{\rm bin} \stackrel{k\rho \ll 1}{\approx} \frac{4n}{\hbar} \left[\Omega_{\rm R}(R_{\rm C}) \right]^2 \cdot \frac{R_{\rm C}^2}{D_{\rm C}} \cdot g_{\rm C}(R_{\rm C})$$
(3.46)

with

$$g_{\rm C} = \left[\left(1 - \left(\frac{R_{\rm B}}{R_{\rm C}} \right)^4 \right) \left(1 - \frac{a_{\rm s}}{R_{\rm C}} - \frac{2}{3} \left(\frac{R_{\rm B}}{R_{\rm C}} \right)^4 \right) \right]^2 \,. \tag{3.47}$$

Recalling $D_{\rm C} \approx \frac{3C_3}{R_{\rm C}^4}$ and the fact that at the Condon point $R_{\rm C}$ the energy of the detuning $\hbar\Delta$ matches in the case of blue detuning and $V_g(R > R_{\rm B}) \approx 0$ the difference between $V_e(R_{\rm C})$ and $V_e(R \to \infty)$, which leads to $\hbar\Delta = \frac{C_3}{R_{\rm C}^3}$, the binary loss rate reduces to

$$\Gamma_{\rm bin} = \frac{4n}{\hbar} \left[\Omega_{\rm R}(R_{\rm C}) \right]^2 \cdot \frac{R_{\rm C}^6}{3C_3} \cdot g_{\rm C}(R_{\rm C}) \tag{3.48}$$

$$= \frac{4n}{\hbar} \left[\Omega_R(R_{\rm C}) \right]^2 \cdot \frac{C_3}{3(\hbar \Delta)^2} \cdot g_{\rm C}(R_{\rm C}) \,. \tag{3.49}$$

At the Condon point the Rabi-frequency is $\Omega_{\rm R}(R_{\rm C}) = b_{\rm C} \Omega_{\rm R,A}$ with the atomic Rabifrequency $\Omega_{\rm R,A} = 2 \Delta (\tilde{\Gamma}_{\rm sc}/\Gamma)^{1/2}$. Furthermore, $C_3 = f_3 \hbar \Gamma (\lambda_A/(2\pi))^3$, where λ_A is the atomic wavelength of the optical pumping transition and $0 < b_{\rm C} < 2/\sqrt{3}$, as well as $0 < |f_3| < 1/2$ are molecular structure factors [63]. This leads to a binary scattering rate for blue detuning

$$\Gamma_{\rm bin}^{\rm blue} = \Gamma_{\rm bin} = \frac{16 f_3 b_{\rm C}^2}{3 \hbar^2} n \left(\frac{\lambda_{\rm A}}{2\pi}\right)^3 \tilde{\Gamma}_{\rm sc} g_{\rm C}(R_{\rm C}), \qquad (3.50)$$

where only $g_{\rm C}$ depends on $R_{\rm C}$ and thus the detuning Δ . The function $g_{\rm C}$ and hence also $\Gamma_{\rm bin}$ have a minimum around the scattering length $a_{\rm s}$, where the first node of the ground state scattering wave function is located.

In the red detuned case, the binary loss rate per atom is modified due to the fact that only at discrete levels a transition to a molecular state is possible [63]:

$$\Gamma_{\rm bin}^{\rm red} = \Gamma_{\rm bin} \sum_{\nu} \frac{\nu_{\nu} \, \Gamma_{\nu}}{(\Delta - \Delta_{\nu})^2 + (\Gamma_{\nu}/2)^2} \tag{3.51}$$

The normalized binary loss rate $\eta \propto \Gamma_{\rm bin}^{\rm color}/(n(\lambda/(2\pi))^3 \tilde{\Gamma}_{\rm sc})$ of reference [63] is plotted for red and blue detuning for sodium in fig. 3.11. The red-detuned binary loss rate is strongly modulated due to bound state resonances, but for both red and blue detuning the suppression of binary loss when the Condon point coincides with the ground state wave function is evident. Between resonances the binary loss rate for red detuning is smaller than for blue detuning. Therefore the experiments of this thesis have been performed with red detuned optical pumping light.



Figure 3.11: Normalized binary loss rate $\eta \propto \Gamma_{\rm bin}^{\rm color}/(n(\lambda/(2\pi))^3 \tilde{\Gamma}_{\rm sc})$ for red (dashed line) and blue detuning (solid line) in a Bose-Einstein condensate compared to losses due to photon scattering (horizontal dashed line). Image taken from [63].



Figure 3.12: The function $g_{\rm C}$, which is proportional to the binary loss rate $\Gamma_{\rm bin}^{\rm blue}$, is depicted a) in a double-logarithmic plot similar to fig. 3.11 and b) in a linear plot. Parameters for the optical pumping transition in ¹⁶⁴Dy and for an electric dipole-dipole coefficient $C_{3,e}$ which varies from the analytical value of [63] ($C_{3,e} = 0.01 E_{\rm h} a_0^3$ (blue curve)) to values corresponding to experimental and maximal theoretical values in ⁵²Cr [46] ($C_{3,e} = 0.103 E_{\rm h} a_0^3$ (red curve) and $C_{3,e} = 0.161 E_{\rm h} a_0^3$ (black curve)) have been chosen. As the linewidth of the optical pumping transition is $\Gamma = 95$ kHz [34], a detuning of $\Delta = 10$ GHz corresponds to approximately $\Delta = 105000 \Gamma$.

In figure 3.12 the function $g_{\rm C}$, which is proportional to the binary loss rate $\Gamma_{\rm bin}^{\rm blue}$, is depicted in a double-logarithmic and in a linear plot for the optical pumping transition in Dysprosium. Hence the chosen parameters are the following. The wavelength of the transition is $\lambda_{\rm A} = 684 \,\mathrm{nm}$, the natural linewidth is $\gamma = 95 \,\mathrm{kHz}$, the van der Waals coefficient is $C_6 = 2003 E_h a_0^6$ [29], and the scattering length is the background scattering length as the experiments are conducted far from any Feshbach resonance, $a_{\rm s} = a_{\rm bg} = 92(8) a_0 [6, 49]$. Here $E_{\rm h} = 4.360 \cdot 10^{-18} \,\text{J}$ is the Hartree energy and $a_0 = 0.05297 \,\mathrm{nm}$ is the Bohr radius. Those are the parameters which have already been measured in Dysprosium. However, the electric dipole-dipole coefficient $C_{3,e}^{Dy}$ of the excited molecule is not known in Dysprosium. Therefore, the function $g_{\rm C}$ is plotted in figure 3.12 for different electric dipole-dipole coefficients $C_{3,e}^{Dy}$ in order to show the broad range of possible node position with the given knowledge about the atomic constants. According to reference [63] the maximal electric dipole-dipole coefficient should be $C_{3,e}^{\max} = f_3 \hbar \Gamma (\lambda_A/(2\pi))^3 = \frac{1}{2} 0.02 E_h a_0^3$. The correspondig g_C function is depicted as a blue curve in fig. 3.12. In chromium $C_{3,e}^{Cr}$ has been determined experimentally to $1.46 E_{\rm h} a_0^3 < C_{3,e}^{Cr} < 1.9 E_{\rm h} a_0^3$ and in an atom loss spectroscopy measurement to determine the node position of the ground state wavefunction it has been fixed in a fit to $C_{3,e}^{Cr} = 1.53 E_{\rm h} a_0^3$ [46]. As $C_{3,e} \propto \Gamma \lambda_{\rm A}^3$ [63], in dysprosium this would correspond to $C_{3,e}^{Dy} = 0.103 E_{\rm h} a_0^3$. The corresponding $g_{\rm C}$ function is shown as a red curve in fig. 3.12. The maximal dipole-dipole coefficient estimated in chromium is $C_{3,e}^{\text{Cr,max}} = 2.4 E_{\text{h}} a_0^3$ It corresponds in dysprosium to $C_{3,e}^{Dy} = 0.16 E_{\rm h} a_0^3$. The corresponding $g_{\rm C}$ function is shown as a black curve in fig. 3.12.

For this range of possible $C_{3,e}$ -coefficients the nodal position of the groundstate wavefunction should be located at a detuning $0.5 \text{ GHz} < |\Delta/(2\pi)| < 8.4 \text{ GHz}$.

The difference between the analytical value in reference [63] and the experimental and theoretical values given in [46] is due to spin-orbit coupling which has been neglected in [63]. In dysprosium, an even stronger spin-orbit coupling is expected as the inelastic dipolar relaxation collisions which quantify this spin-orbit coupling are about one order of magnitude higher in dysprosium than in chromium (see. section 3.2). This is a hint that the node is localized at a larger detuning > 5 GHz.

Furthermore, the validity of the assumption of s-wave scattering (see eq. 3.38) is doubtful due to the strong spin-orbit coupling. Next, in ¹⁶⁴Dy the background scattering length $a_{bg} = 92(8) a_0$ [6, 49] is smaller than the dipole length scale which characterizes the magnetic dipole-dipole interaction strength $D \equiv \frac{\mu_0 \mu^2 m}{8 \pi \hbar^2} \stackrel{Dy}{=} 196 a_0$ (see eq. 3.3) [24]. Hence, the nodal position of the groundstate wavefunction which is located around the scattering length $a_s = a_{bg}$ is still dominated by the magnetic dipole-dipole interaction. Consequently, the approximation of a van der Waals dominated groundstate molecular potential can also be questioned.

An answer could be found by further experimental, i.e. photoassociation spectroscopy, and theoretical investigations, i.e. Movre-Pichler model, which go beyond the scope of this Master's thesis.

Moreover, a description of the system with one possible ground and one possible excited state potential has to be reviewed. This is already the case for ⁵²Cr [46] and of course also for the even more complicated element dysprosium. Different ground state wave functions of different ground state potentials have different node positions. Likewise, different excited state potentials have different Condon-points for a fixed laser detuning Δ . Thus, there is no excitation laser detuning where all Condon-points coincide with the node position. As a result the binary loss rate does not go to zero and the position of its minimum which corresponds to the node position of the ground state wave function is washed out [46].
3.5 Experimental realization

3.5.1 Preparation of the atomic cloud

Before demagnetization cooling several cooling steps are applied to reach starting temperatures of around $30\,\mu$ K. These initial cooling steps have already been explained exhaustively in a Master's thesis [5]. Furthermore, the basics of laser cooling, as well as magneto-optical trapping are explained in several textbooks, e.g. [44]. Hence, we will only give a short overview and then continue the explanations in the optical-dipole trap (ODT), where we perform demagnetization cooling.



Figure 3.13: Schematic view of the apparatus. Starting at the bottom, dysprosium atoms are emitted by an effusion cell. Next, they are transversally cooled by two retro-reflected beams and decelerated in a Zeeman slower (421 nm, blue arrows). In the following, the atoms are captured and further cooled to temperatures around $10\,\mu$ K by a narrow-line MOT (626 nm, orange arrows). Finally, they are loaded into an ODT (red arrow) and transferred into the glass cell by moving the focusing lens ($f = 1250 \,\text{mm}$) of the ODT on an translational stage.

A schematic view of the whole apparatus is shown in figure 3.13. At first, the atoms

are vaporized in a high-temperature effusion cell at temperatures of $1200 - 1250^{\circ}$ C. After being emitted from the oven, they are transversally cooled by two retro-reflected beams which create a two-dimensional optical molasses and then enter the Zeeman slower (ZS). For both cooling techniques we use the broad 421 nm transition with $\Gamma_{421} = 2\pi \cdot 32.2$ MHz [36] and saturation intensity $I_{s,421} = \pi h c \Gamma/3\lambda^3 = 56.4 \text{ mW/cm}^2$. For transversal cooling, the beams are red-detuned at $\Delta = -0.3 \Gamma_{421}$ and elliptically shaped in order to increase the spatial overlap with the atomic beam ($w_z = 6.8 \text{ mm}$ and $w_r = 1.7 \text{ mm}$, where z- and r-direction are chosen with respect to the atomic beam). Typically, we apply a total power of $P_{\text{trans}} = 200 \text{ mW}$ which leads to an intensity $I_{\text{trans}} = 4.5 I_{s,421}$ per beam. Thereby, we gain typically a factor of 2.5 in MOT atom number. For the subsequent ZS, the beams are also red-detuned at $\Delta = -17.5 \Gamma_{421}$ and focused into the effusion cell. At the MOT position its diameter is estimated to 18 mm. Together with a beam power of $P_{\text{ZS}} = 90 \text{ mW}$, we calculate the spread of velocity to be 600 m/s, which allows us to capture around 75 % of the initially Boltzmann-distributed atoms and decrease the velocity of the atoms to $v \approx 10 \text{ m/s}$.

Now, the atoms are slow enough to be captured by the MOT, which is setup with three retroreflected beams. Here, we use the narrow 626 nm transition with $\Gamma_{626} = 2\pi \cdot 136$ kHz [64], saturation intensity $I_{626} = 72 \,\mu \text{W/cm}^2$ and Doppler temperature $T_{\text{D},626} = 3.3 \,\mu \text{K}$. The beams have a diameter of 22.5 mm and an intensity of $I_{\text{MOT}} \approx 240 \, I_{\text{s},626}$ per beam. At a magnetic field gradient of $\nabla B = 3 \,\text{G/cm}$ and a detuning of $\Delta_{\text{MOT}} = -34 \, \Gamma_{626}$, we load typically 10^8 atoms with a temperature around $500 \,\mu \text{K}$ in 4 s into the MOT. Next, we compress the MOT in order to achieve colder temperatures and decrease the trapping volume. For this purpose, we decrease the intensity to $I_{\text{MOT}} \approx 0.24 \, I_{\text{s},626}$, the magnetic field gradient to $\nabla B = 1.5 \,\text{G/cm}$, and the detuning to $\Delta_{\text{MOT}} = -5 \, \Gamma_{626}$ in 170 ms. Thereby, the atomic cloud is spatially compressed and the final temperature is lowered to $12 \,\mu \text{K}$. Besides, the atoms are optically pumped to the energetically lowest Zeeman state M = -8 by the MOT. This prevents inelastic relaxation collisions to lower substates which would lead to a heating of the atomic cloud.

Next, the atoms may be transferred from the MOT chamber to the optical dipole trap in the science chamber. This is done with an optical tweezer. In fact, we use a single-beam optical dipole trap which is created by a broadband fiber laser¹ operating at 1070 nm. We measure a beam waist of $w_0 = (37.3 \pm 1.2) \,\mu\text{m}$ which implies a Rayleigh length of $z_{\rm R} = 4.1 \,\text{mm}$ and estimate the maximal beam power in the glass cell to 72 W. Since the atom cloud in the MOT is a factor of 10 larger than the trapping potential created by the transport beam [5], we load it $\Delta x = 15 \,\text{mm} \approx 3 \,z_{\rm R}$ away from the beam's trap minimum ("funnel" method). Here, the trap depth is decreased and the trapping area is increased, each by the same factor of 17.6 [5]. During the compression of the MOT, we move the trap minimum of the transport beam to the above-mentioned position by moving the $f = 1250 \,\text{mm} \,\text{lens}^2$ which focuses the beam

¹ Laser IPG YLR-100-WP-WC, $\lambda = 1070$ nm, $P_{\text{max}} = 100$ W.

² The lens is mounted on a computer-controlled air-bearing translation stage Aerotech ABL15040 with 40 cm translation range and $0.5 \,\mu$ m accuracy.

on the atoms. After the compression of the MOT, we turn the transport beam on and hold both traps for 120 ms. Next, we turn off the MOT beams and field gradient. In order to prevent the atoms from oscillating after being loaded in an out-of-equilibrium position in the transport trap, we move the trap minimum rapidly to the prior MOT center. However, the sample heats up to about 170 μ K and contains typically 15 \cdot 10⁶ atoms in the transport beam.

Finally, we transfer the atoms from the MOT chamber over a range of 375 mm to the science chamber. As the laser we use for transport cannot be used for evaporative cooling¹, we load $N = 5 \cdot 10^6$ atoms at a temperature of $T = 120 \,\mu\text{K}$ with an efficiency of 50% in ODT 1.



Figure 3.14: Schematic view of the science chamber. The transport beam (red arrow) transfers the atoms from the MOT chamber to the glass cell. Here, they are loaded either in an optical-dipole trap created by ODT 1 or in a crossed optical-dipole trap (cODT) created by ODT 1 and 2. The optical pumping beam (red arrow) for demagnetization cooling is applied in z-direction, parallel to the phase-contrast imaging beam. Alternatively to demagnetization cooling, post-cooling with a laser beam which is red-detuned to the 626 nm transition (orange arrow) and forced evaporative cooling can be applied. For time-of-flight measurements, absorption imaging along the y direction (blue arrow) is used. In addition, phase-contrast imaging (blue arrow) with the high-NA objective along the z direction can be performed for in-situ images. Further, the objective can be used to write tailored potentials with an electro-optical deflector (EOD) system (green arrow) [5]. Above and beneath the glass cell, there are two coils in Helmholtz configuration for generation of magnetic fields up to 600 G.

¹ The laser contains frequency modes which drive two-photon Raman transitions [5]

In figure 3.14 a schematic view of the science chamber and the setup for optical trapping, optical pumping, further cooling and imaging is displayed.

In the science chamber, ODT 1 and eventually a superimposed ODT 2 provide optical trapping of the atoms. When ODT 2 is superimposed to ODT 1, they form a crossed optical-dipole trap. For both traps, a single-mode laser ¹ operating at $\lambda = 1064$ nm is used. In order to form two traps, the beam is split into two paths. The maximum power of ODT 1 in the glass cell is P = 9.7 W. Together with a beam waist in x-and y-direction respectively $(w_x, w_y) = (35.7, 33.5) \,\mu\text{m}$ we obtain maximal trapping frequencies of $(f_x, f_y, f_z) = (5.6, 783, 837)$ Hz. The maximum power of ODT 2 in the glass cell is P = 7.75 W. With beam waists $(w_x, w_y) = (103.5, 39.5) \,\mu\text{m}$ and combined with ODT 1 to a crossed optical-dipole trap we obtain maximal trapping frequencies of $(f_x, f_y, f_z) = (132, 783, 905)$ Hz. After a first evaporation step in the cODT, in which the power of the first ODT is reduced to 25% of the maximal power, we obtain trapping frequencies of $(f_x, f_y, f_z) = (132, 392, 442)$ Hz. For further information on the measurement of the trapping frequencies see [5].

Up to now we use post-cooling with the 626 nm transition, once the atoms are loaded in ODT 1. Here, we illuminate the atoms in the trap with a beam with very low intensity $I = 0.035 I_{sat} = 2.52 \,\mu\text{W/cm}^2$ and a frequency which is red-detuned to the 626 nm transition. This way the atoms are excited with less energy than the transition energy and release in average the difference in energy during spontanous emission. A maximal cooling efficiency of $\chi = 6.7$ is reached by this post-cooling [5] (see sec. 3.5.4.6).

In all experiments described in this thesis, the atom number and temperature is probed by standard time-of-flight absorption imaging.

¹ Laser Coherent Mephisto MOPA 55W

3.5.2 Optical setup of the optical pumping laser system

The full potential of demagnetization cooling may be released by using optical pumping which constantly recycles the atoms from higher Zeeman states to the fully spin-polarized state with $m_J = -J$. In this subsection, I will describe the laser system we set up for optical pumping.

We use a 684 nm transition from the ground state [Xe] $4f^{10} 6s^2 {}^{5}I_8$ with even parity to a $4f^{9}({}^{6}H^{0})5d 6s^2 {}^{5}I_8$ excited state (see fig. 2.1) for optical pumping. In the level scheme of fig. 2.1, it can also be seen that there is a 953 nm transition which can act as a decay channel for the optical pumping transition to the $4f^{10} 6s^2 {}^{5}I_7$ state. However, this 953 nm transition is very narrow [33]. Nevertheless, the efficiency of optical pumping could be restricted by it.

The 684 nm laser system has been set up and studied in a Bachelor's thesis [65], hence I will explain it only briefly and focus on the changes which have been made since then. In figure 3.15 the current laser setup is shown schematically.



Figure 3.15: Scheme of the 684 nm laser setup used for optical pumping.

For the 684 nm light we use a tuneable diode laser¹ with a maximal power of 17 mW. It passes a pair of prisms to obtain a circular instead of an elliptical beam shape. Then, a small fraction of the power is sent to a wavemeter. Next, the beam is split into two main branches. One finally goes on to the experiment and the other branch is used for active frequency stabilization.

¹ Toptica DL 100 pro Design

3.5.2.1 Optical setup to the experiment

Acousto-optic modulator (AOM) In the path for the experiment an AOM¹ with an operating frequency of 100 MHz is used in double pass configuration (see fig. 3.15). Two lenses ensure in this configuration a minimization of the beam shift when the input frequency is changed. Thus, the coupling into the fiber is more stable.



Figure 3.16: Scheme of the 684 nm laser system to the experiment.

Intensity stabilization On the experimental table the light passes a $\lambda/2$ -waveplate and a polarizing beamsplitter in order to ensure π -polarized light (see fig. 3.16). Then we direct 8% of the outcoupled light by a pellicle beamsplitter² to a photo diode³. In the following, we use only the integrating part of a PID controller to stabilize the intensity on the photodiode and thus also on the experiment at a certain value. The intensity is controlled by the power of the rf-frequency for the AOM. As only π -polarized light is used for intensity stabilization, it is guaranteed that the intensity of σ -polarized light is stabilized.

After having passed the pellicle beamsplitter, the π -polarized light is converted into σ^- -polarized light by a $\lambda/4$ -waveplate. The beamwaists are here $w_x = (1.26 \pm 0.02) \text{ mm}$ and $w_y = (1.30 \pm 0.02) \text{ mm}$. Finally, the optical pumping beam is directed over two coupling mirrors and a dichroic mirror⁴ to the science chamber. The dichroic mirror is used, as for optical pumping (684 nm) and for phase-contrast imaging (421 nm) the same optical access to the science chamber is used.

¹ Crystal Technology, 3100-125

² Thorlabs BP108, uncoated for 8:92 (R:T) split ratio

³ Thorlabs PDA36A-EC, switchabel gain detector (0 - 70 dB, 8 steps), 10 MHz bandwidth

⁴ Edmund Optics $#47 - 266, 45^{\circ}$ red dichroic filter

3.5.2.2 Optical setup to ULE cavity

We couple the 684 nm light to an ultra low expansion (ULE) cavity¹ for active frequency stabilization (see sec. 3.5.3). Here, I will describe the optical setup we use for this purpose (see fig. 3.15 on the right-hand side).

After having passed a fiber EOM, we use a 50:50 beamsplitter to couple 50% of the light to a self-made photo diode in order to measure the laser power in front of the ULE. In the following, we utilize another 50:50 beamsplitter and a polarizing beam splitter in order to couple not only the 684 nm, but also the 842 nm and 626 nm laser beams into the ULE cavity. Finally, a power of 0.4 mW of the 684 nm laser is coupled into the ULE cavity.

On the other side of the first 50:50 beamsplitter the back reflected light coming from the ULE is measured with another self-made photodiode. This signal is used to create the error signal for active frequency stabilization with the commercial Pound-Drever-Hall module² (see sec. 3.5.3.2).

3.5.3 Frequency stabilization and adressing of dysprosium atoms

3.5.3.1 Characterization of ULE cavity

In order to characterize the cavity which is used for frequency stabilization, we determine its free spectral range (FSR) and linewidth by measuring the spectrum.

We record the transmission spectrum on a photodiode³ by scanning the laser frequency continuously. Whenever the frequency is on resonance with the cavity, a signal is transmitted and becomes visible as a resonance peak in the spectrum. In order to determine the corresponding relative frequency, the frequency-axis is calibrated on side peaks at ± 25 MHz from the main peak which have been created by a Toptica Pound-Drever-Hall (PDH) module (see sec. 3.5.3.2).

In figure 3.17 a resulting spectrum with two main peaks and their respective side peaks is displayed.

In order to determine the linewidth and the FSR of the cavity, we fit a Lorentzian

$$I_q(\nu) = A_q \frac{\Delta_{1/2}^2}{\Delta_{1/2}^2 + 4(\nu - \nu_q)^2}$$
(3.52)

to each resonance with frequency ν_q by adjusting the amplitude A_q and the full width at half maximum (FWHM) $\Delta_{1/2}$. The difference in frequency between two main peaks corresponds to the FSR. The fit in figure 3.17 reveals a FSR of $\nu_{\rm FSR} = (1527 \pm 1)$ MHz and linewidth of $\Delta_{1/2} = (184.5 \pm 0.2)$ kHz. However, this fit overestimates frequency, because of the discrepancy between fit and data at the first resonance. In order to

¹ Stable Laser Systems

² Toptica PDD 110

³ Thorlabs PDA10A



Figure 3.17: Measurement to determine the resonance frequency.

obtain a more accurate value, several images and fits are taken. The resulting mean values with standard deviation are a FSR of $\bar{\nu}_{\text{FSR}} = (1505 \pm 31) \text{ MHz}$ and a linewidth of $\bar{\Delta}_{1/2} = (181.9 \pm 3.7) \text{ kHz}$.

This value differs only by 0.5% from an expected FSR of $\nu_{\rm FSR} = 1498 \,\mathrm{MHz}$ [65]. Yet, the theoretical linewidth $\Delta_{1/2} = 78 \,\mathrm{kHz}$ of the resonator is about half the experimentally measured linewidth. This is due to the fluctuation of the linewidth of the laser in time $\Delta_{\rm L} > \Delta_{1/2}$, which leads to a broadening of the resonance peaks and a deviation from a Lorentzian shape.

Determining the photon lifetime $\tau_{\rm p}$ of a resonator offers the possibility to determine its linewidth without interfering with the time-dependent laser's linewidth. The photon lifetime is the time it takes to form a standing wave in the resonantor. It can be measured by switching off the frequency stabilized laser very quickly. The decrease in intensity as a function of time is fitted by an exponential decay of the form $I(t) = A \exp(-(t - t_0)/\tau_{\rm p})$ with the fitting parameters A, t_0 and the photon lifetime $\tau_{\rm p}$. We could estimate the photon lifetime by this method to

$$\tau_{\rm p} = (1.41 \pm 0.13)\,\mu{\rm s}\,,$$

which corresponds to a linewidth of

$$\Delta_{1/2} = (113 \pm 11) \,\mathrm{kHz}$$

3.5.3.2 Active frequency stabilization

The linewidth of the optical pumping transition has been determined in a Master's thesis [33, 34] to $\Gamma_{684} = (95 \pm 13)$ kHz. The linewidth of the laser should be narrower than the linewidth of the transition in order to address the transition precisely. Therefore,

active frequency stabilization of the laser is required.

For active frequency stabilization we couple the 684 nm light to an ultra low expansion (ULE) cavity¹ and use a Toptica Pound-Drever-Hall (PDH) module PDD110 to generate smaller side peaks at ± 25 MHz from the main peak (see fig. 3.17). Moreover, we use the PDH module to produce an error signal by multiplying in a mixer the reflected signal of the ULE cavity with the frequency which is used to create the side peaks. A typical error signal can be seen in figure 3.18.



Figure 3.18: Typical errorsignal. Here a moving average over two neighbouring values has been applied.

In the following, we use all three parts of the PID controller of a FALC module² to stabilize the laser on the steap slope of the main peak of the error signal.

By comparing the transmission signal of the frequency stabilized laser with the resonance peak of the ULE cavity, the linewidth of the frequency stabilized laser can be estimated. The transmission signal of the frequency stabilized laser is 0.57 V and varies ± 0.2 V. Comparing this with a resonance peak of the ULE cavity with a linewidth of $\Delta_{1/2} = 113$ kHz and a signal amplitude of 0.59 V the laser linewidth is $\Delta \nu = (35.9 \pm 1.9)$ kHz. Consequently, the linewidth of the frequency stabilized laser is narrower than the transition linewidth ($\Gamma_{684} = (95 \pm 13)$ kHz) as required.

3.5.3.3 Frequency tuning

Furthermore, in order to find an optimal detuning from the transition frequency to suppress light assisted collisions (see sec. 3.4.3 and 3.4.4) we need to frequency tune the laser in the GHz range. However, the resonances of the ULE cavity are fixed at certain frequencies as the length of it is very stable and cannot be changed. Therefore, we use a broad band fiber EOM in order to tune the laser frequency.

¹ Stable Laser Systems

² FALC 110 from Toptica

Electro-optical modulator (EOM) In an EOM an optical medium¹ which exhibits the linear electro-optical effect is used. Crystals of this kind change their refractive index nlinearly in an external electric field. In EOMs the electric field is applied vertically to the light propagation axis and induces a phase shift between light polarized in the same direction and light polarized in the vertical direction. The voltage required for inducing a phase change of π is called half-wave voltage U_{π} . In a fiber EOM a polarization maintaining fiber ensures the application of light which is polarized parallel to the electric field. When we apply a sinusoidally varying potential voltage to the EOM, the resulting sinusoidally varying phase shift produces sidebands at the corresponding frequency and its higher harmonics. The strength of the sidebands depends on the amplitude of the potential voltage U and is given by Bessel functions $J_k(U)$, where knumbers the respective higher harmonic [54, 66].

The typical half-wave voltage of the EOM² used in this thesis is $U_{\pi} = 6$ V. In order to obtain first order side peaks which are approximately as strong as the main peak and invisible second order side peaks, a power of 13 dBm, which corresponds to a peak voltage of U = 1.4 V, is applied to the signal input of the EOM (see fig. 3.19b for the resulting errorsignal). At lower powers the side peaks vanish (see fig. 3.19a) and at higher powers higher order side peaks appear and the main peak vanishs (see fig. 3.19c). A DDS-board provides frequencies between several MHz and 500 MHz. Hence, a tuning from 25 MHz < f < 500 MHz should be possible.

We use an EOM which has been designed for a wavelength of $\lambda = (705 \pm 25)$ nm. We couple 6.7 mW into the EOM and 1.6 mW out of the EOM, which corresponds to an insertion loss of $-6 \,\mathrm{dB}$ and to the value specified in the data sheet of the EOM.

We use the sidebands shown in figure 3.19b as a stable tunable frequency reference. By stabilizing the laser on one of those sidebands and changing the EOM frequency, we tune the laser frequency going on the experiment.

3.5.3.4 Determination of resonance frequency

The laser is stabilized on the blue detuned side peak of the resonance frequency of the ULE cavity which is closest to the wavelength of the pumping transition in vacuum (683.731 nm, [67]). Furthermore the laser frequency is further shifted with a double pass AOM by $+2 \cdot 100$ MHz.

We aligned the optical pumping beam on the atoms in the ODT by maximizing the atom loss at a power of roughly $800 \,\mu\text{W}$ before the $\lambda/4$ -waveplate and an illumination time of 100 ms.

Further, the resonance frequency of the pumping transition has to be found. Therefore, we apply at first a laser power of $400 \,\mu\text{W}$, which corresponds to an optical pumping intensity on the atoms of $I_{\rm op} = (47 \pm 10) \,\frac{\mu\text{W}}{\text{mm}^2}$ for 50 ms. Then we apply for a higher resolution a laser power of $100 \,\mu\text{W}$, which corresponds to an optical pumping intensity on the atoms of $I_{\rm op} = (12 \pm 3) \,\frac{\mu\text{W}}{\text{mm}^2}$ for 10 ms. This way the frequency difference from

¹ here: Lithiumniobat (LiNbO₃)

² phase modulator for $\lambda = 705 \text{ nm} (\text{PM } 705)$ from Jenoptik



Figure 3.19: Errorsignal including the additional sidebands created by the fiber EOM, here at a frequency of $\pm 200 \text{ MHz}$ and including a moving average over two neighbouring values. a) The side peaks are very weak at a power of 10 dBm, which corresponds to a peak voltage of U = 1.0 V. b) At a power of 13 dBm, which corresponds to a peak voltage of U = 1.4 V, the first order side peaks are as strong as the main peak and higher order side peaks vanish. c) The main peak vanishs and higher order side peaks appear at a power of 16 dBm, which corresponds to a voltage of U = 2.0 V.

the cavity resonance to the resonance of the optical pumping transition is determined to 427.2 ± 5 MHz (see fig. 3.20).



Figure 3.20: Measurement to determine the resonance frequency.

3.5.4 Analysis of demagnetization cooling

After having determined the resonance position, we can start with actual demagnetization cooling. Therefore, we load the atoms into ODT 1 (see sec. 3.5.1) with trapping frequencies $\bar{\omega} = 2\pi (5.6 \cdot 783 \cdot 837)^{1/3}$ Hz and wait $t_{\text{evap}} = 1.5$ s in order to ensure that the following temperature evolution does not interfere with residual evaporative cooling (see sec. 3.5.4.6). After that, the sample contains approximately $1.3 \cdot 10^6$ atoms at a temperature of roughly $T_0 = 27 \,\mu\text{K}$ and a mean density of $n_0 = 6 \cdot 10^{17} \,\text{m}^{-3}$. These are the starting conditions of all measurements presented in the next four subsections 3.5.4.1 to 3.5.4.4. Next, we apply optical pumping light with a certain intensity and a certain magnetic field parallel to the optical pumping beam.



Figure 3.21: a) Temperature and b) atom number evolution over time during demagnetization cooling with optical pumping rates of $\Gamma_{\rm op}^{\rm h} = (99 \pm 13) 2\pi$ Hz (green), $\Gamma_{\rm op}^{\rm l} = (20 \pm 3) 2\pi$ Hz (orange), and no optical pumping light (blue) and for a magnetic field of $B_z = 385$ mG. The data is presented as dots. The temperature evolution due to the theoretical cooling rate at B = 385 mG which does not include any heating mechanism, but the experimental densities is shown in a) as solid lines. Temperature and atom number evolution are highlighted by guide to the eyes (dashed line).

In figure 3.21 the temperature and atom number evolution over time is shown for two different optical pumping rates: $\Gamma_{\rm op}^{\rm h} = 2\pi (99 \pm 13) \,\text{Hz}$ and $\Gamma_{\rm op}^{\rm l} = 2\pi (20 \pm 3) \,\text{Hz}$ and no optical pumping light. Here a magnetic field of 385 mG has been applied in z-direction. The theoretical temperature evolution (see figure 3.21 a)), which has been calculated taking an optimal cooling rate¹ at the applied magnetic field B = 385 mG, are also displayed in the graph as solid lines. The difference between the two theoretical temperature evolutions is due to different experimental densities, which have been taken into account in the calculation. Even though we observe a first cooling effect in the experimental data in comparison to no optical pumping light, there is a discrepancy between experimental and theoretical optimal temperature evolution. There are also some possibilities for optimization, which will be presented in the following sections.

3.5.4.1 Optical pumping rate

One possibility for improvement is to find an optimal optical pumping rate Γ_{op} . On the one hand, too many photons lead to heating and atom loss from the trap by giving the atoms energy, eventually enough energy to leave the trap [55]. On the other hand, the optical pumping rate should exceed the rate of dipolar backrelaxation collisions, which cause heating (see sec. 3.4.2). In figure 3.21 the evolution of temperature and atom number over time is shown for two different optical pumping rates $\Gamma_{\rm op}^{\rm h} = (99 \pm 13) 2\pi \, \text{Hz}$ and $\Gamma_{\rm op}^{\rm l} = (20 \pm 3) 2\pi \, \text{Hz}$ and no optical pumping light. For the higher optical pumping rate $\Gamma_{\rm op}^{\rm h}$, one observes a loss of atoms of $(-170 \pm 20) \cdot 10^3 \, \frac{\text{atoms}}{\text{s}}$ (see fig. 3.21 b)). In an exponential fit the lifetime is estimated to $\tau = (6.67 \pm 0.62)$ s. For the lower optical pumping rate Γ_{op}^{l} the atom number stays approximately constant and is comparable to the case of no optical pumping light especially after 3 s of wait time. Due to higher experimental densities for Γ_{op}^{l} , one would expect a faster decrease in temperature as depicted in the theoretical cooling rates in figure 3.21 a). Nevertheless, the experimental data does not show a clear improvement in the case of the lower optical pumping rate Γ_{op}^{l} compared to the higher optical pumping rate $\Gamma_{\rm op}^{\rm h}$. In order to compare the performance of the cooling effect, we calculate the cooling rate $\frac{dT}{dt}$, which estimates the time scale on which cooling takes place. We extract it from the data by a linear fit to the temperature evolution as a function of time. The cooling rate in the first $t_{\text{demag}} = 2 \text{ s performs even}$ better for the higher optical pumping rate with $\frac{dT}{dt}\Big|_{\Gamma_{\text{op}}^{\text{h}}, 2s} = (-4.4 \pm 0.5) \frac{\mu \text{K}}{s}$ than for the lower optical pumping rate $\left(\frac{dT}{dt}\Big|_{\Gamma_{\text{op}}^{1},2s} = (-3.7 \pm 0.6)\frac{\mu\text{K}}{s}\right)$. However, for $t_{\text{demag}} = 3 \text{ s}$ of demagnetization cooling the cooling rates for higher and lower optical pumping rate resemble each other with $\frac{dT}{dt}\Big|_{\Gamma_{\text{op}}^{1,h},3s} = (-3.2 \pm 0.3)\frac{\mu\text{K}}{s}$. Still, as the evolution of the atom number is more favorable in the cost of the d the atom number is more favorable in the case of the lower optical pumping rate, the following experiments in this trap configuration are performed with the lower optical pumping rate $\Gamma_{\rm op}^{\rm l} = (20 \pm 3) 2\pi \, \text{Hz}.$

At the end of subsection 3.5.4.5 heating and loss mechanisms due to optical pumping light are discussed in the scope of a measurement which probes a range of different optical pumping rates.

¹ here optimal means that heating backrelaxations are neglected

3.5.4.2 Magnetic field



Figure 3.22: a) Temperature and b) atom number after $t_{\text{demag}} = 2 \text{ s}$ of demagnetization cooling with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (orange) and without optical pumping light (blue) and for a magnetic field ranging from $B_z = 25$ to 385 mG. The shaded regions indicate the typical variations of the respective parameters.

Another parameter which can be optimized is the magnetic field in z-direction. As it can be seen in figure 3.6, for each temperature there is a magnetic field for which the cooling rate becomes optimal. At a temperature of roughly $T_i = 27 \,\mu\text{K}$, which was the starting temperature in the measurement presented in figure 3.21, the optimal magnetic field without taking any heating meachanism into account is around 400 mG (see fig. 3.6). Therefore we chose a value close to this magnetic field in this first measurement (see fig. 3.21). However, experimentally a lower magnetic field value seems to be more favorable. In figure 3.22 the temperature and atom number after $t_{\text{demag}} = 2 \text{ s of}$ demagnetization cooling with an optical pumping rate of $\Gamma_{\rm op} = (20 \pm 3) 2\pi \, \text{Hz}$ (orange) and magnetic field in z-direction ranging from $B_z = 25$ to $385 \,\mathrm{mG}$ is shown. This graph shows that around $B_z = 200 \text{ mG}$ the temperature has decreased from approximately $T_f = 18 \,\mu\text{K}$ at $385 \,\mathrm{mG}$ to $16 \,\mu\text{K}$ and the atom number has increased from roughly 10^6 to $1.2 \cdot 10^6$ for the same magnetic field values. In comparison with the case of no optical pumping light (blue), the atom number has only decreased marginally between $B_z = 100$ and 350 mG and the temperature is below the value of no optical pumping light in this whole magnetic field range.

Another important parameter is the phase space density. In fact, the objective of a cooling process is not only to decrease temperature or stay at high atom numbers, but



Figure 3.23: Phase space density \mathcal{D} of the atoms after $t_{\text{demag}} = 2 \text{ s}$ of demagnetization cooling with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (orange) and no optical pumping light (blue) and for a magnetic field ranging from $B_z = 25$ to 385 mG. The cooling performs best around $B_z = 225 \text{ mG}$. The shaded regions indicate typical variations of the phase space density.

to increase the phase-space density

$$\mathcal{D} = n \,\lambda_{dB}^3,\tag{3.53}$$

where n = N/V is the atomic density and

$$\lambda_{dB} = \sqrt{\frac{2\pi\hbar^2}{mk_{\rm B}T}} \tag{3.54}$$

is the de Broglie wavelength which depends on the atomic mass m, the Boltzmann constant $k_{\rm B}$ and the temperature T. Hence, the phase-space density unifies a decrease in temperature and an increase in atom number in an increase in phase-space density. In a harmonic trap, the phase space density is given by [68]

$$\mathcal{D} = N \left(\frac{\hbar \bar{\omega}}{k_{\rm B} T}\right)^3 \tag{3.55}$$

where $\bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3}$ is the geometric mean of the trap frequencies in x-, y- and z-direction.

In figure 3.23 the corresponding phase space density \mathcal{D} (3.53) is plotted over the magnetic field in z-direction. It is higher than in the case of no optical pumping light in the whole magnetic field range mentioned-above. Furthermore, it seems to be optimal around $B_z = 225 \,\mathrm{mG}$. Hence, this value is chosen for the following experiments which start at the same initial temperatures.

3.5.4.3 Polarization



Figure 3.24: Level scheme of the lowest Zeeman levels of the optical pumping transition. Not fully σ^{-} , but partially π and σ^{+} -polarized light leads to an excitation to higher Zeeman-substates. The backrelaxation to lower Zeeman-substates leads to a heating of the atomic cloud and eventually to atom loss.

In order to optically repump the atoms in the fully spin-polarized Zeeman-substate, σ^{-} -polarized light has to be applied (see sec. 3.4). If the optical pumping light is not fully σ^{-} , but partially π - and σ^{+} -polarized, it will lead to an excitation of the atoms to energetically higher Zeeman-substates (see fig. 3.24). In the following the backrelaxation to lower Zeeman-substates leads to a heating of the atomic cloud and, if this energy is high enough for the atoms to leave the trap, it leads also to atom loss.



Figure 3.25: Temperature of the atoms after $t_{\text{demag}} = 2 \text{ s}$ of demagnetization cooling at a magnetic field of $B_z = 225 \text{ mG}$ and with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (orange) and with no optical pumping light (blue) for different angles of the $\lambda/4$ -waveplate which converts the formerly π -polarized light to σ^- -polarized light. The cooling performs best around a waveplate angle of 50°. In the previous measurements the waveplate angle was 43° .

Therefore, the polarization of the pumping light has to be optimized. In our experimental setup, a $\lambda/4$ -waveplate converts the formerly π -polarized light into σ^{-} -

polarized light before it is directed into the science chamber and on the atoms. In figure 3.25 the temperature of the atoms after $t_{\text{demag}} = 2 \text{ s}$ of demagnetization cooling is plotted over the angle of the $\lambda/4$ -waveplate. Here, demagnetization cooling was performed at a magnetic field of $B_z = 225 \text{ mG}$ and with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (orange). For comparison the case of no optical pumping light is also displayed (blue). A clear dependence of the cooling effect on the angle of the $\lambda/4$ -waveplate can be observed. The lowest temperature can be achieved at an angle of around 50°. It decreases by $\Delta T \approx 0.5 \,\mu\text{K}$ at this waveplate angle of 50° compared to the previous waveplate angle (43°). Therefore, further measurements are performed at this waveplate angle of 50°.

3.5.4.4 Transversal magnetic fields



Figure 3.26: Sketch of the experimental optical pumping setup. The σ^- -polarized pumping beam should be oriented parallel to the magnetic field. If the magnetic field is slightly tilted as in this figure, the atoms do not see fully σ^- -, but partially π - and σ^+ -polarized light. This leads to heating as depicted in figure 3.24.

As the magnetic field axis defines the quantization axis, the σ^- -polarized pumplight should be oriented parallel to it. If the magnetic field is tilted with respect to the optical pumping beam (see fig. 3.26), the atoms see no longer fully σ^- , but elliptically polarized light, which consists of partially π - and σ^+ -polarized light. This leads to heating as explained in subsection 3.5.4.3.

Consequently, the magnetic field values perpendicular to the optical pumping beam B_x and B_y have to be minimized. In order to do so demagnetization cooling at the optical pumping rate $\Gamma_{\rm op} = (20 \pm 3) 2\pi$ Hz and a magnetic field in z-direction of $B_z = 225$ mG is performed. After $t_{\rm demag} = 2$ s the cooling performance is compared.

In figure 3.27 the temperature after this demagnetization sequence is shown for different cage coil currents. The cage coil current is responsible for the generation of the magnetic field a) in x-direction and b) in y-direction. In x-direction a strong dependence of the temperature on the cage coil current can be observed. Demagnetization cooling seems to perform best at $I_{Bx} = -0.68$ A which is chosen as new standard value. In y-direction the temperature is mostly independent on the cage coil current in the measured range. Therefore we stay at the previous value of $I_{By} = -0.5$ A.

After the optimization of the magnetic field in z-direction, the polarization, and the transversal magnetic fields, the temperature after $t_{\text{demag}} = 2 \text{ s}$ of demagnetization cooling at a optical pumping rate of $\Gamma_{\text{op}} = 2\pi (20 \pm 3)$ Hz has changed from around $T_f =$ $17 \,\mu\text{K}$ to less than $T_{2s} = 15 \,\mu\text{K}$, whereas the atom number has stayed approximately



Figure 3.27: Temperature of the atoms after $t_{\text{demag}} = 2 \text{ s}$ of demagnetization cooling at a magnetic field of $B_z = 225 \text{ mG}$ and with an optical pumping rate of $\Gamma_{\text{op}} = (20\pm3) 2\pi \text{ Hz}$ (orange) and no optical pumping light (blue) dependent on a) the cage coil current which is responsible for the magnetic field in x-direction and b) in y-direction.

constant around $N_{2s} = (1.3 \pm 0.1) \, 10^6$. However, theoretically temperatures beyond $T_{2s} = 7 \,\mu\text{K}$ should be possible. One reason for this discrepancy between theory and experiment could be the confinement which we will address in the next subsection.

3.5.4.5 Confinement

The confinement in a trap ensures rethermalization of the atoms. The timescale on which this rethermalization happens is given by the trapping frequency in the respective direction.

The previous measurements have been done in the single ODT 1 (see sec. 3.5.1) with the following trapping frequencies $\bar{\omega} = (\omega_x \, \omega_y \, \omega_z)^{1/3} = (5.6 \cdot 783 \cdot 837)^{1/3} \, 2\pi \, \text{Hz}$. As the beam propagates in x-direction the trapping frequency in this direction is fairly low ($\omega_x = 5.6 \cdot 2\pi \, \text{Hz}$). We can compare this trapping frequency to the optical pumping rate Γ_{op} and the dipolar relaxation rate Γ_{dr} . The optical pumping rate was set in the previous measurements to $\Gamma_{\text{op}} \geq (20 \pm 3) \, 2\pi \, \text{Hz}$ and the dipolar relaxation rate was $\Gamma_{\text{dr}} = \beta_{\text{dr}} \, n \geq 2 \cdot 10^{-18} \cdot 5 \cdot 10^{17} \cdot 2\pi \, \text{Hz} = 2\pi \, \text{Hz}$. Consequently, the trapping frequency ω_x is close to the dipolar relaxation rate and even lower than the minimal optical pumping rates we applied. Having this in mind, we can assume that we are not limited by the optical pumping rate, but rather the loose confinement in x-direction.

Accordingly, we also used ODT 2 producing a crossed optical dipole trap (cODT, see sec. 3.5.1) in order to confine the atoms better in x-direction. The trapping frequencies in this crossed trap are $\bar{\omega} = (\omega_x \, \omega_y \, \omega_z)^{1/3} = (132 \cdot 392 \cdot 442)^{1/3} \, 2\pi \, \text{Hz}$ and thus in neither direction lower than the minimal optical pumping rate Γ_{op} . However, we perform a first evaporation step in the cODT before the previously mentioned trapping frequencies are reached. Therefore, we start at lower initial temperatures around $T_i = 7.4 \, \mu \text{K}$. As a consequence, an optimal magnetic field value has to be found for this temperature (compare sec. 3.5.4.2 and fig. 3.6).



Figure 3.28: Temperature of the atoms after $t_{\text{demag}} = 3 \text{ s}$ of demagnetization cooling with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (yellow points) and without any optical pumping light (blue points) dependent on a magnetic field which ranges from $B_z = 0$ to 385 mG. The blue curve shows the prediction of the equilibrium temperature in accordance with equation 3.56. This equilibrium temperature is obtained, if no optical pumping light is applied and therefore the cooling of the motional degree of freedom is only due to a transfer of energy from the kinetic to the spin energy reservoir.

In figure 3.28 the temperature T dependent on magnetic field B_z of the corresponding measurement is shown. Here, $t_{\text{demag}} = 3 \text{ s}$ of demagnetization cooling with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (yellow points) and without any optical pumping light (blue points) has been performed at magnetic field values ranging from $B_z = 0$ to 385 mG.

If no optical pumping light is applied and the atoms are initially polarized in the lowest magnetic sub-state, the kinetic and spin degrees of freedom of the atoms will thermalize via dipolar relaxation. The equilibrium temperature T_{eq} after a long time depends on the magnetic field B_z , the initial temperature T_i and the total angular momentum J (for dysprosium J = 8) and can be calculated using [2]

$$3 k_{\rm B} T_{\rm i} = 3 k_{\rm B} T_{\rm eq} + g_J \mu_{\rm B} B \frac{\sum_{j=0}^{2J} j \exp\left(-j\frac{g_J \mu_{\rm B} B}{k_{\rm B} T_{\rm eq}}\right)}{\sum_{j=0}^{2J} \exp\left(-j\frac{g_J \mu_{\rm B} B}{k_{\rm B} T_{\rm eq}}\right)}.$$
(3.56)

Hence, the temperature should be reduced after a sufficiently long hold time, even if no optical pumping light is applied (see e.g. figure 3.28).

However, in figure 3.28 a discrepancy between the data (blue points) and the theoretical equilibrium temperature (blue curve) can be observed. Experimentally a clear minimum in temperature can not be observed in contrast to the theoretical equilibrium temperature which has a minimum at $B_z = 20 \text{ mG}$. On the contrary, experimentally a heating can be observed in both cases with and without applied



Figure 3.29: a) Atom number and b) phase space density \mathcal{D} of the atoms after $t_{\text{demag}} = 3 \text{ s}$ of demagnetization cooling with an optical pumping rate of $\Gamma_{\text{op}} = (20 \pm 3) 2\pi \text{ Hz}$ (yellow points) and without any optical pumping light (blue points) dependent on a magnetic field which ranges from $B_z = 0$ to 385 mG.

optical pumping light at magnetic field values $B_z < 100 \text{ mG}$. Up to now, we could not explain this heating at low magnetic field values, which we observe not only in the cODT, but also in ODT 1 (see fig. 3.22).

The experimental temperature evolution with and without optical pumping light is very similar. Even though a lower temperature can be reached if optical pumping light is applied and thus, an effect of demagnetization cooling can be observed, there is not a clear improvement compared to the case of no optical pumping light. Furthermore, in figure 3.29 it can be seen that the loss in atom number is even higher with optical pumping light. Consequently, the phase space density is lower in the case of optical pumping light. In this measurement heating and atom loss mechanisms, e.g. due to optical pumping photons [55], outperform the effect of dipolar relaxation collisions which lead to a cooling. A possible reason is the fact that we start already at a lower initial temperature than in the previous measurement in ODT 1 (see sec. 3.5.4.1 to 3.5.4.4). However, the frequency of dipolar relaxation collisions diminishs at lower temperatures, as lower temperatures are accompagnied by lower velocities and thus less collisions happen in a certain time interval (see sec. 3.3.1). Furthermore, evaporative cooling makes a major contribution to the overall cooling effect and hence demagnetization cooling can not be evaluated independently this way. In order to circumvent those problems demagnetization cooling should be performed in the crossed optical dipole trap without applying a first evaporation step (see sec. 3.5.1). This way

the initial temperature should be higher which is favorable for demagnetization cooling as explained before.



Figure 3.30: a) Temperature, b) atom number and d) phase space density after $t_{\text{demag}} = 3 \text{ s}$ of demagnetization cooling with two different opical pumping intensities (blue points: $I_{\text{op}} = (12 \pm 2) \frac{\mu W}{\text{mm}^2}$, yellow points: $I_{\text{op}} = (35 \pm 6) \frac{\mu W}{\text{mm}^2}$) and at a detuning ranging from $\Delta = -100$ to -35 MHz with a resolution of 5Γ . For comparison the respective temperature and atom number without optical pumping light is plotted as a blue, dashed line. c) Optical pumping rate for the same detuning range and intensities (see eq. 3.37).

However, heating and atom loss due to optical pumping light can be analyzed qualitatively in this trap configuration. Since the optical pumping rate $\Gamma_{\rm op}$ depends not only on the applied intensity $I_{\rm op}$, but also on the detuning Δ from resonance (3.37), different optical pumping rates can be probed, if one stays at a certain intensity and changes the detuning. As depicted in figure 3.30c, we could probe an optical pumping rate ranging from $\Gamma_{\rm op} = 202\pi$ Hz up to over $\Gamma_{\rm op} = 5002\pi$ Hz by applying optical pumping light once with an intensity of $I_{\rm op} = (12\pm 2) \frac{\mu W}{mm^2}$ and second $I_{\rm op} = (35\pm 6) \frac{\mu W}{mm^2}$ and by tuning the frequency of the optical pumping light from $\Delta = -100$ to -35Γ . In figure 3.30a, b and d temperature, atom number and phase space density are plotted for those two different optical pumping intensities (blue points: $I_{\rm op} = (35\pm 6) \frac{\mu W}{mm^2}$) and the mentioned detuning range. For comparison, we show the case of no optical pumping light as a dashed blue line. The discussed data has been taken after $t_{\rm demag} = 3$ s wait time, during which a magnetic field of

 $B_z = 137 \,\mathrm{mG}$ was applied.

In figure 3.30 it becomes apparent that the atom number and the phase space density is highest in the case of no optical pumping light. Solely, the temperature shows an improvement when adding optical pumping light. For $\Gamma_{\rm op} < 100 \cdot 2\pi$ Hz the temperature is lower with optical pumping light than without. Hence, an effect of demagnetization cooling in terms of temperature can still be observed. This tendency has already become apparent in the previous measurement in the confined trap (see fig. 3.28 and 3.29). Further, the higher the optical pumping rate, the lower the atom number and phase space density. For $\Gamma_{\rm op} > 100 \cdot 2\pi$ Hz heating due to the optical pumping light outperforms cooling due to dipolar relaxation collisions and the temperature is higher with applied optical pumping light than without it. Further, for $\Gamma_{\rm op} > 60 \cdot 2\pi$ Hz the higher the optical pumping rate, the higher the temperature.

As it can be seen in figure 3.30c and d, at two different detunings and the same optical pumping rate ($\Gamma_{\rm op} = 100 \cdot 2\pi \, \text{Hz}$), the phase space density is approximately the same $(\mathcal{D} = 10^{-3})$. Hence, in this first prove of principle measurement a dependence of the performance of demagnetization cooling on the detuning can not be observed. However, systematic measurements have to be taken in order to confirm or disprove this first guess. Furthermore, a larger frequency range has to be probed. In subsection 3.4.4, we have seen that suppression of light assisted collisions can be achieved at the detuning which corresponds to the node of the ground state wavefunction. In Dysprosium the corresponding detuning is $|\Delta| > 0.5 \cdot 2\pi \text{ GHz} \approx 5\,000\,\Gamma$ (see fig. 3.12). However, the maximal detuning in the measurements which have been taken in the scope of this thesis is $|\Delta| = 10 \cdot 2\pi \,\mathrm{MHz} \approx 100 \,\Gamma$ and thus at a much lower value. In order to further optimize demagnetization cooling, systematic measurements of its performance with constant optical pumping rate and a detuning which ranges up to $|\Delta| = 10.2\pi \,\text{GHz} \approx 100\,000\,\Gamma$ have to be taken. In order to localize the node position of the ground state wavefunction blue detuning is favorable, as in contrast to red detuning the binary loss rate is not modulated by bound state resonances (see fig. 3.11).

3.5.4.6 Efficiency

The most important figure of merit of any cooling technique is the efficiency. It is defined as

$$\chi = -\frac{\log\left(\mathcal{D}_{\rm f}/\mathcal{D}_{\rm i}\right)}{\log\left(N_{\rm f}/N_{\rm i}\right)} \tag{3.57}$$

and relates the gain in phase-space density to the loss in atom number. In a double logarithmic plot of the phase-space density over the number of atoms, the efficiency corresponds to the negativ slope of the curve which describes the evolution of the phase-space density and the atom number.

In figure 3.31 phase space density and atom number evolution of two exemplary measurements are plotted in a double-logarithmic plot. One measurement shows the case of no optical pumping light (green points) and the other measurement shows the evolution with an optical pumping rate $\Gamma_{\rm op} = (16.8 \pm 2.2) 2\pi \,\mathrm{Hz}$ (blue points).



Figure 3.31: Double logarithmic plot of phase space density over atom number evolution. Both data sets show the temporal evolution during to t = 2 s in ODT 1 at a magnetic field $B_z = 385 \text{ mG}$ and at a detuning of $\Delta = 20 \Gamma$. For comparison, there is one data set without and one with applied optical pumping light (green and $\Gamma_{\text{op}} = (16.8 \pm 2.2) 2\pi \text{ Hz}$, blue points respectively).

Both have been performed in ODT 1, at a magnetic field $B_z = 385 \,\mathrm{mG}$, a starting temperature of $T_0 \approx 38 \,\mu\mathrm{K}$ and at a detuning of $\Delta = 20 \,\Gamma$.

The gain in phase space density is clearly higher for $\Gamma_{\rm op} = (16.8 \pm 2.2) 2\pi$ Hz than for $\Gamma_{\rm op} = 0$. However, during the first second both cases show a similar behaviour. The slope of the case with optical pumping light is slightly higher, but the loss in atom number is approximately the same. A possible explanation is that the cooling process is dominated by evaporative cooling during the first second. Therefore, the measurements presented in subsection 3.5.4.1 to 3.5.4.4 are taken after a wait time of $t_{\rm evap} = 1.5$ s in order to ensure that the observed effect is dominated by demagnetization cooling. In figure 3.31 we can see that after approximately one second the atom number and the gain in phase space density saturates in the case of $\Gamma_{\rm op} = (16.8 \pm 2.2) 2\pi$ Hz the loss in atom number decreases after approximately one second, but a gain in phase space density can still be observed. Here, the effect of demagnetization cooling becomes visible. A fit to the data gives the efficiencies listed in table 3.1.

In evaporative cooling efficiencies of $\chi \approx 4$ can only be reached with substantial effort [69]. Hence, demagnetization cooling clearly outperforms evaporative cooling

Efficiency χ	$t = 01 \mathrm{s}$	$t = 12 \mathrm{s}$
$\Gamma_{\rm op} = 0$	1.22 ± 0.24	-
$\Gamma_{\rm op} = (16.8 \pm 2.2) 2\pi {\rm Hz}$	4.46 ± 0.54	7.95 ± 1.58

Table 3.1: Efficiencies obtained by a fit to the data presented in fig. 3.31 in the time interval indicated in the first line of the table.

already in this first dysprosium proof of principle experiment.

However, by a post-cooling step, which I mentioned at the end of section 3.5.1, we can reach a maximal cooling efficiency of $\chi = 6.7$ [5]. Hence, this post-cooling has a similar cooling efficiency as our first proof of principle demagnetization cooling experiment.

CHAPTER 4

Broad Feshbach resonances

Feshbach resonances are a very powerful tool in atomic and molecular physics in order to control the interactions between the atoms [10]. Adjusting the interaction strength in an ultracold Fermi gas of atoms led to the direct observation of a molecular Bose-Einstein condensate and the cross-over from a Bardeen-Cooper-Schrieffer (BCS)-type superfluid to a Bose-Einstein condensate (BEC) [70]. Furthermore, Efimov three-body physics and related few-body phenomena could be observed by tuning the interactions in the resonant and universal regime [71].

In this chapter, I will briefly explain the elementary properties of a Feshbach resonance. Following this, I will report on our measurement of two broad Feshbach resonances in Dysprosium and their characterization by atom loss spectroscopy and magnetic field modulation spectroscopy. Furthermore, we analyzed the loss dynamics on resonance using a universal model and the temperature dependence of the atom loss, which I will present in the last section.

4.1 Feshbach resonance

The physical origin and theoretical details of Feshbach resonances are very well explained and dicussed in [10]. Here, I will focus on the elementary properties of a Feshbach resonance.

Magnetic Feshbach resonances can be understood in a simple picture of two colliding atoms, which experience a molecular potential (see fig. 4.1). The molecular potential can support discrete rovibrational bound molecular states. At finite magnetic field the potential splits into Zeeman-substates. In order to understand the basic principle of a Feshbach resonance, we can concentrate on two of those molecular potential curves (see fig. 4.2). For collision processes with small scattering energy $E_{\rm sc}$, only one of those potentials is energetically accessible. This potential represents the energetically open channel, which is also called background potential $V_{\rm bg}(R)$. For large internuclear distances R, it connects asymptotically to two free atoms. The other potential which is energetically inaccessible for the atoms is referred to as the closed channel $V_{\rm c}$. It supports discrete rovibrational bound molecular states, which can eventually be located



Figure 4.1: Exemplary explanation of the emergence of different potential curves. a) Molecular potential curve which is experienced by two colliding atoms (b) at zero magnetic field. The potential supports several rovibrational states. c) At finite magnetic field $B \neq 0$ the Zeeman-substates at large interatomic separations $R \rightarrow \infty$ split as well as the corresponding potentials.

near the threshold of the open channel.



Figure 4.2: Basic two-channel model for a Feshbach-resonance. a) Molecular potential curve of the closed $V_{\rm c}(R)$ and open channel $V_{\rm bg}$ with respective magnetic moments $\mu_{\rm c}$ and $\mu_{\rm bg}$. The open channel is in contrast to the closed channel energetically accessible for two atoms colliding with energy $E_{\rm sc}$. b) A resonance occurs when two atoms colliding with the energy $E_{\rm sc}$ couple resonantly to a bound molecular state in the closed channel. The coupling is symbolized by a green arrow.

A Feshbach resonance occurs when the energy one of those bound molecular states is equal to the energy of the scattering state in the open channel. Then even a weak coupling between closed and open channel can lead to a strong mixing between the scattering state and the bound molecular state. The energy difference between closed and open channel can be tuned by a magnetic field, if the magnetic moments of closed and open channel are different. We refer to this type of Feshbach resonance as a magnetic Feshbach resonance.



Figure 4.3: Properties of a Feshbach resonance. Graph taken from [10]. a) Comparison of contact (grey sphere) and dipolar (green-red magnet) interaction, b) scattering length a_s and c) molecular state energy E close to a magnetically tuned Feshbach resonance. The binding energy $E_{\rm b} = -E$ is defined to be positive. In the inset E is shown in the universal regime close to the resonance at large and positive scattering length a_s .

The energy E of a weakly bound molecular state near a magnetic Feshbach resonance is shown in figure 4.3c relative to the threshold of two free atoms with zero kinetic energy. It can be described in two different regimes. At magnetic field values away from resonance, where the energy of the bound molecular state is lower than the threshold of the open channel, the molecular state energy E increases linearly with the magnetic field, where the slope is given by the difference in magnetic moments of the open and closed channel $\delta \mu = |\mu_{\rm bg} - \mu_{\rm c}|$. Closer to the resonance, the coupling of the two channels mixes in contributions of the open channel and bends the molecular state: $E_{\rm b} \propto (B - B_0)^2$.

In the vicinity of the resonance the scattering of two free atoms is greatly enhanced due to the presence of the near-threashold bound state. It can be drawn an analogy to resonances in other systems, which are characterized by a phase shift of π across the resonance. The same phase shift in the wave-function appears across a Feshbach resonance and enhances the scattering. Without going into further details of scattering theory [10], the tuning of interactions at a Feshbach resonance can be described by the variation of the s-wave scattering length a_s as a function of the magnetic field B [72],

$$a_s(B) = a_{\rm bg} \left(1 - \frac{\Delta_B}{B - B_0} \right). \tag{4.1}$$

This relation is visualized in figure 4.3b. As already explained in chapter 3.4.1, the scattering length a_s characterizes the interaction strength. The scattering length which is associated with the background potential $V_{\rm bg}(R)$ and therefore called background scattering length $a_{\rm bg}$ represents its value far away from resonance. It is related to the energy of the last-bound vibrational level. The resonance is located at the magnetic field B_0 where the scattering length diverges $(a \to \pm \infty)$. The resonance width Δ_B depends on the coupling strength between the open and closed channel. The zero crossing of the scattering length is located at the magnetic field $B = B_0 + \Delta_B$.

The binding energy close to the resonance, where the scattering length diverges and the two channels are strongly coupled, can also be expressed in terms of the scattering length:

$$E_{\rm b} = \frac{\hbar^2}{2\,\mu_{\rm red}\,a_s^2} \propto (B - B_0)^2\,,\tag{4.2}$$

where $\mu_{\rm red}$ is the reduced mass of the atom pair. The bend shown in the inset of fig. 4.3c reflects that the binding energy depends quadratically on the magnetic detuning from resonance $B - B_0$. Since in this region the details of the interaction become irrelevant, it is called universal regime [10, 73].

As already outlined in chapter 2, in Dysprosium both the contact interaction and the magnetic dipole-dipole interaction play an important role. Furthermore, we introduced the dipolar length D in order to characterize the dipolar interactions (see eq. 3.3 and [6, 40]). Comparing the background scattering length with the dipolar length, we saw that in the regime of the background scattering length, and hence away from resonance, the interaction is dominated by dipolar interactions (see eq. 3.31, 3.32, and 3.3). However, tuning the scattering length with a Feshbach resonance allows one to choose a contact- or dipolar-interaction-dominated regime (see fig. 4.3a).

4.1.1 Strength of a Feshbach resonance

In order to characterize the strength of a resonance, a dimensionless resonance strength parameter

$$s_{\rm res} = \frac{a_{\rm bg} \,\delta\mu \,\Delta_B}{\bar{a} \,\bar{E}} \tag{4.3}$$

can be defined [10]. Here, $\bar{a} = \frac{4\pi}{\Gamma^2(1/2)} \left(\frac{5}{4}\right)^{1/4} R_{\rm B} \approx 1.01\,082\,R_{\rm B}$ is the mean scattering length [74], $\bar{E} = \frac{\hbar^2}{2\,\mu_{\rm red}\bar{a}^2}$ is a corresponding energy scale, and $R_{\rm B} = \left(\frac{\mu_{\rm red}\,C_6}{10\,\hbar^2}\right)^{1/4}$ is the

van der Waals length scale (see sec. 3.4).

There are two limiting cases. First, if $s_{\rm res} \gg 1$, the resonances are open-channeldominated. The bound state is universal over a large fraction of the resonance width Δ_B and the binding energy $E_{\rm B}$ can be well approximated by equation 4.2.

Second, if $s_{\rm res} \ll 1$, the resonance is closed-channel-dominated. The bound state is universal only over a small fraction of the resonance width Δ_B near $B = B_0$ and the binding energy $E_{\rm B}$ needs to be modeled by a coupled-channel description [10].

The interaction strength can be described in both cases by the scattering length given by equation 4.1.

4.2 Universal loss dynamics

Near a Feshbach resonance interactions are strongly enhanced, leading to an increased collision rate. This increase in collision rate leads to an increase of atom losses by two main mechanisms: three-body recombinations and 2-body evaporation.

Two- and three-body collision losses can be quantified by rate equations for atom number and temperature. Those rate equations will be discussed in the following subsections on the basis of [75].

4.2.1 Rate equation for atom number

4.2.1.1 Three-body recombination loss

In a three-body collision, two colliding atoms form a dimer and the binding energy is transfered into kinetic energy of the dimer and the third atom. Usually the kinetic energy is larger than the trap depth, which allows all three atoms to leave the trap. Since three-body recombination processes occur more often at high densities and thus at the center of the trap were the temperature is the lowest, atoms with small potential energy are lost. Therefore, three-body losses lead to a net heating of the atomic cloud [75].

The loss rate of atoms due to three-body recombination is given by [75]

$$\frac{\mathrm{d}N}{\mathrm{d}t} = -3\int \frac{L_3 n^3(r)}{3} \,\mathrm{d}^3 r = -L_3 \,\langle n^2 \rangle \, N \,, \tag{4.4}$$

where $L_3 n^3(r)/3$ is the three-body recombination rate and the factor of 3 in front of the integral reflects the fact that all three atoms are lost per recombination event [75].

At finite s-wave scattering length $(a_s k \ll 1)$ in the universal regime, L_3 can be expressed in the form [76]

$$L_3 = C(a_s) \frac{\hbar}{m} a_s^4, \qquad (4.5)$$

where a general a_s^4 -scaling has been separated from an additional dimensionless dependence $C(a_s)$, which is given in [76] and for three interacting bosonic dipoles in [77]. For a Boltzmann distributed gas in a harmonic trap:

$$\langle n^2 \rangle = \left(\frac{m\bar{\omega}^2}{2\sqrt{3}\pi k_{\rm B}}\right)^3 \frac{N^2}{T^3} \tag{4.6}$$

and hence the rate equation for three-body losses of bosonic dipoles at finite s-wave scattering length a_s in the universal regime reads

$$\frac{\mathrm{d}N}{\mathrm{d}t} = -\gamma_3^{a_s} \frac{N^3}{T^3} \,, \tag{4.7}$$

where

$$\gamma_3^{a_s} = C(a_s) \frac{\hbar}{m} a_s^4 \left(\frac{m\bar{\omega}^2}{2\sqrt{3}\pi k_{\rm B}} \right)^3 \,. \tag{4.8}$$

 L_3 is proportional to a_s^4 , but when approaching the resonance, the scattering length diverges and becomes larger than any other length scale $(a_s k \gg 1)$, where k is the relative wavenumber of the colliding particles). In this so-called unitary regime we may express a_s by 1/k, where $k \propto \sqrt{T}$. Therefore, a naiv guess would be that L_3 shows a T^{-2} behaviour. Indeed, [75] showed that in the unitary regime L_3 can be approximated by

$$L_3 \approx \frac{\hbar^5}{m^3} 36\sqrt{3}\pi^2 \frac{1 - e^{-4\eta_*}}{(k_{\rm B}T)^2} = \frac{\lambda_3}{T^2}, \qquad (4.9)$$

where η_* is the inelasticity parameter which characterizes the strength of the short range inelastic processes, i.e. the efficiency with which three atoms in contact recombine to a dimer and a free atom, and λ_3 is a temperature-independent constant. Finally the rate equation reads

$$\frac{\mathrm{d}N}{\mathrm{d}t} = -\gamma_3^{\mathrm{u}} \frac{N^3}{T^5} \,, \tag{4.10}$$

where

$$\gamma_3^{\rm u} = \lambda_3 \left(\frac{m\bar{\omega}^2}{2\sqrt{3}\pi k_{\rm B}}\right)^3. \tag{4.11}$$

As the change of atom number in time scales with T^{-5} in temperature, three-body atom loss in the unitary regime depends more strongly on temperature than for finite scattering length. Finally, the dependence of atom loss on temperature and the inelasticity parameter can be studied in the unitary regime.

4.2.1.2 Two-body losses

If two atoms collide with the initial momenta p_1 and p_2 , they emerge from the collision with the momenta p_3 and p_4 . If one of them acquires a momentum $|p_3| \ge \sqrt{2mU}$, where U is the potential depth, it will leave the trap. This mechanism is called evaporation.

The loss of atoms due to evaporation can be described by the following rate-equation for the atom number [75, 78]:

$$\frac{\mathrm{d}N}{\mathrm{d}t} = -n_0^2 \,\sigma \,\bar{v} \,e^{-\eta} \,V_{\mathrm{ev}} \equiv -\Gamma_{\mathrm{ev}} N\,, \qquad (4.12)$$

Here, $n_0 = N/V_e$ is the peak density and V_e is the effective volume of the sample. In a harmonic trap V_e is given by $V_e = \left(\frac{2\pi k_{\rm B}T}{m\bar{\omega}}\right)^{\frac{3}{2}}$ [75]. $\bar{v} = \sqrt{\frac{8k_{\rm B}T}{\pi m}}$ is the mean velocity, $\eta = U/k_{\rm B}T$ is the ratio of potential depth and thermal energy, and $V_{\rm ev}$ is the effective volume for elastic collisions which lead to evaporation. Its expression can be found in [78]. The low-energy scattering cross section is given by

$$\sigma(k) = \frac{8\pi}{k^2 + a_s^{-2}},\tag{4.13}$$

where k is the relative momentum of the colliding partners. For $a_{\rm s}k \ll 1$ it reduces to

$$\sigma_{a_s} = 8\pi a_s^2 \,. \tag{4.14}$$

For a finite scattering length and additionally a finite dipolar length $D \neq 0$ the scattering cross-section becomes [40]

$$\sigma_{a_s,D} = 8\pi \left(a_s^2 + \frac{4}{45} D^2 \right) \,. \tag{4.15}$$

At unitarity $(a_{\rm s}k \gg 1)$ the scattering cross-section reduces to

$$\sigma = 8\pi/k^2 \tag{4.16}$$

Note that $a_s \stackrel{\text{unitarity}}{\to} 1/k$ applies here, too. The relative momentum of the colliding partners is given by $|\mathbf{k}| = k = \sqrt{\frac{mU}{2\hbar^2}}$ [75]. This results in the following collisional cross-section at unitarity

$$\sigma_U = \frac{16\pi\hbar^2}{m\,U} \tag{4.17}$$

which is independent of the energy of the colliding partners.

Combining in the unitary limit three-body (eq. (4.10) and (4.11)) and two-body

loss, the total atom number loss rate equation can be expressed as follows

$$\frac{\mathrm{d}N}{\mathrm{d}t} = -\gamma_3^{\mathrm{u}} \frac{N^3}{T^5} - \gamma_2^U e^{-\eta} \frac{V_{\mathrm{ev}}}{V_{\mathrm{e}}} \frac{N^2}{T} , \qquad (4.18)$$

where

$$\gamma_2^U = \frac{16}{\pi} \frac{\hbar^2 \bar{\omega}^3}{k_{\rm B} U} \,. \tag{4.19}$$

At finite scattering length and finite dipolar length taking into account three-body recombination loss (see eq. (4.7) and (4.8)) and evaporative loss, the total atom number loss rate equation can be expressed as

$$\frac{\mathrm{d}N}{\mathrm{d}t} = -\gamma_3^{a_s} \frac{N^3}{T^3} - \gamma_2^{a_s, D} e^{-\eta} \frac{V_{\mathrm{ev}}}{V_{\mathrm{e}}} \frac{N^2}{T} , \qquad (4.20)$$

where

$$\gamma_2^{a_s,D} = \frac{8}{\pi} \frac{m \,\bar{\omega}^3}{k_{\rm B}} \, \left(a_s^2 + \frac{4}{45} \, D^2 \right) \,. \tag{4.21}$$

4.2.2 Rate equation for temperature

4.2.2.1 Three-body recombination heating

Reference [75] points out that the lost energy per lost atom in a three-body recombination process at unitarity is

$$\frac{\dot{E}}{\dot{N}} \stackrel{3b}{=} \frac{4}{3} k_{\rm B} T \,, \tag{4.22}$$

whereas at finite scattering length it is higher

$$\frac{\dot{E}}{\dot{N}} \stackrel{3b}{=} 2 k_{\rm B} T \,. \tag{4.23}$$

The mean energy per atom in the harmonic trap is $3 k_{\rm B}T$. Therefore, in each threebody recombination event a loss of an atom is associated at unitarity with an excess of $(3 - 4/3) k_{\rm B}T = 5/3 k_{\rm B}T$ and at finite scattering length [79] with an excess of $(3 - 2) k_{\rm B}T = k_{\rm B}T$ of energy that remains in the sample.

The equations (4.22) and (4.23) combined with the fact that for a non-interacting gas $E = 3 N k_{\rm B} T$ and eq. (4.10) at unitarity and (4.7) at finite scattering length respectively can be transformed into the rate equation for the rise of temperature per

lost atom. At unitarity, it is given by [75]:

$$\frac{\mathrm{d}T}{\mathrm{d}t} = \frac{5}{3} \frac{T}{3} \gamma_3^{\mathrm{u}} \frac{N^2}{T^5} \,, \tag{4.24}$$

whereas for finite scattering length, we obtain

$$\frac{\mathrm{d}T}{\mathrm{d}t} = \frac{T}{3}\gamma_3^{a_s}\frac{N^2}{T^3}\,.$$
(4.25)

4.2.2.2 Evaporative cooling

In two-body collisions, mostly atoms with a high potential energy leave the trap. This mechanism is used in forced evaporative cooling, where the trap depth is decreased in order to remove high-energy atoms from the thermal cloud, which leads to a net cooling effect after thermalization [80].

In reference [75] the temperature rate equation is derived as follows

$$\frac{\mathrm{d}T}{\mathrm{d}t} = -\Gamma_{\mathrm{ev}} \left(\eta + \tilde{\kappa} - 3\right) \frac{T}{3} \,, \tag{4.26}$$

where in a harmonic trap [78]

$$\tilde{\kappa} = 1 - \frac{P(5,\eta)}{P(3,\eta)} \frac{V_{\rm e}}{V_{\rm ev}}$$

$$\tag{4.27}$$

and $P(a,\eta)$ is the incomplete Gamma function. $\Gamma_{\rm ev}$ depends on the scattering cross section σ (see eq. (4.12)). If the respective scattering cross-section is plugged in, the temperature rate equation can be expressed either in the unitary regime or for finite scattering and dipolar length. Presented in a similar manner as in the previous section the temperature rate equation reads

$$\frac{\mathrm{d}T}{\mathrm{d}t} = -\gamma_2 \, e^{-\eta} \, \frac{V_{\mathrm{ev}}}{V_{\mathrm{e}}} \, (\eta + \tilde{\kappa} - 3) \frac{N}{T} \frac{T}{3} \,, \tag{4.28}$$

where

$$\gamma_2 = \begin{cases} \gamma_2^{a_s, D} & \text{if } a_s k \ll 1\\ \gamma_2^U & \text{if } a_s k \gg 1 \end{cases}.$$

$$(4.29)$$

Combining the process of recombination heating and evaporative cooling at unitarity (eq. 4.24 and 4.28) the temperature rate equation reads [75]:

$$\frac{\mathrm{d}T}{\mathrm{d}t} = \frac{T}{3} \left(\frac{5}{3} \gamma_3^{\mathrm{u}} \frac{N^2}{T^5} - \gamma_2^U e^{-\eta} \frac{V_{\mathrm{ev}}}{V_{\mathrm{e}}} (\eta + \tilde{\kappa} - 3) \frac{N}{T} \right) \,, \tag{4.30}$$

whereas at finite scattering and dipolar length we get (eq. 4.25 and 4.28)

$$\frac{\mathrm{d}T}{\mathrm{d}t} = \frac{T}{3} \left(\gamma_3^{a_s} \frac{N^2}{T^3} - \gamma_2^{a_s,D} e^{-\eta} \frac{V_{\mathrm{ev}}}{V_{\mathrm{e}}} (\eta + \tilde{\kappa} - 3) \frac{N}{T} \right) \,. \tag{4.31}$$

The atom loss dynamics are described at unitarity by the set of coupled rate equations (4.18) and (4.30) and at finite scattering and dipolar length by the set of coupled rate equations (4.20) and (4.31).
4.3 Characterization of two broad Feshbach resonances in dysprosium

In dysprosium, the complex electronic structure leads to a large magnetic dipole moment (see chapter 2), which in turn gives rise to a complex molecular spectrum. Figure 4.4 shows an example of the potential energy curves of 164 Dy + 164 Dy collisions which dissociate to the six energetically lowest Zeeman states. The high number of potentials indicates the high number of possible Feshbach resonances [81].



Figure 4.4: Longrange potential energy curves of collisions of two ¹⁶⁴Dy atoms in a magnetic field B = 50 G as a function of internuclear separation. Image taken from [81]. The red dashed line at zero energy indicates the energy of the open channel. A resonance occurs when a bound state has the same energy as the open channel. Here m_j is the projection of the sum of the total angular momentum of the two colliding atoms $j = J_1 + J_2$ on the quantization axis The graph shows the channels with even $l \leq 10$. The value of l is indicated for the $m_j = -16$ curves.

In alkali atoms the coupling between open and closed channel is hyperfine induced [10]. However, the bosonic dysprosium isotopes have a nuclear spin I = 0. The complex electronic structure of dysprosium leads not only to a large magnetic dipole moment, but also to a strong anisotropy of the van der Waals interaction. Similar to Erbium, this gives rise to a very dense spectrum of narrow Feshbach resonances. The nearest neighbour distribution can be attributed to a chaotic distribution commonly observed in multi-level systems with strong inter-level interactions [28, 29].

Further, we could observe two broad Feshbach resonances in the fully spin-polarized ground state of ¹⁶⁴Dy, J = 8, M = -8. In this section, I will present our respective inelastic loss spectroscopy measurements up to a magnetic field B = 600 G, which reveals loss features with widths of several Gauss. Besides, we did magnetic field modulation spectroscopy in order to measure the binding energy $E_{\rm b}$, which I will present in the second subsection.

4.3.1 Atom loss spectropsopy

The most frequently observed signature of Feshbach resonances are resonant losses. These losses may be induced by three-body collisions and occur as internal energy of the colliding atoms is released into kinetic energy. This happens when the colliding atoms end up in a lower internal state or when molecules are formed. The frequency of inelastic collisions and thus inelastic loss is strongly enhanced close to Feshbach resonances. This is due to a strong coupling of the Feshbach bound states to inelastic outgoing channels. Therefore, near a Feshbach resonance, the probability that the atoms gain kinetic energy due to inelastic collisions is strongly enhanced [10].

Consequently, atom loss spectroscopy permits to determine the position and width of Feshbach resonances.

4.3.1.1 Experimental procedure

We prepared the atomic cloud in the crossed optical dipole trap as explained in subsection 3.5.1 and [29, 32]. By evaporation at a magnetic field close to $1 \,\mathrm{G}$ we reached temperatures of $T \approx 0.6 \,\mu\text{K}$. Then, we recompressed the trap by increasing the power of the beams by a factor of approximately 2 in 100 ms. This was done for two reasons: First, we suppressed additional atom loss due to residual evaporation, which may be induced by increased elastic collisions. Second, we created a strong bias magnetic field with a pair of coils in Helmholtz configuration (see fig. 3.14), which created a gradient field weakening the trap confinement. Recompression minimizes this weakening of the trap confinement. In order to calibrate the magnetic field, we performed radio-frequency (RF) spectroscopy between the two lowest Zeemansubstates of ¹⁶⁴Dy. For this purpose, we applied electromagnetic radiation in the radio-frequency range to a sample of atoms in the lowest Zeeman substate at a certain unkown magnetic field value. If the applied frequency matches the Larmor frequency, we drive the excitation between the two lowest Zeeman-substates and observe this resonance as an increase in atom loss. As a consequence, the energy difference between the two lowest Zeeman-substates is determined by the resonance. Furthermore, we know the Zeeman-splitting as a function of the magnetic field. Therefore, we could deduce the applied magnetic field value by this measurement.

The magnetic field was ramped up in 15 ms to the target value. If we turned on the magnetic field, we observed a residual heating. This heating was partially due to an unintentional change of the magnetic field direction when we ramped to high magnetic fields. When the magnetic field changes its sign, the cloud is in a magnetic field B = 0, which allows the atoms to slightly depolarize. When the field is increased inelastic dipolar relaxation collisions cause a heating of the cloud. Probably, the residual magnetic field gradient contributed to this heating as well. Finally, we obtained typically samples which contained 10^5 atoms in their lowest Zeeman sublevel M = -8 at a temperature of $T = 2.4 \,\mu$ K.

For atom loss spectroscopy, the atoms are now held in the optical dipole trap at a constant magnetic field value during a certain wait time. During this time the atoms undergo elastic and inelastic collisions. The latter leads to atom loss from the trap as explained at the beginning of this subsection 4.3.1.

4.3.1.2 Feshbach resonance spectrum



Figure 4.5: Feshbach resonance spectrum of ¹⁶⁴Dy in J = 8, M = -8 with 100 mG resolution by atom loss spectroscopy at a temperature of 2.4 μ K. The atom number is normalized with respect to a reference atom number $N_{\rm ref}$ which was taken at low field every 30 images. [6]

Feshbach resonance spectrum of ¹⁶⁴Dy at $T = 2.4 \,\mu\text{K}$ In figure 4.5 the final atom number is plotted for magnetic field values ranging from B = 60 to 580 G with a resolution of 100 mG. In the corresponding measurement the wait time in the dipole trap at the respective constant magnetic field value was $t = 2 \,\text{s}$, the initial temperature of the atoms was $2.4 \,\mu\text{K}$ and the starting atom number was $N_0 = 1.8 \cdot 10^5$.

We can observe a very high energy-level density as it was observed in [29, 82]. However, several broader features form an irregular pattern on top of the narrow resonances. Two especially broad features appear in $^{164}\mathrm{Dy}$ around $80\,\mathrm{G}$ and around $180\,\mathrm{G}.$

In the next sections, we focus on investigation techniques to determine typical parameters of Feshbach resonances. Magnetic field modulation spectroscopy is a tool to map weakly bound molecular states. By this method, we show in the next section that the broad loss features can be attributed to a single molecular state and hence the losses can be analyzed in terms of universal scattering behaviour.

4.3.2 Magnetic field modulation spectroscopy

In order to address the question whether those broader features correspond to single resonances with a large magnetic field width, we performed magnetic field modulation spectroscopy [83, 84].



Figure 4.6: Scheme of molecular association by magnetic field amplitude modulation. If the oscillation of the magnetic field amplitude is on resonance with the binding energy of a dimer, the respective molecules form.

In magnetic field modulation spectroscopy molecules are associated by modulating sinusoidally the magnetic field amplitude. The energy of the oscillating magnetic field stimulates an atom pair to lose energy emitting a photon with very low frequency and thereby decay to a bound molecular state at a lower energy. However, molecules can only be associated by this technique, if the frequency of the oscillating magnetic field is near the binding energy of the dimer [83]. Either the molecules are no longer trapped or they induce increased three-body losses ($Dy_2 + Dy \rightarrow 3Dy + energy$). Therefore, a decay in atom number indicates that molecules have been formed and hence, the energy of the oscillating magnetic field is close to the binding energy of the respective molecules.

Consequently, we can measure the energy of the bound states as a function of the magnetic field by this technique of molecular association by magnetic field modulation.

We are especially interested in the broader features around 80 G and 180 G. We focus for magnetic field modulation spectroscopy on the regions $55 \text{ G} \le B \le 75 \text{ G}$ and $160 \text{ G} \le B \le 180 \text{ G}$.

Experimentally, we realize this technique by modulating the magnetic field around its bias value B at the respective radio frequencies during a time between 100 ms and 500 ms and with a modulation amplitude between 100 mG and 500 mG. The modulation is applied with an additional single coil which is placed near the atom cloud and an oscillating current.

In figure 4.7 exemplary magnetic-field modulation spectra are shown for different magnetic field values. In those spectra the final atom number is plotted over the applied radio-frequency. The kinetic energy distribution of the associated atoms leads to a broadening of the spectrum and is given by a Maxwell-Boltzmann distribution. Therefore, by fitting a Maxwell-Boltzmann distribution we can extract the binding energy of the weakly bound molecular state E_b at the chosen magnetic field value. The



Figure 4.7: Exemplary magnetic-field modulation spectra at two different magnetic field values around each broad resonance. We extract the binding energy E_b by a fit with the Maxwell-Boltzmann distribution (solid line) to the data.

width of the distribution is in qualitative agreement with our temperature.

In figure 4.8 several of those spectra are plotted over a magnetic field range $B \in [72.0 \text{ G}, 72.6 \text{ G}]$ and a modulation frequency range $\nu_{\text{RF}} \in [-200 \text{ kHz}, -10 \text{ kHz}]$. The normalized atom number is visualized by a blue colour bar. Typically, we observe two kinds of features, which are exemplarily depicted in this figure. First, a spectrally narrow feature, which varies slowly in binding energy when the magnetic field is changed and second, spectrally broad features which are localized at a certain magnetic field. We associate very broad spectra to bound states with a high slope of the binding energy versus the magnetic field. Those bound states are responsible for narrow resonances. The slow variation of the narrow feature in binding energy may be associated with the weakly bound state which is responsible for the broad resonance.

As explained in figure 4.7, we extract the binding energy $E_{\rm b}(B)$ of this weakly bound state and show it in figure 4.9b (red dots).

Additionally, we focus on those two broad features with atom loss spectroscopy. In



Figure 4.8: Part of the magnetic field modulation spectroscopy data for $B \in [72.0 \text{ G}, 72.6 \text{ G}]$. The atom number N normalized by the atom number without field modulation N_0 is depicted as a function of frequency modulation ν_{RF} and magnetic field B. We observe a signature which varies slowly in frequency with magnetic field and another signature which high frequency-vs.-magnetic field slope. The former is associated to a weakly bound state responsible for the broad Feshbach resonance at 76.9 G and the latter to a bound state which creates a narrow resonance.

figure 4.9a the respective two data sets with a magnetic field resolution of 20 mG are presented. Those atom-loss spectra were obtained at a temperature of 500 nK and an initial atom number of $N_0 = 1.0 \cdot 10^5$ after a wait time of 500 ms.

Comparing the atom loss spectrum with the magnetic field modulation spectrum, it becomes apparent that the binding energy appears only on the low-field side of the two broad resonances. Besides, for both broad resonances, the binding energy $E_{\rm b}(B)$ varies very slowly with the magnetic field ($< \frac{1\,{\rm MHz}}{10\,{\rm G}}$). Furthermore, it shows a quadratic behaviour over a range of several gauss. As explained in section 4.1 this quadratic behaviour implies a coupling of the bound state with the open channel [10]. The bound state seems to have an open-channel dominated character over a broad field range. Despite many crossings with a dense sea of narrow resonances, the molecular bound state remains clearly visible. In fact, in figure 4.10 the signatures of avoided crossings with other bound states become apparent and are highlighted by a blue background. They reveal a coupling of the weakly bound dimer to another state which is stronger than usual. Furthermore, they can be associated to certain corresponding "broader-than-usual" narrow resonances.

By theoretical analysis in cooperation with Paul S. Julienne and Krzysztof Jachymski



Figure 4.9: Focus on the two broad Feshbach resonances of ¹⁶⁴Dy. a) Atom loss spectroscopy at 500 nK with a resolution of 20 mG. We observe two broad resonances that overlap with a sea of narrow resonances. b) The binding energy of weakly bound dimers extracted from the spectra of magnetic field modulation spectroscopy (red circles). The solid (dashed) lines are obtained by fitting the coupled-channel calculations (universal expression (4.2)) to the data. We extract $a_{\rm bg} \Delta_{\rm bg}$ from these fits (see table 4.1). c) The scattering length $a_s(B)$. The red circles have been obtained by converting the binding energy data of b) using coupled-channel calculations for $a_s(E_{\rm b})$. The solid lines are a fit to this data using the universal formula (4.1). The dashed lines represent the scattering length which results from the fit of the binding energy $E_{\rm b}(B)$ by the universal formula (4.2) assuming $a_{\rm bg} = 91 a_0$.

we could get further insight into the properties of the two broad Feshbach resonances. First, we could extract a high lower bound of the pole strength $s_{\rm res} \geq 10$ for both broad resonances by this analysis. This large $s_{\rm res}$ value indicates that the corresponding resonance is open channel dominated and shows a universal behaviour (see subsec. 4.1.1) [10]. Second, this analysis provides evidence that the open-channel dominated character of the broad Feshbach resonance persists across many narrow resonances and that away from them we can express the scattering length simply by equation (4.1). Furthermore, up to binding energies $E_{\rm b} \leq 1 \,\mathrm{MHz}$, the binding energy may be approximated by the universal expression (4.2). Next, Paul S. Julienne and Krzysztof Jachymski deduced the binding energy dependence on the scattering length $E_{\rm b}(a_s)$ and $a_{\rm bg} \Delta_B$ by coupled channel (CC) calculations. In figure 4.9c the consequent scattering lengths corresponding to the experimental measured binding energies are plotted as red circles. Additionally, we can extract $a_{\rm bg} \Delta_B$ by a fit of the universal formula (4.1) to the scattering lengths (see fig 4.9c, solid line). The parameters obtained by fitting both the universal expressions and CC calculations to the data are depicted in table 4.1. The fits to our binding energy data by coupled channel calculations and the



Figure 4.10: Focus on the atom loss spectrum and the binding energy on the low-field side of the resonance at $B_0 \approx 80$ G. In the blue areas, we observe strong avoided crossings in the binding energy and quite strong narrow Feshbach resonances. In those regions, we do not know the scattering length, whereas appart from these areas there are only narrow resonances which modify a_s only locally.

	resonance 1		resonance 2		
	B_0	$a_{ m bg}\Delta_B$	B_0	$a_{ m bg}\Delta_B$	
universal numerical CC	$76.9(5) \mathrm{G}$ $76.8(5) \mathrm{G}$	$\begin{array}{c} 2810(100)\mathrm{G}a_{0}\\ 2700(100)\mathrm{G}a_{0} \end{array}$	178.8(6) G 179.1(6) G	$\begin{array}{c} 2150(150)\mathrm{G}a_{0}\\ 2540(110)\mathrm{G}a_{0} \end{array}$	

Table 4.1: Resonance parameters which were obtained by fits of the binding energy to the universal expression (4.2) and to a numerical coupled channels calculation.

universal quadratic dependence $E_{\rm b} = -\left(\frac{B-B_0}{a_{\rm bg} \Delta_B}\right)^2$ are represented in figure 4.9b by a solid line and a dashed line respectively. Comparing the extracted parameters in table 4.1 and the fits of the binding energy, we see that the coupled-channel theory and the universal fit are in close agreement with each other for the lower-field resonance and in reasonable agreement for the higher-field one.

However, we can only extract the product $a_{\text{bg}}\Delta_B$. Therefore, a second method is required in order to determine the background scattering length a_{bg} and the resonance width Δ_B independently (see sec. 4.3.2.1).

Assuming the value of $a_{\text{bg}} \Delta_B$ which was obtained by the universal fit, we obtain the following scattering length for the lower-field resonance:

$$a_s(B) = a_{\rm bg} - \frac{2810\,{\rm G}\,a_0}{B - 76.9\,{\rm G}} \tag{4.32}$$

and for the higher-field resonance:

$$a_s(B) = a_{\rm bg} - \frac{2150\,{\rm G}\,a_0}{B - 178.8\,{\rm G}}\,.\tag{4.33}$$

This scattering length is also plotted in figure 4.9c as a dashed line (here: $a_{bg} = 91 a_0$, see subsubsec. 4.3.2.1). It is in close agreement with the CC results. However, it is only valid at magnetic field values away from the narrow resonances (see fig. 4.10) as they perturb the scattering length locally. In figure 4.10 the regions were the extracted scattering length $a_s(B)$ does not hold are highlighted by a blue background.

More information on the theoretical analysis of the broad Feshbach resonances is given in reference [6] and its supplemental material.

4.3.2.1 Localization of the zero crossing

With the methods presented up to now, we could only extract the product $a_{\text{bg}}\Delta_B$ from the data. By determining the zero-crossing of the scattering length, we can derive independently the resonance width $\Delta_B = B(a_s = 0) - B_0$ and thus also the background scattering length a_{bg} from the universal model (4.1).

Therefore, in this section several methods are explained in order to indicate the localization of this zero crossing.

In order to do so, we will at first focus on an exemplary Feshbach resonance (fig. 4.11b) Here, we can see that at one side of the resonance the atom number shows a local maximum. This peak in atom number corresponds to the zero-crossing of the scattering length. If the scattering length is zero $a_s = 0$, the collision and thus also the loss rate is minimal¹. This leads to the asymmetry of the lineshape. The fact that the zero-crossing is located at a magnetic field $B > B_0$ indicates that the background scattering length is positive $a_{bg} > 0$ (see fig. 4.3). For a negative scattering length $a_{bg} < 0$ this zero crossing would be on the other side of the Feshbach resonance $B < B_0$.

Hence, the asymmetry of the lineshape of narrow resonances gives a hint on the localization of the zero-crossing of the scattering length . The scattering length which orginates from a broad feature can be interpreted as a "local" background scattering length $\tilde{a}_{\rm bg}$ of a narrow resonance. If the local zero-crossing is on the high- (low-) field side of the resonance, $\tilde{a}_{\rm bg} > 0$ ($\tilde{a}_{\rm bg} < 0$).

In figure 4.11, we see several resonances that are sufficiently broad to extend over several data points. Here the asymmetry of their lineshapes can be observed. In figure 4.11a and b the lineshape corresponds to $\tilde{a}_{bg} > 0$. Both corresponding resonances are located on the low-field side of the broad resonance at $B_0 = 76.9$ G. Figure 4.11c and d show resonances on the high-field side of the broad resonance. Here, the symmetry

¹ The two-body collision rate is proportional to the elastic scattering cross-section (4.12), which is given in a dipolar gas by $\sigma_{a_s,D} = 8\pi \left(a_s^2 + \frac{4}{45}D^2\right)$ (4.15). Therefore, the dipolar interaction still induces collisions even when a_s vanishes. However, as D is independent of the magnetic field, a minimum in scattering length a_s corresponds to a minimum in scattering cross-section and thus also a minimum in collision rate.



Figure 4.11: Focus on narrow resonances [6]. (a) and (b) are examples of resonances at magnetic fields $B < B_0 = 76.9 \text{ G}$ which show a zero-crossing on their high-field side $(\tilde{a}_{\text{bg}} > 0)$. (c) and (d) show resonances with inverted asymmetry $(\tilde{a}_{\text{bg}} < 0)$, which occur for $B > B_0$. Thereafter, at 109 G (e) we observe a resonance with symmetric lineshape which corresponds to a \tilde{a}_{bg} close to zero.

of the lineshape is inverted, which corresponds to $\tilde{a}_{\rm bg} < 0$. In figure 4.11e, we see a resonance at $B = 109 \,\mathrm{G}$ with a very symmetric lineshape, which suggests that $\tilde{a}_{\rm bg}$ is close to zero. We conclude that the zero-crossing occurs in this magnetic field region. Between B = 110 and 120 G we do not observe sufficiently large resonances to conclude on the symmetry of the lineshape. Higher fields are so close to the other broad resonance at 178.7 G that it might provide the background for the narrow resonances. Therefore, we cannot withdraw any information on the broad resonance at 76.9 G in this magnetic field region.

Another hint on the localization of the zero-crossing of the scattering length give the final background atom number and temperature in between narrow resonances in the magnetic field range, where we expect the zero-crossing of the scattering length to occur. As already pointed out, the zero crossing corresponds to a minimal collision rate.



Figure 4.12: Focus on the atom number and temperature after a wait time of 500 ms around the area where a zero-crossing of the scattering length of the lower Feshbach resonance is expected [6]. A maximum in atom number is visible at $B_{\max,N} = (109 \pm 5)$ G. Likewise, a maximum in temperature can be located at $B_{\max,T} = (107 \pm 5)$ G. The corresponding area is highlighted by a gray background.

A low collision rate includes a low evaporative two-body collision rate and hampers thermalization. Therefore, at the zero-crossing a maximum in both the atom number and the temperature is expected.

In figure 4.12 the atom number and temperature are displayed as a function of magnetic field in the atomloss measurement. It becomes apparent that neglecting narrow resonances the background value of those observables varies slowly with magnetic field and shows a maximum. The maximum in atom number can be located at a magnetic field of $B_{\max,N} = (109 \pm 5)$ G. Likewise, a maximum in temperature can be located at $B_{\max,T} = (107 \pm 5)$ G.

Together with the analysis of the maximum in background atom number and temperature, we conclude that the scattering length in between narrow resonances varies slowly and crosses zero at a magnetic field $B = (108 \pm 5)$ G.

Consequently, the universal model (4.1) implies a background scattering length of $a_{\rm bg} = 91(15) a_0$ and a width of $\Delta_B = 31(6)$ G for the broad resonance at $B_0 = 76.9$ G.

4.4 Evidence of universal loss dynamics

In the previous section, we have characterized two Feshbach resonances and we have seen that $s_{\rm res} \gg 1$. Therefore, we expect to observe effects which are characteristic for broad resonances (see subsec. 4.1.1). One of those effects are universal loss dynamics, which can be observed at the center of broad resonances (see sec. 4.2). In this section, we show evidence that our observations of atom loss close to the poles of the resonances are indeed compatible with universal loss dynamics.

4.4.1 Estimation of the trap depth

In order to analyse the data on universal loss dynamics, the trap depth is an essential parameter. It can be expressed by the trap ratio $\eta = U/k_{\rm B}T$ and estimated analytically by estimating the optical dipole trapping potential $U_{\rm dip} = \frac{\hbar\Gamma}{8} \frac{\Gamma}{\Delta} \frac{I}{I_s}$ [44], where the knowledge of the beam waists enters with the intensity $I = \frac{2P}{\pi w_x w_y}$. Further the effect of gravity has to be estimated. However, the accuracy of the analytical estimation depends strongly on the knowledge of the trapping potential and since dysprosium is a strongly magnetic element, the residual magnetic field gradients may change the trap depth.

Statistical analysis using loss dynamics We improved our estimation of the trap depth by statistically analysing the respective Feshbach resonance spectrum in terms of final atom number and final temperature. For this purpose we used the rate equations for temperature and atom loss which we derived in section 4.2 by using loss dynamics equations adapted from [75] for finite *s*-wave scattering and including universal dipolar scattering. This was done by taking final experimental temperature and atom number in a certain magnetic field interval and checking the overlap with the final temperature and atom number obtained by solving numerically eq. (4.20) and (4.31) for different $\eta = U/k_{\rm B}T$. As the initial temperature and atom number are not known with certainity, they are chosen arbitrarily within a gaussian distribution around the estimated values.

For ¹⁶⁴Dy around the broad resonance at lower magnetic field, we estimated the scattering length by $a_s(B) = a_{bg} \left(1 - \frac{\Delta_B}{B - B_0}\right)$ (4.1) with $a_{bg} = 115 a_0$, $B_0 = 76.8$ G, and $\Delta_B a_{bg} = 2752.3 a_0$ G. Around the broad resonance at higher magnetic field, we chose $a_{bg} = 115 a_0$, $B_0 = 179.5$ G, and $\Delta_B a_{bg} = 2752.3 a_0$ G. These values were estimated to our best knowledge before the characterization of the broad resonances was complete. However, the slight deviation from the values we obtained finally should not have a considerable influence on the estimation of the trap depth. Taking the uncertainity of the initial temperature and atom number and the fluctuations in final temperature and atom number into account, we obtain finally a knowledge of η with an uncertainity of 20%.

4.4.2 Temperature dependence of atom loss

In order to study the temperature dependence of the universal losses, we do atomloss spectroscopy at different temperatures. In figure 4.13, the respective atom loss spectra are shown. The data presented in green is the same as in figure 4.5 and was taken at an initial temperature of $T_0 = 2.4 \,\mu\text{K}$. The data presented in blue was taken at $T_0 = 500 \,\text{nK}$ and the one in red at $T_0 = 300 \,\text{nK}$. Further information on the experimental conditions for the atom loss spectra are summarized in table 4.2.



Figure 4.13: Atom loss spectroscopy for different experimental starting conditions and wait times. The atom number N is normalized by the initial atom number N_0 and shown as a function of the magentic field B, with a resolution of 100 mG. The experimental starting conditions are given in table 4.2. The temperature dependence is well reproduced by the model of universal loss dynamics of a unitary Bose gas (see [75]) (solid horizontal lines). The shaded areas represent the uncertainity of the results of the model. Here a one-standard deviation enters the model on all experimental parameters (temperature, atom number, trap frequency and depth).

data set	$N_0 \ [10^5]$	$n_0 \; [10^{18} \mathrm{m}^{-3}]$	$T_0 [\mathrm{nK}]$	$\nu_{x,y,z}$ [Hz]	η	$t_{\rm wait} \ [{\rm ms}]$
green	1.8(2)	6.1(2.0)	2400(200)	70(10)	7.0(1.5)	2000
				123(15)		
				263(25)		
red	0.46(2)	4.6(1.6)	500(100)	44(5)	4.0(1.0)	500
				87(10)		
				145(15)		
blue	1.0(2)	3.7(1.3)	450(100)	52(5)	5.0(1.0)	500
	. ,			23(5)		
				260(25)		

Table 4.2: Experimental conditions under which the atom-loss spectra presented in figure 4.13 and 4.14 have been taken. N_0 and T_0 are the initial atom number and temperature respectively. $\eta = U/k_B T$, where U is the trap depth, is estimated as explained in 4.4.1. $\nu_{x,y,z} = \omega_{x,y,z}/2\pi$ are the trapping frequencies in the three dimensions and t_{wait} is the wait time before we record the final atom number N.

In figure 4.13, we observe that the final atom number reaches a minimum that is the same on the two broad resonances. Furthermore, for a lower temperature, the saturation is reached in a narrower field region and at a lower final atom number.

At first, we compare our findings with the loss dynamics model of a Bose gas in the unitary regime where $a_s k \gg 1$, which has been developed in [75]. The respective coupled rate equations for atom number (4.18) and temperature evolution (4.30)predict the final atom number taking into account two-body evaporation and threebody recombination. In order to obtain the expected final atom numbers, we solve numerically these coupled equations using the parameters given in table 4.2. These observables are calibrated on our experiment and the only unknown parameter of the model is the inelasticity parameter η_* (see eq. (4.9)), which describes the recombination efficiency for three atoms in contact to recombine to a dimer and a free atom. We then compare the final atom numbers obtained by this model with a certain η_* to the three data sets shown in figure 4.13. The shaded area represents the range of final atom number, which we obtain by varying the input parameters of the model by one standard deviation. The respective uncertainities are given in table 4.2. Since the saturation level is similar for the low- and high-field resonance at the same temperature $T_0 = 2.4 \,\mu\text{K}, \eta_*$ appears to be identical for both broad resonances. Therefore, we choose a fixed η_* to compare the model to the three data sets. Thereby, we find that this model reproduces well our final atom number at resonance within experimental uncertainities, when the following inelasticity parameter is chosen:

$$\eta_* = 0.07^{+0.17}_{-0.05} \,. \tag{4.34}$$

This value is comparable to the lowest reported values for alkali-metal atoms [76, 85–89]. This implies that lifetimes in dysprosium are comparable to alkalis in the unitary regime. This experimental estimate of the inelasticity parameter is valuable, since no theoretical prediction is available. However, a more precise measure could be obtained by a systematic study of $L_3(T)$ at resonance.

4.4.3 Loss traces

Subsequently, we compare our findings not only with the loss dynamics model of a Bose gas in the unitary regime, but also at finite scattering length a_s . In this combined model, we decide whether either the conditions of the unitary regime (coupled rate equations (4.18) and (4.30)) or of finite scattering length are fulfilled ((4.20) and (4.31)). For this purpose, we decide, whether γ_2^U (eq. (4.19), unitary regime) or $\gamma_2^{a_s,D}$ (eq. (4.21), finite scattering length) is smaller and plug the respective value into the coupled rate equations in order to describe evaporative two-body losses. Similarly, we analyze, whether γ_3^u/T_0^2 (eq. (4.11), unitary regime) or $\gamma_3^{a_s}$ (eq. (4.8), finite scattering length) is smaller and choose either coupled rate equations for a Bose gas at unitarity or for finite scattering length respectively in order to describe three-body recombination.

This way, the model represents that the atom losses depend on the scattering length. However, for a diverging scattering length in the vicinity of the resonance pole the



atom losses are limited within the unitary regime.

Figure 4.14: Atom loss spectroscopy for the same experimental starting conditions and wait times as in figure 4.13 (see table 4.2). The loss traces in between narrow resonances are in good agreement with the model of universal loss dynamics for finite scattering length and at unitarity respectively (see ref. [75] and section 4.2) using the scattering length which was obtained by fitting the binding energy data to the universal formula (4.2) (dashed lines) or to a numerical coupled channels calculation (solid lines) (see table 4.1).

Despite the somewhat sharp transition between unitarity and the regime of finite scattering length, the experimental data is in good agreement with the final atom numbers calculated numerically using this model (see fig. 4.14). Since we compare this model to the green and blue data sets of figure 4.13, we use also the parameters depicted in table 4.2 for the respective numerical calculation. Furthermore, we use the inelasticity parameter $\eta_* = 0.07$ (see subsec. 4.4.2) and the scattering length $a_s(B)$, which was obtained by fitting the binding energy data to the universal formula (4.2) ("universal scattering length", dashed lines) or to a numerical coupled channels calculation ("CC scattering length", solid lines) (see table 4.1 for the respective scattering length parameters).

At the lower resonance the difference between the universal scattering length and the CC scattering length is small, for the higher resonance, it is somewhat stronger pronounced.

On the low-field side of the resonance, the calculated final atom numbers are lower than on the high-field side of the resonance, since we assume the zero-crossing of the scattering length to be on the high-field side of the resonance (see subsubsec. 4.3.2.1 and 4.3.1.2). For the higher resonance these calculated final atom numbers reproduce well the experimentally observed final atom numbers in between narrow resonances for $B \ge B_0 \pm 10$ G. On the high-field side of the lower resonance the loss traces in between narrow resonances are also well reproduced by the model, whereas on its low-field side the above described tendency is not that clear. The calculated loss traces lie beneath the normalized experimental atom numbers. However, unphysically normalized atom numbers $N/N_0 > 0$ can be observed in this region. This is probably due to the fact, that the estimated initial atom number N_0 is too low in this magnetic field region.

Moreover, the smaller the relative trap depth η , the lower the final calculated atom number on the low-field side of the resonance. This is due to enhanced evaporative loss in the energetically more shallow trap.

A signature of the formation of Effimov trimers [76] would be the experimental observation of the dip in atom number on the high-field side of the resonance, which can be observed in the calculated loss traces. However, this could not yet be observed in the data.

Last but not least, the experimental observation that for a lower temperature the saturation is reached in a narrower field region is well reproduced by the combined model we used here.

CHAPTER 5

Conclusion and outlook

Conclusion

In this thesis, we have seen the experimental realization of demagnetization cooling of ultracold dysprosium atoms. We could observe a cooling rate of $\frac{dT}{dt} = -(3.2 \pm 0.3) \frac{\mu K}{s}$ during 3s of demagnetization cooling at an initial temperature of $T_0 = 27 \,\mu\text{K}$. The atom number stayed approximately constant during this time. Consequently, we could cool the sample down to approximately $17 \,\mu K$ within 2s in this first proof of principle measurement. After the optimization of the magnetic field amplitude parallel to the optical pumping beam to $225 \,\mathrm{mG}$, a temperature of around $16\,\mu\mathrm{K}$ could be reached after 2s of demagnetization cooling. Subsequently, we optimized the polarization and minimized the magnetic field amplitude perpendicular to the optical pumping beam. Thereby, we reached a temperature close to $15 \,\mu \text{K}$ within 2 s. However, theoretically temperatures beyond $7\mu K$ could be possible within 2s of demagnetization cooling. In order to ensure a fast rethermalization and hence a large trapping frequency compared to the optical pumping and dipolar relaxation rate, we adjusted the confinement from a single optical dipole trap to a crossed optical dipole trap. However, as the initial temperatures were quite low around $T_0 = 7.4 \,\mu\text{K}$ after a first evaporative step in this trap configuration, only a little cooling effect of around $1\,\mu\text{K}$ could be observed and atom loss outperformed the effect of demagnetization cooling. Nevertheless, we could qualitatively probe the dependence of temperature and atom number on the optical pumping rate and the detuning in this trap configuration. We observed a decrease in atom loss and temperature when coming to larger detunings and lower optical pumping rates. Finally, we compared the efficiency of non-optimized demagnetization cooling with evaporative cooling. We obtain at an initial temperature of $38\,\mu\mathrm{K}$ an efficiency of demagnetization cooling of $\chi_{\text{demag}} = 7.95 \pm 1.58$ within 1 s, which clearly outperforms the efficiency of evaporative cooling of $\chi_{evap} \approx 4$ that can only be reached with substantial effort. However, our actual cooling scheme includes post-cooling with the 626 nm transition in the single ODT 1. By this cooling step we reach a similar cooling efficiency of $\chi_{626} = 6.7$ [5]. It might be possible to improve our efficiency of demagnetization cooling further by the measures substantiated in the outlook. However, this might imply some effort.

Moreover, we have presented two broad Feshbach resonances, which emerge from a background of an apparent chaotic distribution of Feshbach resonances. We characterized both broad resonances by atom loss spectroscopy and magnetic field modulation spectroscopy. In cooperation with Paul S. Julienne and Krzysztof Jachymski we analyzed the data in terms of coupled channel calculations and the universal expressions. Thereby, we deduced for both resonances and by both methods the position of the resonance B_0 and the product of background scattering length and resonance width $a_{\rm bg}\Delta_B$. The respective values are given in table 4.1. For the lower field resonance, CC calculations and the universal expressions are in close agreement and for the higher field resonance in reasonable agreement. Furthermore, we evaluated the symmetry of the lineshape of the lower field Feshbach resonance in terms of temperature and atom number in order to localize the zero-crossing of the scattering length and derive the resonance width and background scattering length independently from that. We supported our analysis additionally by considering the symmetry of the lineshape of neighbouring narrower Feshbach resonances. Thereby, we could localize the zerocrossing of the lower field Feshbach resonance at a field B = 108(5) G, which implies a width $\Delta_B = 31(6)$ G and $a_{bg} = 91(15) a_0$. Assuming the same background scattering length for both Feshbach resonances, the scattering length $a_s(B)$ in between narrow resonances could be approximated by the following universal expression for the lower field resonance

$$a_s(B) = 91 a_0 - \frac{2810 \,\mathrm{G} \,a_0}{B - 76.9 \,\mathrm{G}} \tag{5.1}$$

and for the higher field resonance

$$a_s(B) = 91 a_0 - \frac{2150 \,\mathrm{G} \,a_0}{B - 178.8 \,\mathrm{G}} \,. \tag{5.2}$$

Moreover, we obtained an $s_{\rm res} \gg 1$ for both resonances. Therefore, we expected effects which are characteristic for broad resonances and saw indeed evidence that the loss traces we observed by atom loss spectroscopy in the vicinity of the resonances are compatible with universal loss dynamics (see figure 4.14). We could even extract an estimate of the trap depth $\eta = U/k_{\rm B}T$ with an uncertainity of 20% by comparing the model of universal loss dynamics with our Feshbach spectrum in terms of final atom number and temperature. The estimated values are given in table 4.2 for the different data sets. By evaluating the temperature dependence of three-body losses in the unitary regime, we could additionally estimate the elasticity parameter, which characterizes the efficiency with which three atoms in contact recombine to a dimer and a free atom: $\eta_* = 0.07^{+0.17}_{-0.05}$. This value is close to the lowest values, which have been reported for alkali-metal atoms [76, 85–89] and implies similar lifetimes in the unitary regime. Last but not least, we could use the universal description of our broad resonances presented here, in order to create pure Bose-Einstein condensates with up to $25 \cdot 10^3$ atoms and lifetimes > 1 s.

Outlook

In order to optimize demagnetization cooling the following steps can be taken. First, the optical pumping rate Γ_{op} can still be optimized systematically. Next, an optimal magnetic field ramp can be found by analyzing the effect of demagnetization cooling subsequently at different starting temperatures and magnetic field values. Furthermore, the effect of demagnetization cooling in the confined trap, but without an initial evaporative cooling step can be evaluated. This would permit a higher initial temperature and thus a higher initial cooling rate. Moreover, in chromium the efficiency of demagnetization cooling could be improved from $\chi_{Cr} \approx 6$ to > 17 by localizing the node position of the groundstate wavefunction and thus suppressing light assisted collisions (see sec. 3.4.4) [4, 46]. Therefore, probing the dependence of the performance of demagnetization cooling on the detuning up to approximately 8 GHz $\approx 84\,000\,\Gamma$ might help to improve the efficiency of demagnetization cooling. Up to now, our maximal detuning from resonance was $\Delta = -100$ MHz.



Figure 5.1: Microscopic droplet crystal growth. a) Representative droplet patterns with droplet numbers ranging from two to ten. b) Scheme of the experimental procedure: We prepared a stable BEC at $a_s \approx 3/2 D$. When decreasing the scattering length to $a_s \approx a_{\rm bg}$, the atoms clustered to droplets in a triangular pattern. c) Mean number of atoms in dependence of visible droplets, with the standard deviation as error bars. Image taken from [24].

Recently, we could observe discrete and stable droplets arranging in a triangular structure (see figure 5.1a) by creating a BEC and then tuning the scattering length a_s with a magnetic Feshbach resonance such that $a_s < 2/3 D$ (see figure 5.1b) [24]. The number of the droplets we observe is growing with increasing atom number (see

figure 5.1c). Therefore, we might be able to observe some kind of droplet crystal increasing the atom number by means of demagnetization cooling.

Besides, the estimate of the scattering length, which we obtained by analyzing the two broad Feshbach resonances, can be improved by the measurement of the density distribution of Bose-Einstein condensates. Furthermore, in the analysis of the loss traces in section 4.4.3 the prospect of studying the few-body physics of bosonic dipoles discernibled. Conceivably, it might be possible to investigate for example the impact of the magnetic dipole-dipole interaction on the energy of Efimov states [77].

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Danksagung

Schließlich möchte ich an dieser Stelle all denen danken, die mir zur Seite gestanden, mich unterstützt und damit maßgeblich zum Gelingen dieser Arbeit beigetragen haben. Mein besonderer Dank gilt dabei

- **Prof. Dr. Tilman Pfau** für die Aufnahme am 5. Physikalischen Institut, den Vorschlag meines Vortrags über "Quantum Chaos in ultracold collisions of gas-phase erbium atoms", der mein Interesse für das Dysprosiumexperiment geweckt hat, und seinen ansteckenden Enthusiasmus für die Wunder der Physik.
- **Prof. Dr. Clemens Bechinger** in seiner Funktion als Mitberichter und für sein Interesse an meiner Arbeit.
- Jahn Rührig and the whole Dysprosium team including Thomas Maier, Holger Kadau, Matthias Schmitt, Matthias Wenzel and Igor Ferrier-Barbut for the warm welcome in the PI5 dipolar gases family. Vielen Dank für die gute Laune, viele Späße an der Grenze des politisch Korrekten und vor allen Dingen die unermüdliche Hilfe. Merci beaucoup pour toutes les explications sur la physique et les cours de Mathematica.
- Allen Mitarbeitern des 5. Physikalischen Instituts für die Hilfsbereitschaft und die angenehme Arbeitsatmosphäre
- Mona für die entspannenden Joggingrunden und die damit verbundenen Gespräche.
- Joni, Marco, Matze, Tobi und Katha für die Begleitung durch das Physikstudium.
- Meinen **Eltern**, meiner **Schwester** und **Samir** für die liebevolle, großzügige und langjährige Unterstützung in jeder Hinsicht.