## Laser Cooling Towards a Dipolar Quantum Gas of Erbium Atoms

Nathaniel B. Vilas Emmanuel College

Professor Zoran Hadzibabic, Supervisor

A thesis submitted for the degree of Master of Philosophy

Department of Physics Cavendish Laboratory July 2018



### Declaration

This thesis is the result of my own work and includes nothing which is the outcome of work done in collaboration except where specifically stated otherwise. I have not submitted, nor do I plan to submit, the same or similar work for another degree at the University of Cambridge or any other institution. The length of this thesis adheres to the limit of 15,000 words, exclusive of tables, footnotes, bibliography, and appendices, set out by the Degree Committee for the Faculty of Physics and Chemistry.

> Nathaniel B. Vilas July 2018

### Abstract

This thesis describes the design and construction of an apparatus meant to achieve degenerate quantum gases of erbium atoms. Erbium possesses a large magnetic dipole moment of  $7\mu_B$ , making it an ideal candidate for the study of long-range interactions in many-body quantum systems, and the present experiment aims to combine the study of long-range dipole-dipole interactions with quasi-uniform trapping geometries to study, e.g., roton physics and supersolidity. The starting point is an atomic beam of Er atoms heated to 1200°C, which we will collimate via transverse laser cooling and decelerate with a Zeeman slower, home-built over the course of this work. Both processes take place on a broad, 30 MHz transition at 401 nm, allowing for efficient cooling and a compact Zeeman slower only 40 cm long. They are designed with several tunable degrees of freedom, which we can systematically optimize for maximum slow-atom flux at the experiment chamber. The next step is to load a magneto-optical trap (MOT) operating on a narrow, 200 kHz transition at 583 nm, which will enable the production of fully spin-polarized atom clouds at temperatures as low as 5  $\mu$ K. Throughout this work we have paid particular attention to the control and compensation of magnetic fields around the experiment chamber, leaving the door open to future experiments requiring magnetic field tuning at near-milliGauss level. This level of control is particularly important for erbium, which along with its large magnetic moment also possesses many narrow, closelyspaced Feshbach resonances. The results presented here represent a good starting point from which to work towards MOT and optical trap loading, followed by forced evaporation down to quantum degeneracy.

### Acknowledgements

Ultracold atom experiments, as it turns out, require an enormous investment of time and resources, and without others around me on the erbium experiment, in the Hadzibabic-Smith group, and in AMOP more broadly, we wouldn't have made half as much progress in the past year, nor would I be half as sane. I unquestionably have to start at the top, with Zoran and Rob. Eighteen months ago, Zoran took on the role of being my supervisor only later to find out that erbium would be fleeing Cambridge after only a year; despite the hard deadline on the machine officially being 'his,' he has consistently offered insightful advice and asked the right questions to help push the experiment forward. Rob, meanwhile, has been a daily presence who almost inevitably offers clarity on even the most confounding problems, typically with a few moments of humor thrown in for good measure. I owe them both a great deal of gratitude for trusting that I knew what I was doing when they offered me a place in the group.

I next have to thank my 'other half' on the erbium experiment, Milan Krstajic, who not only single-handedly designed and built the vacuum system but also consistently took on the most laborious tasks in the lab<sup>1</sup> and tolerated my (often rather unenlightened) questions all year. Also working on erbium for several months this past year, Tobias Schaich and Tanish Satoor made valuable contributions to the experiment during their Part III projects, as has Sean Seet over the last month. Furthermore, I thank Kevin Mott and Gavin Ross, respectively, for assistance in the workshop and with 3D printing.

Finally, I thank the other students and postdocs in the group, too numerous to list here, for many enlightening conversations, for equipment sharing,<sup>2</sup> and for a friendly atmosphere that makes the group buzz. Along with the physics, I expect recollections of 'burger day', the Cavendish custard, and football to remain with me well after leaving Cambridge, and that is a testament to the exceptional collection of people assembled here.

 $<sup>^1{\</sup>rm Hours}$  spent aligning and realigning (...and realigning yet again) a wayward frequency doubling cavity from scratch come to mind.

<sup>&</sup>lt;sup>2</sup>One and a half power meters are apparently not enough for four experiments.

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# Chapter 1 Introduction and Background

Over the past 25 years, ultracold atoms have emerged as an exciting platform with which to study many-body quantum physics [1]. Experimental advances have made it possible to implement precise Hamiltonians with dynamical control of both interparticle interactions and confining potentials in such systems, and techniques for time-resolved extraction of thermodynamic quantities and other properties have enabled breakthrough studies of, for instance, the superfluid to Mott insulator transition in Bose-Hubbard models [2] and the BEC-BCS crossover in fermionic gases [3–5]. The ability of these systems to realize and study a range of Hamiltonians recalls Feynman's idea of an analog quantum simulator [6].

Until recently, interactions in ultracold atom systems were exclusively of a shortrange, van der Waals character. In recent years, however, the introduction of long range, anisotropic dipole-dipole interactions (DDI) has led to the observation of qualitatively new phenomena. This thesis describes the first year of construction of a new such ultracold atom experiment, with highly magnetic erbium atoms, which will ultimately allow for the study of *homogeneous* many-body physics with dipolar interactions.

#### 1.1 Bose-Einstein Condensation

A necessary starting point to establish the "quantumness" of a many-body system, the phenomenon of Bose-Einstein condensation (BEC) was first proposed theoretically by Einstein [7] and Bose [8] in the 1920s and occurs when a nontrivial fraction of bosons in an ensemble "condense" into their quantum-mechanical ground state. The transition from a classical, thermal ensemble to such a macroscopically "quantum" system occurs when the average inter-particle spacing becomes comparable to the quantum length scale associated with each particle – namely, the thermal de Broglie wavelength,

$$\lambda_T = \frac{h}{\sqrt{2\pi m k_B T}} , \qquad (1.1)$$

where  $h = 2\pi\hbar$  is the Planck constant, *m* is the particle mass,  $k_B$  is the Boltzmann constant, and *T* is the temperature of the system.

Such a condensate can be formed out of a classically thermal system by increasing the particle number density, n = N/V (where N is the total particle number and V is the system volume), decreasing the temperature so as to increase  $\lambda_T$ , or by some combination of the two. A convenient parameter used to quantify the relationship of these competing values is the so-called *phase-space density*,

$$\rho = n\lambda_T^3 , \qquad (1.2)$$

which denotes the average number of particles per "quantum volume" given by the cube of the de Broglie wavelength. The BEC phase begins to emerge when  $\rho$  is of order unity. For a non-interacting gas in a uniform potential (i.e. a box), the phase-space density at condensation is  $\rho \approx 2.612$ .<sup>1</sup>

BEC manifests itself in a variety of systems, including superfluid <sup>4</sup>He [10] and superconducting metals (in which bosonic pairs of fermionic electrons form the condensate) [11]. Due to their ready manipulability by electromagnetic fields (specifically laser radiation and magnetic coils), BECs of dilute atomic vapor have further established themselves as fertile grounds for research over the past 20 years. To achieve these gaseous condensates requires low particle densities between  $10^{12}$  and  $10^{15}$  cm<sup>-3</sup>; at higher densities a cold gas will attempt to return to its global ground state – a solid – through rapid three-body recombination.<sup>2</sup> Low temperatures on the order of  $10^{-7}$  K are necessary to 'balance' this low spatial density and achieve high-enough  $\rho$  for condensation.

BEC in cold atoms was first realized in 1995 with <sup>87</sup>Rb [13], <sup>23</sup>Na [14], and <sup>7</sup>Li [15, 16] through the use of a variety of laser cooling and trapping techniques developed over the preceding two decades (see chapter 2) [17–23]. To date, approximately twenty distinct bosonic species have been cooled to quantum degeneracy (see table 1.1).

<sup>&</sup>lt;sup>1</sup>See, for instance, [9] for a more complete treatment of this case, as well as that of the harmonic oscillator potential used so frequently in experiments.

<sup>&</sup>lt;sup>2</sup>For comparison, room temperature air has a density of  $\sim 10^{19}$  cm<sup>-3</sup>, and typical solids have densities near  $10^{23}$  cm<sup>-3</sup> [12].

<sup>&</sup>lt;sup>3</sup>The Innsbruck group later condensed both <sup>166</sup>Er and <sup>170</sup>Er [41], though the accomplishment

<sup>87</sup> Rb [13]	$^{23}$ Na [14]	<sup>7</sup> Li [15, 16]	${}^{1}\mathrm{H}$ [24]	$^{85}$ Rb [25]	$^{4}$ He [26, 27]
1995	1995	1995	1998	2000	2001
$^{41}$ K [28]	$^{133}Cs$ [29]	$^{174}$ Yb [30]	<sup>52</sup> Cr [31]	$^{39}$ K [32]	$^{40}Ca$ [33]
2001	2003	2003	2005	2007	2009
$^{84}$ Sr [34, 35]	<sup>86</sup> Sr [36]	<sup>164</sup> Dy [37]	<sup>168</sup> Er [38] <sup>3</sup>	<sup>160</sup> Dy [39]	<sup>162</sup> Dy [39]
2009	2010	2011	2012	2015	2015

**Table 1.1:** Bosonic atomic species condensed to date, along with year achieved and relevant publication. Dipolar species are indicated in boldface. Adapted from [12, 40].

#### 1.1.1 BEC in a Uniform Potential

For purposes of experimental ease, the vast majority of ultracold atom experiments are conducted in approximately harmonic potentials – such is the form created by the Gaussian beams used in optical dipole traps (ODT) and by standard magnetic traps.<sup>4</sup> Yet from a theoretical perspective, it is not the harmonic but the uniform potential in which BEC (and many other phenomena) can most readily be understood.<sup>5</sup> As compared with uniform potentials, harmonic traps introduce an additional length scale to physical phenomena related to the various trapping frequencies  $\omega_x$ ,  $\omega_y$ , and  $\omega_z$ ,<sup>6</sup> and atomic cloud densities vary with position in inhomogeneous potentials. Indeed, considerable effort has been put into the extraction of uniform system phenomena from harmonic experiments, for instance through use of the local density approximation (LDA) or by selectively studying only the center of a harmonically trapped gas [44–47]. Despite these efforts, there remain important many-body phenomena that simply cannot be adequately simulated in harmonic systems [48-51]. In particular, the LDA breaks down near phase transitions, making the study of critical behavior difficult in harmonic traps; furthermore, the imposition of a harmonic oscillator length becomes inconvenient in scenarios where other fundamental length scales disappear (e.g. the study of *unitary* Bose gases, where the scattering length a, which characterizes the strength of atomic interactions, diverges).<sup>7</sup> Fi-

was never specifically highlighted in a publication.

<sup>&</sup>lt;sup>4</sup>See, for instance, [42]. More generally, all trapping geometries are approximately harmonic for small displacements from the minimum, except in those potentials where the quadratic term has been specifically engineered away, e.g. in box potentials or those with high power laws  $r^n$ , n > 2.

<sup>&</sup>lt;sup>5</sup>Note the treatment of BEC in any undergraduate thermal and statistical physics text like [43], as compared with the treatment in a specialized text like [9].

<sup>&</sup>lt;sup>6</sup>This characteristic length is defined by  $l_i = \sqrt{\hbar/(m\omega_i)}$  along each axis.

<sup>&</sup>lt;sup>7</sup>One nice feature of homogeneous unitary gases should be that the *only* relevant length scale is the average interparticle spacing  $n^{-1/3}$ .

nally, in a uniform potential the BEC spreads out over most of the trap<sup>8</sup> (unlike in a harmonic trap, where the condensate tends towards the potential minimum and spatially distinguishes itself from the thermal atoms), which allows for lower densities and thereby minimizes three-body losses in, e.g., strongly interacting Bose gases.

Recently in our group, the first demonstration of Bose-Einstein condensation in a uniform potential was performed through the use of a spatial light modulator (SLM) to alter the phase characteristics of trapping light [52, 53]. Early subsequent experiments quickly revealed phenomena unobservable in harmonic systems [49–51]. The realization of a box potential has since been replicated in ultracold Fermi gases [54], and homogeneous ultracold gas experiments have become increasingly common, including in two-dimensional systems [55]. Building on these successes, it is our goal in the experiment described here to implement uniform trapping geometries for a dipolar gas.

#### 1.1.2 Atomic Interactions and Dipolar BEC

Due to the relative simplicity of their internal energy structure and the corresponding accessibility of closed cooling transitions, alkali atoms such as Rb, Na, Li, and K have been the focus of a significant proportion of ultracold atom experiments to date. In quantum gases of these atoms, interactions are purely *s*-wave and are therefore describable as isotropic, short range contact interactions quantified by a single parameter, the scattering length  $a_s$ .<sup>9</sup> The strength of these interactions can be tuned with magnetic bias fields through the use of Feshbach resonances (FRs) [56– 58], and much of current research in alkali systems takes advantage of this convenient phenomenon.

Studies with Rydberg atoms and polar molecules have allowed the introduction of long-range interactions to the cold atom toolbox, though both of these systems typically have short lifetimes. An alternative approach is to use atoms with large magnetic dipole moments. Though their interactions exhibit weaker longrange character than those of polar molecules, systems of ultracold dipolar atoms are technically easier to realize and are comparatively long-lived. The dipole dipole interaction (DDI) potential between such atoms takes the form (with the dipole

<sup>&</sup>lt;sup>8</sup>At least in the presence of interactions. For a = 0 it should approach the shape of the trap ground state.

<sup>&</sup>lt;sup>9</sup>A detailed description of this theory, as well as that of Feshbach resonances, can be found in [9].



Figure 1.1: Schematic representation of the dipole-dipole interaction. (a) Plot of equation 1.3, where red sections correspond to repulsive interactions and blue sections correspond to attractive interactions. Regions inside the black curves are made homogeneous for plotting purposes; in reality the potential diverges towards  $\mathbf{r} = 0$ . Note that the potential vanishes at  $\theta_{\rm crit} = \arccos(1/\sqrt{3}) = 54.7^{\circ}$ . (b) Schematic representation of the geometry of two polarized dipoles.

moments polarized along one axis) [59]

$$U_{\rm dd}(\mathbf{r}) = \frac{\mu_0 \mu^2}{4\pi} \frac{1 - 3\cos^2\theta}{r^3} , \qquad (1.3)$$

where  $\mu_0$  is the vacuum permeability,  $\mu$  is the magnetic dipole moment, **r** gives the relative atomic position, and  $\theta$  is the angle between **r** and the polarization axis (figure 1.1). It can immediately be seen that the DDI is long range and anisotropic, in direct contrast to the contact potential described above.<sup>10</sup> For direct comparison to the *s*-wave contact interaction we define a 'dipolar length'

$$a_{\rm dd} = \frac{\mu_0 \mu^2 m}{12\pi\hbar^2} , \qquad (1.4)$$

where m is the particle mass [59]. The ratio of dipolar and contact length scales,

$$\varepsilon_{\rm dd} = a_{\rm dd}/a_s , \qquad (1.5)$$

thus becomes a convenient quantity for keeping track of the important interactions in a BEC (larger  $\varepsilon_{dd}$  implies more strongly dipolar characteristics). In particular,  $a_{dd}$  in equation 1.4 is defined so that a uniform, 3D condensate becomes unstable

<sup>&</sup>lt;sup>10</sup>A potential of the form  $r^{-n}$ , with  $n \leq D$ , is typically considered long range for a *D*-dimensional system [59, 60]; thus the DDI is long-range in three dimensions but short-range in one or two. Contrast this with the van der Waals potential,  $-C_6/r^6$ .

to collapse when  $\varepsilon_{\rm dd} \geq 1$  [59].<sup>11</sup>

Quantum degeneracy with strongly dipolar atoms was first achieved in 2005 using <sup>52</sup>Cr ( $\mu = 6\mu_B$ , where  $\mu_B$  is the Bohr magneton) by the group of Tilman Pfau at the University of Stuttgart [31], and later in the lanthanide elements Dy ( $\mu = 10\mu_B$ ) and Er ( $\mu = 7\mu_B$ ) by groups at Stanford University (Ben Lev) [37, 61] and the University of Innsbruck (Francesca Ferlaino) [38, 62] in 2011 and 2012, respectively. Experiments with chromium immediately demonstrated recognizably dipolar phenomena, including changes to the aspect ratio of expanding atomic clouds due to DDI [63, 64], a demagnetization cooling method related to the coupling of spin and motional degrees of freedom [65], dependences in cloud stability on trapping geometry [66, 67], and a dramatic *d*-wave collapse displaying the signature anisotropy of the DDI [38, 68].

Despite these observations, in <sup>52</sup>Cr the relative strength of DDI compared with the s-wave contact interaction is small ( $\varepsilon_{\rm dd} \approx 0.16$ ) [69]. By using a Feshbach resonance to decrease the latter to strengths comparable to the DDI, strong dipolar effects could be observed [70]; nonetheless, in heavier, more strongly dipolar elements like Dy and Er, DDI can dominate ( $\varepsilon_{dd} \geq 1$ ) without the need to tune  $a_s$ [37, 71], thereby avoiding the three-body losses associated with experiments near a FR [37, 70]. Experiments in these systems have demonstrated exciting phenomena such as dense Feshbach spectra indicative of quantum chaotic scattering [72, 73] and the appearance of metastable 'quantum droplets' at large  $\varepsilon_{dd}$  [71, 74–76], including the observation of a 'scissors' oscillation mode in such a droplet [77]. Further recent advances include the study of an extended Bose-Hubbard model with DDI [78] and an observation of first experimental signatures for the roton mode [79]. Finally, ultracold atoms with DDI were recently shown as an effective tool to study thermalization behavior in near-integrable systems [80]. Various novel phenomena predicted in dipolar BECs remain largely unexplored in experiments, making magnetic elements like erbium particularly attractive choices for new experiments like that described in this thesis.

<sup>&</sup>lt;sup>11</sup>Technically, for  $\varepsilon_{\rm dd} \ge 1$  the Bogoliubov phonon-excitation frequencies become imaginary [59]. Intuitively, this dipolar collapse can be understood from the anisotropy of equation 1.3: polarized dipoles sitting 'head-to-tail' ( $\theta = 0$ ) experience an attractive potential and collapse in the same way as a BEC with negative  $a_s$ .

Isotope	<sup>162</sup> Er	$^{164}\mathrm{Er}$	$^{166}\mathrm{Er}$	$^{167}\mathrm{Er}$	$^{168}\mathrm{Er}$	$^{170}\mathrm{Er}$
Statistics	boson	$\operatorname{boson}$	$\operatorname{boson}$	fermion	boson	$\operatorname{boson}$
Abundance	0.14%	1.61%	33.6%	23.0%	26.8%	15.0%

**Table 1.2:** Natural abundances and quantum statistical behavior for all stable Er isotopes. Adapted from [41].

#### 1.2 Erbium

Along with its large magnetic dipole moment of  $\mu = 7\mu_B$  (corresponding to  $a_{dd} = 66a_0$ , where  $a_0$  is the Bohr radius), erbium has several properties that make it an appealing candidate for studies of ultracold dipolar gases. In particular, it has a number of abundant bosonic and fermionic isotopes (table 1.2) which allow access to different quantum statistics as well as a range of background *s*-wave scattering lengths<sup>12</sup>; these can readily be tuned via an assortment of Feshbach resonances available at low field [72]. Further, Er possesses several (nearly) closed electronic transitions on which to efficiently laser cool.

Erbium has atomic number Z = 68 and, in the ground state, consists of a full 6s valence shell along with twelve 4f electrons (these partially filled, high-l orbitals lend Er its large magnetic moment [41]). The ground state has total orbital angular momentum L = 5, total spin S = 1, total angular momentum J = 6, and even parity; it is notated [Xe] $4f^{12}6s^2$  ( ${}^{3}H_6$ ). The energy level structure of erbium, up to states 25000 cm<sup>-1</sup> above the ground state, is shown in figure 1.2. The first laser cooling work with erbium came in the mid 2000s from the group of Jabez J. McClelland at NIST (USA), who identified five promising  $J \rightarrow J + 1$  laser cooling transitions [81] before demonstrating broad [82] and narrow-line [83] magneto-optical traps (MOTs) of Er. For the remainder of this thesis we focus on only the two highest-energy cooling transitions identified, at 401nm and 583nm, respectively.

Because the fine structure (spin-orbit interaction) is comparable to the gross structure (Coulomb interaction) in highly excited valence shells for lanthanides like Er, the LS coupling scheme breaks down and we instead employ a special case of jj-coupling. In this  $J_1J_2$ -coupling picture, the n = 6 valence electrons LS-couple into a state with total angular momentum  $J_2$ ; all inner electrons similarly couple to a total angular momentum  $J_1$ . These two values then undergo the standard vector addition to form a total angular momentum J; the overall state is notated  $(J_1, J_2)_J$ 

<sup>&</sup>lt;sup>12</sup>For <sup>166</sup>Er, the background scattering length  $a_s$  is such that  $\varepsilon_{\rm dd} \approx 1$ , and  $\varepsilon_{\rm dd}$  can be tuned about this value with modest fields between 0 and 3 G [71]. Meanwhile in <sup>168</sup>Er, which was first condensed,  $a_s$  is about double  $a_{\rm dd}$ .



Figure 1.2: Energy level structure of erbium, up to states with energy 25000 cm<sup>-1</sup>. States of even parity are shown in red; those with odd parity are black. The two laser cooling transitions used in the present work, at 401nm and 583nm, are highlighted. This figure is loosely based on those in [41, 81].

[41]. For both the 401nm and the 583nm transitions, one 6s electron is excited to a 6p orbital. The inner electrons couple to an L = 5 triplet state with  $J_1 = 6$ . Meanwhile, the two n = 6 electrons couple to an L = 1 state with  $J_2 = 1$ . For the broad 401nm (~ 30 MHz natural linewidth) transition the spin state remains a singlet (S = 0), while the 583nm transition is an intercombination line in which the spin flips to a triplet state; its semi-forbidden nature<sup>13</sup> lends the transition a narrow linewidth of ~ 200 kHz. Properties of these two excited states, along with the ground state and electronic transitions between, are given in table 1.3.

A cursory glance at figure 1.2 suggests a dire outlook on the suitability of laser cooling at 401nm. Indeed, there are over 10 intermediate states with J = 6 or 8 and even parity into which the excited state can decay, indicating a transition that is far from closed. Luckily, it turns out that few of these decay pathways are strongly E1-coupled [81], and an experimentally measured value for the decay rate to all metastable states [82] gives a branching ratio of  $9.1 \times 10^{-6}$ . Meanwhile, theoretically expected decay rates to two even-parity intermediate states for the

<sup>&</sup>lt;sup>13</sup>The transition is electric dipole (E1)-allowed insofar as the  $\Delta J$  and opposite-parity conditions are satisfied. Traditional E1 transitions also require  $\Delta S = 0$ , which fails here.

State	$[Xe]4f^{12}6s^2(^{3}H_6)$	$[Xe]4f^{12}({}^{3}H_{6})6s6p({}^{1}P_{1}^{o})$	$[Xe]4f^{12}({}^{3}H_{6})6s6p({}^{3}P_{1}^{o})$
$J_1 J_2$ -coupling	J = 6	$(6,1)_7^o$	$(6,1)_7^o$
$g$ -factor $g_J$	1.163801(1)	1.160	1.195
energy $(\text{cm}^{-1})$	0	24943.272	17157.307
wavelength	—	$400.91~\mathrm{nm}$	$582.84~\mathrm{nm}$
nat. linewidth	_	29.7(6)  MHz	186(10)  kHz
sat. intensity	—	$60.2 \mathrm{~mW/cm^2}$	$0.12 \mathrm{~mW/cm^2}$
Doppler temp.	—	713 $\mu K$	$4.5 \ \mu K$
uses	g.s.	ZS, TC, Imaging	MOT

**Table 1.3:** Properties of the erbium states and cooling transitions of interest in this thesis. The *g*-factor quantifies the Zeeman energy shift under the presence of an external magnetic field B,  $\Delta E_Z = m_J g_J \mu_B B$ , where  $m_J$  is the magnetic quantum number and  $\mu_B$  is the Bohr magneton. The Doppler temperature is given by  $T_D = \hbar \Gamma/2k_B$ , where  $\Gamma$  is the natural linewidth; it reflects the lowest temperature attainable with the simplest laser cooling techniques (see section 2.1.2). The saturation intensity is  $I_{\text{sat}} = \pi h c \Gamma/3\lambda^3$ , where c is the speed of light and  $\lambda$  is the transition wavelength. Here g.s. stands for ground state, ZS for Zeeman slowing, TC for transverse cooling, and MOT for magneto-optical trap. Numbers taken (with minor rounding corrections) from [41].

583nm transition are on the order of  $10^{-3}$  s<sup>-1</sup> and therefore negligible [81].<sup>14</sup>

Following the precedent set by the Innsbruck Er experiment [41], we use the broad 401nm transition for initial laser cooling and slowing of an Er atomic beam due to its high scattering rate. We operate our magneto-optical trap (MOT) on the narrow 583nm line, however, which yields lower temperatures.<sup>15</sup>

Other current ultracold erbium experiments include those at Innsbruck and Universität Bonn (group of Martin Weitz), both of which have achieved BEC [38, 84]. An Er quantum gas microscope is also being implemented at Harvard University (group of Markus Greiner).<sup>16</sup>

 $<sup>^{14}</sup>$ The branching ratio here is on the order of  $10^{-9}$ ; i.e. one out of every billion scattered photons leads to decay into a metastable state.

<sup>&</sup>lt;sup>15</sup>An alternative approach was used, for instance, in the dysprosium experiment at Stanford, which first loaded a MOT on the broad blue transition before transferring to a second MOT on a particularly narrow line of  $\sim 2$  kHz [37].

<sup>&</sup>lt;sup>16</sup>Experiments on dysprosium, erbium's closely-related (but slightly *more* dipolar) cousin, are more numerous. These include groups at Stanford, Stuttgart (group of Tilman Pfau) and LENS/University of Florence (Massimo Inguscio), all of which have achieved BEC. Further Dy experiments in progress include those at Innsbruck (Ferlaino group, whose 2nd generation experiment has achieved BEC of Er and Dy in the same apparatus), MIT (group of Wolfgang Ketterle), Paris (Jean Dalibard and Sylvain Nascimbene), and Mainz (Patrick Windpassinger).

#### 1.3 Outline of the Thesis

This thesis is structured as follows. Chapter 2 presents the theoretical background for the various laser cooling and trapping techniques towards which the bulk of this thesis has pushed. Chapter 3, meanwhile, describes the design of the experimental apparatus and details of its physical construction. The material presented there comprises the vast majority of the work undertaken over the course of this thesis, and portions of this work are still ongoing. Finally, chapter 4 presents the first measurements of an erbium atomic beam observed in the new apparatus, briefly discusses possible future research projects suitable to the completed machine, and concludes the thesis.

## Chapter 2

## Theory of Laser Cooling and Trapping

In this chapter we describe the theoretical background necessary to understand the principal aim of the work contained in this thesis: namely, to take atoms at or above room temperature ( $\sim 10^2 - 10^3$  K) and cool them to temperatures near  $10^{-6}$  K. Section 2.1 describes the cooling and slowing methods relevant to this work while 2.2 describes the technique used to introduce spatial trapping of an atomic cloud.<sup>1</sup>

#### 2.1 Laser Cooling

Traditionally, the consideration of atom-light interactions involves the discussion of their effect on *internal* atomic states – in particular, radiation resonant with an atomic transition will excite the atom to a higher energy level. Yet it is by exploiting the effect of such a process on *external* momentum states that we are able to slow, and thereby cool, atoms to low temperature.

The intuitive picture of this phenomenon is as follows. When an atom is excited to a higher internal energy state, it absorbs a photon of momentum  $\hbar k$ , where  $k = 2\pi/\lambda$  is the wavenumber of the photon. By momentum conservation, the atom's momentum increases by  $\hbar k$  in the direction of laser propagation. The photon is then spontaneously emitted, applying another momentum kick to the atom in some arbitrary direction. Over time, these momenta from spontaneous emission average to zero, meaning that the only net force is from absorption: the atom is 'pushed' in the direction of laser propagation.

<sup>&</sup>lt;sup>1</sup>The 1997 Nobel Prize in physics was awarded to William Phillips, Steven Chu, and Claude Cohen-Tannoudji in recognition of their contributions to the development of these processes.

This process is quantitatively described for a two-level system in terms of a scattering force,<sup>2</sup>

$$F_{\text{scatt}} = \frac{\hbar k \Gamma}{2} \frac{s}{1 + s + (2\delta/\Gamma)^2} , \qquad (2.1)$$

where  $\Gamma$  is the linewidth of the transition,  $\delta$  is the total detuning of the radiation from resonance, and the saturation parameter,  $s = I/I_{\text{sat}}$ , is the ratio of laser intensity to saturation intensity. Note that the maximal force,  $F_{\text{max}} = \hbar k \Gamma/2$ , corresponds to photon scattering at *half* the natural transition rate – this is because at saturation  $(s \to \infty)$  the population of the excited state is 1/2.

#### 2.1.1 Zeeman Slowing

We begin by considering atomic slowing in one dimension. In practice, such a technique is used to cool hot, collimated atomic beams to velocities at which other cooling and trapping methods become effective, and it is an essential first step in the erbium cooling process described in this work.

The principle of one-dimensional slowing is simple: propagate a resonant laser beam opposite the direction of the atomic beam and let the scattering force decelerate the atoms to near-zero velocity. Due to the Doppler effect, however, the laser frequency in the atomic frame changes as a function of velocity. To maintain the resonance condition necessary for optimal slowing, we introduce a spatially varying magnetic field tuned so that, by virtue of the Zeeman effect, the transition frequency exactly matches the atomic-rest-frame laser frequency. The resulting device is aptly named a Zeeman slower.<sup>3</sup>

The resonance condition in such a system<sup>4</sup> is

$$\underbrace{\omega + kv(x)}_{\text{laser frequency}} = \underbrace{\omega_0 + \frac{\mu' B(x)}{\hbar}}_{\text{transition frequency}} , \qquad (2.2)$$

where  $\omega$  is the laser frequency (in the lab frame),  $\omega_0$  is the unperturbed transition frequency, B(x) and v(x) are the magnetic field and velocity as functions of position along the slower, and  $\mu' = (g_e m_e - g_g m_g)\mu_B$  is the effective magnetic moment along the transition, defined in terms of Landè g-factors g and magnetic quantum numbers m of the ground and excited states, as well as the Bohr magneton  $\mu_B$ . For slowing

 $<sup>^{2}</sup>$ See, for instance, [42] for a derivation of this ubiquitous formula.

<sup>&</sup>lt;sup>3</sup>The following section is based in part on a similar discussion in [41].

<sup>&</sup>lt;sup>4</sup>For an atomic beam propagating in the +x direction and laser beam opposite.

on the 401 nm  $|J, m_J\rangle = |6, -6\rangle \rightarrow |7, -7\rangle$  transition in erbium,  $\mu' = -1.1372$ .

Assuming this condition is satisfied, the atomic motion has constant acceleration (from equation 2.1)

$$a_{\rm ZS} = \frac{F_{\rm scatt}}{m} \equiv \eta a_{\rm max} \le a_{\rm max} , \qquad (2.3)$$

where

$$a_{\max} = \frac{\hbar k \Gamma}{2m} \tag{2.4}$$

is the maximum possible acceleration (with m the mass of the atom) and

$$\eta = \frac{s}{1 + s + (2\delta/\Gamma)^2} \tag{2.5}$$

is the 'security parameter.' Since the detuning is zero given the resonance condition in equation 2.2,  $\eta$  has its maximal value of  $\eta_{\text{max}} = \frac{s}{1+s}$ .

In practice, to account for experimental limitations such as imperfect magnetic coil windings and varying laser intensity, as well as quantum fluctuations in the scattering process, it is best to assume suboptimal scattering and therefore choose a value of  $\eta$  less than  $\eta_{\text{max}}$ . According to equation 2.5, this enforces a constant total 'security' detuning<sup>5</sup>

$$\delta_{\rm sec} = -\frac{\Gamma}{2} \sqrt{\frac{s}{\eta} - s - 1} , \qquad (2.6)$$

and we therefore modify the resonance condition (equation 2.2) to read

$$\Delta + kv(x) - \frac{\mu'}{\hbar}B(x) = \delta_{\text{sec}} , \qquad (2.7)$$

where  $\Delta = \omega - \omega_0$  is the (lab frame) laser detuning.

Taking our constant-acceleration slower to bring atoms to zero velocity over a length  $x_0$ , we use simple kinematics to solve for the atomic velocity as a function of position along the slower:

$$v(x) = v_c \sqrt{1 - \frac{x}{x_0}}$$
, (2.8)

where

$$v_c = \sqrt{2a_{\rm ZS}x_0} = \sqrt{2\eta a_{\rm max}x_0} \tag{2.9}$$

is the capture velocity (i.e. the maximum initial velocity that will be meaningfully slowed by the device).

<sup>&</sup>lt;sup>5</sup>The sign here ensures the that the slower is stable: for negative  $\delta$ , if v(x) is larger than its ideal value at a given x, the detuning increases towards zero (eqn. 2.7), thereby causing the atom to scatter more photons and slow more quickly.

To solve for the magnetic field, we insert equation 2.8 into equation 2.7 and rearrange. The result is a square-root dependence of the magnetic field on position:

$$B(x) = B_b + B_0 \sqrt{1 - \frac{x}{x_0}} , \qquad (2.10)$$

where the bias field  $B_b$  and amplitude  $B_0$  are given by

$$B_b = \frac{\hbar}{\mu'} \left( \Delta - \delta_{\text{sec}} \right) \quad \text{and} \quad B_0 = \frac{\hbar k}{\mu'} v_c \ . \tag{2.11}$$

Evidently we are able to adjust the capture velocity (or, equivalently,  $\eta$ ) by tuning  $B_0$ . The bias field and laser detuning, meanwhile, are mutually interdependent, and their values can be chosen for experimental convenience.<sup>6</sup> For a more detailed treatment of Zeeman slowing theory, see appendix A.

There are three regimes in which the Zeeman slower can be operated (figure 2.1): increasing-field  $(B_b \approx -B_0)$ , decreasing-field  $(B_b \approx 0)$ , and spin-flip (0 <  $|B_b| < |B_0|$ ) [85]. In the first, the magnitude of the field is greatest at the end of the slower, meaning that atoms quickly fall out of resonance with the slowing light after the slower and proceed to the MOT unhindered. At the same time, though, the laser detuning must equal the Doppler shift of the fastest atoms to be slowed (a technically difficult 1.25 GHz for 500 m/s erbium atoms), and the large field at the MOT end of the slower makes reducing its magnitude to zero at the MOT center difficult. Decreasing-field slowers, meanwhile, require small laser detunings and small fields near the MOT, but atoms fail to quickly fall off resonance with the laser past the end of the slower, resulting in additional slowing that can, in a worst-case scenario, reverse the direction of the atoms before they reach the MOT.

Our choice of slower – in the spin-flip configuration – represents a happy medium between the two: the laser detuning is modest and achievable with one double-pass AOM ( $\Delta \approx -2\pi \times 540$  MHz in our setup), the atoms quickly drop off resonance at the end of the slower, and the magnetic field at the MOT end is not so large as to significantly interfere with MOT operation (provided we introduce modest compensation fields). The disappearance of a well-defined quantization axis at the zero-crossing of the magnetic field presents problems in principle, but practically

<sup>&</sup>lt;sup>6</sup>The  $\eta$ -dependent detuning  $\delta_{\text{sec}}$  is typically small compared to  $\Delta$  and is fully determined by  $B_0$ ; i.e. the value of  $B_b$  cannot change  $\eta$ . Altering the relationship between  $B_b$  and  $\Delta$  can be useful, however: for a given  $B_b$ , adjusting  $\Delta$  in equation 2.7 necessitates a compensating 'velocity offset' term  $k\Delta v$  to maintain the equality, and in this way we get a non-zero final velocity (another way to achieve this is to simply have the field drop to zero before  $x = x_0$ ). See section A.2 in the appendix.



Figure 2.1: Magnetic field profiles for various types of Zeeman slowers: (a) increasing-field, (b) spin-flip, and (c) decreasing-field. Notice that here  $B_0$  is negative, corresponding to the fact that  $\mu' < 0$  for the cycling transition  $(|6, -6\rangle \rightarrow |7, -7\rangle)$  used in our slower.

speaking atoms are optically pumped to the desired cycling transition quickly enough upon reestablishment of a field axis that they do not come far off resonance in the interim.

#### 2.1.2 Optical Molasses

While Zeeman slowing allows us to cool and slow an atomic beam along its axis,<sup>7</sup> it does little to alter the transverse profile. In order to also cool the radial dimension(s), we employ a simple technique known as optical molasses. Making use of the scattering force, we simply point pairs of counter-propagating laser beams, of frequency  $\omega$ , along the axes where cooling is desired. In order to ensure that an atom of nonzero velocity preferentially interacts with the photons opposing its motion (rather than those that will tend to speed it up, which exist in equal numbers), we red-detune the laser radiation, so that  $\omega < \omega_0$ . By virtue of the Doppler effect, a moving atom will then preferentially scatter photons that tend to slow it down – in an ensemble of many atoms, this leads to a narrowing of the velocity distribution, i.e. a cooling of the cloud.<sup>8</sup>

 $<sup>^{7}\</sup>mathrm{ZS}$  not only slows but also cools because atoms of different initial velocities exit the slower with similar final velocities, narrowing the distribution.

<sup>&</sup>lt;sup>8</sup>Note that, while it's easier to think of temperature in an ensemble of atoms (which is Maxwell-Boltzmann distributed at any given time) an individual laser-cooled atom can also be said to have

We make this picture quantitative by once again considering the scattering force acting on the atom.<sup>9</sup> Along each principal axis, this is given by<sup>10</sup>

$$F_{\rm mol}(v) = \hbar k R_{\rm scatt}(\omega - \omega_0 - kv) - \hbar k R_{\rm scatt}(\omega - \omega_0 + kv) , \qquad (2.12)$$

where

$$R_{\text{scatt}}(\delta) = \frac{\Gamma}{2} \frac{s}{1 + s' + (2\delta/\Gamma)^2} , \qquad (2.13)$$

 $s' = 2NI/I_{\text{sat}}$  represents a crude attempt at accounting for saturation effects from all 2N lasers, and  $\delta$  is the total detuning. For  $kv \ll \Gamma$  (slow atomic velocities), expanding eqn. 2.12 gives

$$F_{\rm mol}(v) \approx -\alpha v , \qquad (2.14)$$

with the damping coefficient given by

$$\alpha = \frac{4\hbar k^2}{\Gamma} \left(\frac{\partial R_{\text{scatt}}}{\partial \delta}\right)_{\delta=\Delta} = -4\hbar k^2 s \frac{2\Delta/\Gamma}{\left[1 + s' + (2\Delta/\Gamma)^2\right]^2}$$
(2.15)

For negative detuning, this is a linear damping force similar to that acting on an object moving through a viscous fluid, lending the cooling method its name.

The damping force implies an exponential decrease in the atom's kinetic energy with time constant  $\tau_{cool} = m/2\alpha$ :

$$\left(\frac{dE}{dt}\right)_{\text{cool}} = \frac{d}{dt} \left(\frac{p^2}{2m}\right) = vF_{\text{mol}}(v) = -\alpha v^2 = -\frac{2\alpha}{m}E \qquad (2.16)$$
$$\implies E(t) \propto e^{-t/\tau_{\text{cool}}}$$

This cooling, which would otherwise asymptotically lead to zero temperature, is balanced by heating due to the random-walk character of single-photon absorption and emission events, which each give momentum-space random walks  $\langle p^2 \rangle \propto \hbar^2 k^2 R_{\text{scatt}} t$ . As discussed in [42, 86], these processes lead to heating given in three dimensions

a temperature insofar as its time-averaged velocity distribution is thermal [86].

<sup>&</sup>lt;sup>9</sup>The discussion below roughly follows those of [42, 86].

<sup>&</sup>lt;sup>10</sup>This treatment assumes that the two laser beams act independently, i.e.  $s \ll 1$ . Nonetheless, below we (at least initially) crudely account for saturation effects from the other beams. A more precise treatment requires the consideration of counter-propagating beams as standing waves and atoms as multi-level (rather than two-level) systems – these effects lead to various sub-Doppler cooling processes not considered here [42, 86].

 $by^{11}$ 

$$\left(\frac{dE}{dt}\right)_{\text{heat}} = \frac{1}{2m} \frac{d\langle p^2 \rangle}{dt} = 2\hbar^2 k^2 R_{\text{scatt}}/m \qquad (2.17)$$

Equating the heating (eqn. 2.17) and cooling (eqn. 2.16) rates yields an equilibrium temperature<sup>12</sup>

$$k_B T = mv^2 = \frac{2\hbar^2 k^2 R_{\text{scatt}}}{\alpha} = \frac{\hbar\Gamma}{4} \frac{1 + (2\Delta/\Gamma)^2}{2|\Delta|/\Gamma}$$
(2.18)

where we have assumed  $s \sim s' \ll 1$  (reflecting a necessary condition for treating the cooling beams independently). This is minimized when the detuning is half the linewidth,  $\Delta = -\Gamma/2$ , and gives the *Doppler cooling limit*:

$$k_B T = \frac{\hbar\Gamma}{2} \tag{2.19}$$

The minimum temperature achievable through optical molasses, or Doppler cooling, is therefore equal to half the linewidth of the cooling transition.

#### 2.2 The Magneto-Optical Trap (MOT)

Though optical molasses provides effective cooling of atomic clouds, it acts exclusively in *velocity* space, meaning that atoms are free to spatially diffuse wherever they please. In order to add spatial confinement, we can introduce a position-dependent magnetic field; the result is the aptly named magneto-optical trap (MOT) [21], which has become a ubiquitous workhorse in atomic physics as the starting point for a vast number of experiments.

The fundamental principle of the MOT can be understood through a 1D model consisting of a two-state atom with ground state angular momentum F = 0 and excited angular momentum F' = 1.<sup>13</sup> As for optical molasses, the atom interacts with two counter-propagating, red-detuned lasers. Here, however, the laser propagating in the + (-) x direction has  $\sigma^+ (\sigma^-)$  polarization so as to drive transitions from  $m_F = 0 \rightarrow m_{F'} = +1$  (-1). Finally, we introduce a constant magnetic field gradient B(x) = b'x, which splits the magnetic sublevels in a spatially dependent

<sup>&</sup>lt;sup>11</sup>There should be slightly less heating in two dimensions, given that some photon emission events occur along the third axis and heat the atoms in a direction we can effectively ignore.

<sup>&</sup>lt;sup>12</sup>Using the equipartition theorem with one degree of freedom for the one axis under consideration.

<sup>&</sup>lt;sup>13</sup>This is easily generalized to any system for which F' = F + 1.



Figure 2.2: Principle of MOT operation. Because of the lasers' red detuning and the magnetic field gradient, atoms at positions x < 0 are more likely to absorb light from beam 1 ( $\sigma^+$ -polarized and propagating in the +x direction), while atoms at x > 0 will more probably scatter from beam 2. This leads to a spatially dependent force that tends to confine the atoms near the origin.

manner (they are degenerate at the origin).<sup>14</sup>

As can be seen from figure 2.2, the red-detuning of the cooling lasers, combined with transition selection rules and the spatial dependence of the *B* field, ensures that an atom offset from the origin will always preferentially scatter photons from the laser that tends to push it back to x = 0. Furthermore, the red detuning also typically allows Doppler cooling as described above, so atoms experience feedback in both position *and* velocity space. The MOT is easily extended to three dimensions by adding two more pairs of counterpropagating beams and using a quadrupole magnetic field, which results in gradients  $-2\frac{dB_x}{dx} = -2\frac{dB_y}{dy} = \frac{dB_z}{dz}$  along the three principal axes (see section 3.5.1).

#### 2.2.1 Narrow-Line Trapping

For a typical MOT used in, e.g., cold atom experiments with alkali atoms, the cooling transition has a linewidth of order several MHz and the laser detuning is set such that  $|\Delta| \approx \Gamma$ . This leads to a radiative force that depends linearly on position near the center of the trap, and the behavior of the MOT is well described

<sup>&</sup>lt;sup>14</sup>Here and in figure 2.2 we adopt a lab-frame picture in labeling atomic states and laser polarizations. In the atom frame at positions x < 0, the quantization axis flips to align with  $\vec{B}$ , and the lowest-energy state is still  $m_F = -1$ . However, because the quantization axis is reversed, the laser polarization is also reversed, so that now the +x-propagating light is  $\sigma^-$  polarized, and vice-versa. This choice therefore amounts to a labeling convention, while the underlying physics is unchanged.

by a damped harmonic oscillator (see, for instance, [21, 42, 87]). In contrast to this typical scenario, in the present experiment – motivated by a desire to achieve lower final MOT temperatures<sup>15</sup> and thereby avoid sub-Doppler cooling steps later in the experimental sequence – we choose to trap on the narrow intercombination line of erbium at 583 nm, with natural linewidth  $\Gamma/2\pi = 190$  kHz. Partly because of the fact that our laser linewidth is of order  $\Gamma$  (so that for  $|\Delta| < \Gamma$  some of the laser light will be blue detuned and therefore anti-confining), for such a narrow transition we operate in the regime where  $|\Delta| \gg \Gamma$  throughout the MOT loading and compression process. In this regime, the intuitive picture of MOT operation is entirely distinct from that of more traditional broad-line traps.

We begin by considering the relevant forces in the problem. In most wideline MOTs, the scattering force is large enough to neglect both gravitational and magnetic forces, but in our case we must proceed more carefully. Here, the ratio of scattering to gravitational forces is about 250 (compared with ~  $5 \times 10^4$  for the 401nm transition), and due to the large ground-state magnetic moment of Er, the ratio of scattering to magnetic forces is also small compared with a standard MOT (about 340 in this case). In the regime where  $|\Delta| > \Gamma$  it is important that we take both of these conservative forces into account.

The force on an atom in our trap along the z axis (defined opposite gravity) is given by [41, 83, 88]

$$F_{\text{MOT}}(z, v_z) = F_{\text{scatt}} + F_{\text{grav}} + F_{\text{mag}} = \frac{\hbar k \Gamma}{2} \left[ \frac{s}{1 + s' + 4(\Delta - kv_z - |\mu'\partial_z B_z| z)^2 / \Gamma^2} - \frac{s}{1 + s' + 4(\Delta + kv_z + |\mu'\partial_z B_z| z)^2 / \Gamma^2} \right] - mg - m_g g_g \mu_B \partial_z B_z , \quad (2.20)$$

where  $\partial_z B_z = \partial B_z / \partial z$  is the magnetic field gradient,  $\mu' = (m_e g_e - m_g g_g) \mu_B / \hbar$  is the effective magnetic moment of the cooling transition,<sup>16</sup>  $s' \geq s$  accounts for saturation due to the beams along other axes, and g = 9.8 m/s is the gravitational acceleration. Figure 2.3a shows plots of this force for fixed s = s' = 12 and several values of

<sup>&</sup>lt;sup>15</sup>Recall that the final MOT temperature is roughly given by the Doppler temperature,  $T_D \propto \Gamma$ , which is roughly a factor of 150 smaller (4.6  $\mu$ K vs. 714  $\mu$ K) for the 583 nm transition than for the broad one at 401 nm.

<sup>&</sup>lt;sup>16</sup>We take the absolute value of this and the gradient in equation 2.20 for clarity in establishing the sign of various terms. Note also that here we define  $m_e$  and  $m_g$  as projections along the +z axis, even though at times this is opposite the direction of  $\vec{B}$ .



**Figure 2.3:** (a) Vertical MOT force (eqn. 2.20) as a function of position for  $v_z = 0$ , s = s' = 12,  $dB_z/dz = 3$  G/cm, and two values of the detuning  $|\Delta| > \Gamma$  and  $|\Delta| = \Gamma$ . In the large-detuning regime, the atoms effectively fly in free space surrounded by two hard walls. Gravity and the magnetic force break the symmetry, however, leading to a small negative force in the free-space regime (not visible on this scale) that causes a 'gravitational sag' (b). Part (b) reproduced, with small modifications, from [41].

 $\Delta$ . The atoms preferentially interact with the upward-pointing beam because of the non-negligible gravitational potential, which drags them downwards. For  $\sigma^$ polarization (corresponding to a *negative* vertical gradient, in which the field points down for z > 0 and up for z < 0), the atoms are therefore pumped into the  $m_J = -6$ ground state, and the magnetic force becomes anti-confining, summing with the gravitational force to push the atoms further down.

Because of these combined gravitational and magnetic forces, the atoms sit below the x-y plane, on an equipotential of constant Zeeman shift where  $F_{\text{MOT}} = 0$ , i.e. where the scattering force from the upward-pointing beam balances the combined gravitational and magnetic forces (figure 2.3b). Ignoring the downward-pointing beam (which is far detuned for  $v_z \approx 0$  and z < 0), we can solve equation 2.20 for the equilibrium position  $z_0$ , given by

$$\frac{\Delta - |\mu'\partial_z B_z| z_0}{\Gamma} = -\frac{\sqrt{Rs - s' - 1}}{2} \tag{2.21}$$

where

$$R = \frac{\hbar k\Gamma}{2\left(mg - m_g g_g \mu_B \left|\partial_z B_z\right|\right)} \tag{2.22}$$

is the ratio of the maximum scattering force to the combined magnetic and gravitational forces.

To determine the equilibrium temperature of the MOT we take the same ap-



Figure 2.4: MOT temperature calculated from equation 2.25 for the narrow 583 nm line as a function of the saturation parameter s = s'. A small gradient dependence arises from the magnetic force term in R (equation 2.22); this results in a horizontal offset of the temperature minimum (inset) but no notable temperature increase. In all cases for saturation parameters around s = 0.01 we achieve temperatures approximately equal to  $T_D = 4.46 \ \mu \text{K}.$ 

proach as in section 2.1.2, beginning by finding the damping coefficient  $\alpha_z$  about this equilibrium point ( $z = z_0, v_z \approx 0$ ). To first order, expanding  $F_{\text{MOT}} \approx -\alpha_z v_z$ gives [88]

$$\alpha_z = -2\hbar k^2 \frac{\sqrt{Rs - s' - 1}}{R^2 s} \tag{2.23}$$

The heating rate is described by a single-beam random walk due to absorption and scattering events, with diffusion coefficient  $[86, 88]^{17}$ 

$$D_p = \frac{1}{2} \frac{d\langle p^2 \rangle}{dt} = \hbar^2 k^2 R_{\text{scatt}} = \frac{\hbar^2 k^2 \Gamma}{2R}$$
(2.24)

where in the last equality we have substituted equation 2.21 into  $R_{\text{scatt}}$  (eqn. 2.13). Balancing damping and diffusion gives the (z-axis) equilibrium temperature [88]

$$T_z = -\frac{D_p}{k_B \alpha_z} = \frac{\hbar \Gamma \sqrt{s}}{2k_B} \frac{R}{2\sqrt{R - s'/s - 1/s}} \approx \left(\frac{\hbar \Gamma_s}{2k_B}\right) N_R = T_D N_R \tag{2.25}$$

where  $\Gamma_s = \Gamma \sqrt{1+s}$  is the power-broadened linewidth and we have approximated  $\sqrt{1+s} = \sqrt{s}$  for large s.  $T_D$  is the Doppler temperature (equation 2.19) and  $N_R$  is a numerical factor of order 1 (figure 2.4).

<sup>&</sup>lt;sup>17</sup>This describes purely 1D diffusion, which is a reasonable approximation for atoms sitting near equilibrium, since they hardly interact with the x-y beams.

As discussed in [89], because the atoms sit in a nonzero bias field at equilibrium due to their z-axis offset from the origin, we have to treat the radial (x-y plane) beam polarizations in terms of their respective projections onto the quantization axis, which is tilted at a small angle  $\alpha \approx (x^2 + y^2)^{1/2}/z_0 = r/z_0$  from the z axis. Whereas in the absence of a bias field the x-y behavior of the atoms is characteristic of free space motion bounded by hard walls (located at  $x_0 = \Delta/\mu' \partial_x B_x$ ), its introduction means that each radial laser has a  $\sigma^-$  polarized component (defined with respect to the tilted quantization axis) and is no longer far-detuned even at small r.<sup>18</sup> As it turns out, this component is always largest for the laser tending to push the atom back towards the origin, and the result is a harmonic potential in the radial directions. Moreover, the radial equilibrium temperature is actually (though not obviously) equal to the vertical temperature (equation 2.25) [89].

While low Doppler temperatures and straightforward spin polarization are definitive advantages of the narrow MOT, these boons come with the consequence of reduced capture velocity  $v_c$ . Indeed, an upper bound on  $v_c$  is set by the maximum scattering acceleration  $a_{\text{max}} = \hbar k \Gamma/2m$  so that  $v_{c, \text{max}} = \sqrt{2a_{\text{max}}l}$ , where l is the interaction distance roughly given by the MOT beam diameter [89].<sup>19</sup> For 28.5 mm beam diameters, this gives about 180 m/s for the 401 nm transition but only 11.6 m/s for the 583 nm line. Nonetheless, velocities under 10 m/s are achievable with a well-designed Zeeman slower, so this low capture velocity is not particularly problematic.

We finally turn to the realization of this MOT in the laboratory, in particular by determining the MOT parameters we expect to use and the capture velocities we expect to attain. In the first phase of operation we aim to simultaneously optimize the capture velocity of the MOT and allow a gravitational sag large enough so that the atoms sit below the Zeeman slowing beam waist of  $w_{\rm ZS} \sim 5 \text{ mm.}^{20}$  This is given by the condition  $|z_0| \leq w_{\rm ZS}$  with  $z_0$  from equation 2.21. Further, it is shown in

<sup>&</sup>lt;sup>18</sup>In the absence of a bias field the quantization axis is defined by  $\vec{r}$  for any radial offset  $r \neq 0$  from equilibrium, and the small *B* value at small *r* means that the Zeeman shift is small compared to the laser detuning; i.e. the total detuning is large. With the bias field, however, the Zeeman shift at position  $(r, z_0)$  for small *r* is approximately equal to its value at equilibrium  $(0, z_0)$ ; therefore the total detuning is relatively small and the atoms immediately interact with the  $\sigma^-$  component of the radial lasers.

<sup>&</sup>lt;sup>19</sup>One should be careful to note that this maximum capture velocity is never achievable with a MOT: as we have seen above, to maintain constant, maximum acceleration over the whole slowing process requires a square-root B profile, whereas a MOT is bound to a linear profile.

<sup>&</sup>lt;sup>20</sup>Corresponding to a 10 mm  $1/e^2$  beam diameter.

appendix A.3 that in order to capture any atoms we need a minimum radial gradient

$$\frac{dB_x}{dx} \gtrsim \frac{2\left|\Delta\right|}{\left|\mu'\right|l} \tag{2.26}$$

where l is again the interaction length, which for sufficiently large s is given by approximately twice the  $1/e^2$  MOT beam diameter.<sup>21</sup> Combining these two conditions, for realistic parameters of s = 12, beam diameter  $2w_0 = 28.5$  mm and  $l = 1.8 \times (2w_0)$ ,<sup>22</sup> this gives a minimum detuning  $\Delta = -6.5$  MHz =  $-35\Gamma$  and corresponding vertical gradient of  $-2\frac{dB_x}{dx} = -2.6$  G/cm. Figure 2.5a shows numerical simulations of the capture behavior<sup>23</sup> for parameters of  $\Delta = -7$  MHz and  $\frac{dB_z}{dz} = -3.2$  G/cm, which allow some leeway over the absolute minimum results and still give  $z_0 = -5.2$  mm  $< -w_{\rm ZS}$ .<sup>24</sup> For this particular configuration we expect a capture velocity of 7.2 m/s, though in other configurations we can achieve  $v_c$  above 9 m/s (figure 2.5b-d).

While we will load the MOT in this far-detuned, high-s regime, from eqn. 2.25 and figure 2.4 it is clear that minimizing the temperature requires very low  $s \sim 0.01$ . Therefore, in a later stage we will ramp down the intensity, gradient, and detuning (following [41]) to move the atoms up towards z = 0 and cool them towards the Doppler temperature. It is also likely necessary to apply a small bias field, which offsets the quadrupole minimum and pushes the atoms the rest of the way up to the x-y plane.

<sup>&</sup>lt;sup>21</sup>As the atom moves forward the saturation intensity decreases but the atom also slows, and the resulting increased interaction time per position interval dx compensates for the decrease in maximum deceleration  $\eta_{\text{MOT}} = s/(1+s) = \frac{s_0 \exp(-2x^2/w_0^2)}{1+s_0 \exp(-2x^2/w_0^2)}$ . Accordingly, the slowing continues effectively even to very small  $\eta_{\text{MOT}}$  on the order of  $1/e^4$  of its maximum value; for  $s_0 = 12$  this occurs at  $x = 1.8w_0$ , where  $w_0$  is the MOT beam waist.

<sup>&</sup>lt;sup>22</sup>See previous note.

 $<sup>^{23}</sup>$ Simply obtained by numerically solving the *x*-axis equivalent of equation 2.20, which omits the gravity term.

 $<sup>^{24}</sup>$ The elevated gradient also leads to a slightly higher capture velocity – see appendix A.3 for a simple argument why this might be the case.



Figure 2.5: Numerically simulated narrow-line MOT capture behavior. (a) Atom trajectories for various initial velocities in a MOT with  $\Delta = -7$  MHz,  $dB_z/dz = -3.2$  G/cm and s = s' = 12. All captured atoms fall onto the same trajectory, which roughly follows the linear behavior of B(x). Gray arrows represent the vector field in phase space corresponding to  $F_{MOT}(x, v_x)$ . (b) Dependence of the capture velocity on radial gradient  $dB_x/dx = -0.5 dB_z/dz$  and laser detuning  $\Delta$  for s = s' = 12. The dashed line roughly corresponds to the critical gradient condition in equation 2.26 with  $l = 1.5 \times (2w_0)$  and a small vertical offset, reflecting the fact that equation 2.26 is only approximately correct. (c) Cut of the parameter space in part (b), showing the maximum capture velocity as a function of the radial gradient, with the laser detuning that maximizes  $v_c$ . We see that the maximum capture velocity for this MOT at s = s' = 12 is 9.1 m/s, achieved at  $dB_z/dz = -4.4$  G/cm and  $\Delta = -10.8$  MHz, corresponding to  $z_0 = -8.6$  mm. (d) Maximum MOT capture velocity as a function of the saturation intensity, calculated by finding the highest point in the gradient-detuning parameter space of part (b) for each value of s. The curve, somewhat expectedly, goes roughly like s/(1+s), meaning that higher saturation intensities yield diminishing returns.

## Chapter 3

## Experimental Design and Implementation

In this chapter, we summarize the design and construction of the erbium experiment. Section 3.1 gives an overview of the full experimental design, many pieces of which are currently realized and some of which are planned for the near future. Section 3.2 outlines the system used to achieve high (HV) and ultra-high (UHV) vacuum in several chambers needed for the cooling, trapping, and manipulation of ultracold Er. Section 3.3 describes the lasers and optics used for Er cooling and trapping. Section 3.4 details the implementation of a Zeeman slower that decelerates Er atoms, in one dimension, from  $\sim 500$  m/s to near zero velocity. Finally, section 3.5 discusses the design and implementation of the various magnetic fields required at and around the MOT chamber.

#### 3.1 Experimental Overview

At the highest level, our erbium experiment is designed following the successful approach taken at Innsbruck [41], with several key modifications tailored to our specific scientific goals. The starting point is an atomic beam of Er atoms, created by heating erbium metal to  $\sim 1200^{\circ}$ C in order to overcome low vapor pressures at room temperature.<sup>1</sup> The beam is mechanically collimated through a series of apertures and then radially cooled in a 2D molasses before passing through a Zeeman slower (section 3.4) for capture in a MOT. While the 2D molasses (or "transversal cooling") and Zeeman slowing are run on the broad 401 nm transition for maximum cooling

 $<sup>^1{\</sup>rm The}$  melting point of Er is  $1522^{\rm \circ}{\rm C}$  , which is suggestive of the impossibility in achieving a substantial gaseous sample near room temperature.



Figure 3.1: Overview of the main experimental apparatus, with important sections labeled and discussed in more detail below. The two main vacuum sections are labeled HV and UHV for (ultra-)high vacuum. TC stands for transverse cooling, ZS for Zeeman slower, and MOT for magneto-optical trap. All pictured components have been assembled with the exception of the science cell, which will be added in several months' time.

efficiency, the MOT is operated on the narrow, 190 kHz-natural-linewidth transition at 583 nm in order to achieve low temperatures around 10  $\mu$ K. This eliminates the need for sub-Doppler cooling techniques and allows us to load directly into an optical dipole trap (ODT) formed from high-intensity red-detuned light at 1030 nm.<sup>2</sup> Using focus-tunable lenses [90], we will be able to spatially match the size of the ODT beam to the MOT for optimal loading and then translate the ODT roughly 30 cm from the MOT chamber to an all-glass science cell [91]. Here we will cool the atoms to degeneracy via forced evaporation; the all-glass chamber will allow sufficient optical access for the implementation of arbitrary (typically uniform) potentials and optical lattices. Finally, we have allowed room to eventually connect a 2D MOT to the main MOT chamber, which will be used to introduce an alkali species for, e.g., impurity studies in dipolar gases. Figure 3.1 gives an overview of the complete apparatus.

#### 3.2 Vacuum System

The vacuum system contains two main sections: a high-vacuum (HV,  $\sim 10^{-9}$  mbar) side containing the source oven and a radial pre-cooling stage, and an ultra-high vacuum (UHV,  $\sim 10^{-11}$  mbar) section which includes the MOT chamber and, even-tually, a science cell.<sup>3</sup>

<sup>&</sup>lt;sup>2</sup>We choose to trap at 1030nm rather than in the more common 1064-1070nm range because of an Er transition at 1069.5nm that could lead to unwanted heating in the trap, particularly in the fermionic <sup>167</sup>Er isotope (this phenomenon is observed in [41], though no explanation is offered).

<sup>&</sup>lt;sup>3</sup>The vacuum system was designed and constructed almost exclusively by Milan Krstajic [92].

On the HV side, we start with erbium metal in a tantalum crucible which is heated to 1100-1200°C in a high-temperature oven.<sup>4</sup> The atoms exit through a forward-facing 3mm aperture and are mechanically collimated by another 8mm aperture slightly downstream. They then pass through a cube with four 35mm viewports through which we can apply 2D molasses/transversal cooling light. After the TC stage there is a six-way cross with three viewports which can be used for beam characterization. A valve separates the HV section from the remainder of the vacuum system should it be necessary to bring one side up to room pressure (e.g. to load more erbium into the oven).

Between the HV and the UHV sections is a 56cm-long, 8mm-diameter tube, around which we wind and water-cool the Zeeman slower coils (section 3.4). One can think of vacuum systems in analogy with electrical circuits: pressures are voltages, leaks and out-gassing are current sources, pumps are ground, and intermediate components have different conductivities. The long, narrow 'differential pumping' tube in our system has a very low conductivity and accordingly allows a pressure drop of up to two orders of magnitude between the HV and UHV sections.

The centerpiece of the UHV section is the MOT chamber, which has six 40mm viewports (two in the vertical direction and four in the radial plane) for the MOT beams, as well as several 16mm viewports for imaging and other optical access (e.g. for the optical transport beam that will later be implemented). It is machined from non-magnetic 316LN stainless steel,<sup>5</sup> and an indentation is made around both vertical viewports to allow the installation of coils near the atoms. A 16mm-diameter,  $\sim$  30cm-long tube connects the MOT chamber to the (planned) science cell. Another arm of the chamber is currently valved off but will eventually allow the introduction of a 2D MOT with an alkali species. A final arm, opposite the differential pumping/Zeeman slowing tube, contains an aluminum mirror which is used to align the ZS beam into the chamber. This is used in lieu of a window with direct access, which would quickly be coated with Er metal and rendered opaque. It has been shown in similar experiments that in-vacuum mirrors coated with metal exhibit only slightly diminished reflectivity [41, 93].

We maintain vacuum using three combined NEG/ion pumps. The fastest pump<sup>6</sup> is placed adjacent to the MOT chamber, while two weaker pumps<sup>7</sup> are placed in the HV section and between the MOT chamber and science cell, respectively. We expect

<sup>&</sup>lt;sup>4</sup>Dual Filament Cell, model DFC-40-10-WK-2B-SHE, from Createc.

<sup>&</sup>lt;sup>5</sup>Nominal relative permeability < 1.005.

<sup>&</sup>lt;sup>6</sup>Nextorr D 300-5, from SAES Getters.

<sup>&</sup>lt;sup>7</sup>Nextorr D 100-5, from SAES Getters.

with this system to achieve pressures of  $< 3 \times 10^{-11}$  mbar in the MOT chamber and  $< 8 \times 10^{-12}$  mbar in the science cell. See [92] for more details on the design and assembly of the vacuum system.

#### 3.3 Optical System

The optical systems for Er laser cooling must accomplish three principal tasks for light at each of the cooling wavelengths. The first is to stabilize the laser frequency, with respect to both short-time fluctuations and long-term drifts, at a known position relative to the transition frequency.<sup>8</sup> The second is to precisely tune the frequency of individual beams to achieve the detunings required for cooling and slowing. The third is to distribute cooling light to the vacuum chamber and sculpt the beams' spatial profile prior to interaction with the atoms. Sections 3.3.1 and 3.3.2 describe the laser and optical setups for 401 nm and 583 nm light, respectively, while section 3.3.3 describes the the frequency stabilization (laser locking) technique used for both wavelengths.

#### 3.3.1 401 nm Optical Setup

The 401 nm light is derived from an M-Squared laser system consisting of a titaniumsapphire (Ti:sapph) crystal and frequency doubling cavity. The narrow-linewidth Ti:sapphire laser<sup>9</sup> is pumped by 15W of 532 nm light from a diode-pumped solid state laser<sup>10</sup> and produces approximately 5W at 802 nm. This is then sent to a frequency-doubling cavity<sup>11</sup> to deliver approximately 2W of 401 nm light.

For 401 nm laser cooling and imaging we need light resonant with the erbium transition for spectroscopy/locking and absorption imaging, light detuned by approximately -10 MHz ( $\approx \Gamma/3$ ) for transverse cooling/2D molasses (TC), and light detuned by approximately -540 MHz for Zeeman slowing (ZS). This is achieved with 2 acousto-optic modulators (AOMs)<sup>12</sup> operated at 270 MHz in double-pass configuration. The first double-pass setup shifts the frequency by the ZS detuning,  $-\Delta\nu_{ZS} \approx +540$  MHz (see section 3.4), before sending it to the spectroscopy setup

<sup>9</sup>M-Squared SolsTiS module, consisting of the Ti:sapph crystal and resonant bow-tie cavity.

<sup>&</sup>lt;sup>8</sup>This is typically done by locking directly to the atomic transition, as is the case here, though sometimes a lock is achieved on an ultra-stable cavity mode offset from the transition.

<sup>&</sup>lt;sup>10</sup>Lighthouse Photonics Sprout G-15W.

<sup>&</sup>lt;sup>11</sup>M-Squared ECD-X second-harmonic generation (SHG) cavity and crystal.

<sup>&</sup>lt;sup>12</sup>Gooch & Housego, AOMO 3270-125,  $f_0 = 270$  MHz. These are driven with voltage-controlled oscillators (VCOs) fed into ~ 28 dB amplifiers (Mini-Circuits ZFL-1000VH2B+).

for locking (section 3.3.3). This means that the direct laser output is detuned by  $\Delta\nu_{\rm ZS}$  and can be immediately fiber-coupled and sent to the experiment table for Zeeman slowing. Another branch of laser light is sent through a second double-pass AOM setup and used for both TC and imaging, as these frequencies are similar and we never require simultaneous transverse cooling and atom-cloud imaging. For TC operation, the frequency shift is  $\Delta\nu_{TC} - \Delta\nu_{\rm ZS} \approx +530$  MHz and for imaging it is equal to  $-\Delta\nu_{\rm ZS}$ . Shutters are used to select the appropriate beam path following the AOM setup. Figure 3.2a illustrates this portion of the optics.

After frequency shifting, the various beam paths are coupled into polarizationmaintaining fiber and sent to the experiment table. Because of the small beam diameters (around 3  $\mu$ m) required for single-mode fiber coupling at 401 nm, as well as photodarkening effects/color center formation in optical fiber at UV and near-UV wavelengths [94, 95], special caution has to be taken in fiber selection. In particular, we use standard, low-power Thorlabs fiber<sup>13</sup> for imaging and spectroscopy light (both of which can be limited to under 10 mW), but we use 3 higher-power fibers<sup>14</sup> to separately couple ~ 100 mW for each of the (two) TC beams and (one) ZS beam.

In order to maximize the interaction time of atoms with the transverse cooling (TC) beams we send them through a series of spherical and cylindrical lenses<sup>15</sup> that produce elliptical beams with tunable dimensions in a range around 3mm wide  $\times$  35mm long (the length is limited by the CF40 viewports used for TC). While the TC light polarization is not important from an atomic-energy perspective (all levels are effectively degenerate at this point), we use quarter-waveplates to produce circularly polarized light and therefore avoid the creation of standing waves/intensity minimums in the chamber, which will diminish cooling effectiveness.

For the Zeeman slower we use an f = 100mm lens, placed a distance  $L \gtrsim f$ from the diverging fiber output, that allows us to create a slightly convergent beam with the desired diameter through the slower. We can tune the angle of convergence by adjusting L and change the beam diameter through the ZS by translating the whole setup (on a cage system) with respect to the chamber. A  $\lambda/4$ -waveplate and, optionally, a linear polarizer in the beam path enforce  $\sigma^-$  polarization. Both TC and ZS optical setups are shown in figure 3.3.

<sup>&</sup>lt;sup>13</sup>Thorlabs PM-S405-XP Panda-style fiber with FC/APC connectors.

<sup>&</sup>lt;sup>14</sup>Schäfter+Kirchoff, PMC-E-400Si-2.8-NA011-3-APC.EC-700-P, pure silica core fiber with  $300\mu$ m end caps for increased power handling at the coupling interface. The limiting power-handling factor for these fibers is stimulated Brillouin scattering, which caps power transmission at about 150 mW.

<sup>&</sup>lt;sup>15</sup>This lens system was designed by Milan Krstajic, though I built the optical setup.



Figure 3.2: Distribution and frequency shifting setup for (a) 401 nm and (b) 583 nm laser cooling light. In the above, AOM stands for acousto-optic modulator, TC (V/H) for vertical/horizontal transverse cooling, ZS for Zeeman slower, spect. for spectroscopy/locking setup,  $\lambda/2$  for half waveplate, and  $\lambda/4$  for quarter waveplate. In the key, PBS is polarizing beamsplitter, PX is plano-convex lens, f.c. is fiber coupler, Ap. is aperture/iris, and Shutt. is mechanical shutter. Note that optical path lengths are not necessarily to scale.



Figure 3.3: Basic optical paths of cooling beams used for (a) transverse cooling, (b) Zeeman slowing, and (c) the MOT. For TC, the two spherical lenses can be adjusted to obtain collimated beams with  $1/e^2$  diameter between 2 and 5 mm. The cylindrical lense setup consists of two plano-concave and one plano-convex lens and can be tuned to achieve single-axis magnification between  $8 \times$  and  $12 \times$ . This setup is used for each of the two TC beams. For ZS, the single 100 mm plano-convex lens can be moved to change the beam convergence through the slower, while the fiber coupler and lens can be moved together to change the beam diameter through the ZS. The single 150 mm plano-convex lens for each of the MOT beams collimates the fiber output to 28.5 mm  $1/e^2$  diameter; this setup is repeated six times.
#### 3.3.2 583 nm Optical Setup

The 583 nm light used for the erbium MOT is produced by a Toptica system<sup>16</sup> consisting of a tunable diode laser (DL pro) and tapered amplifier (TA pro) that together can produce  $\sim 2$  W at 1166 nm, followed by a SHG frequency doubling cavity that outputs about 600 mW of 583 nm radiation as presently operated.<sup>17</sup> The beam immediately passes through a  $2/5 \times$  telescope that reduces its size to below that of the AOM apertures used for frequency shifting.<sup>18</sup>

The detunings required for the MOT can be anywhere between 0 and 15 MHz, which are not achievable with a single AOM. Therefore, we use two 80 MHz AOMs<sup>19</sup> in double-pass configuration to first shift the light away from resonance and then bring it nearly all the way back. This is accomplished in two 'arms' of the optical setup, in a manner analogous to that used for the TC light at 401 nm. In the first arm, we shift the light used for spectroscopy and locking by  $\Delta \nu_{\rm spect} \approx +160$  MHz before fiber coupling<sup>20</sup> and sending it to the spectroscopy setup. The second branch is shifted via another double-pass AOM setup by  $\Delta \nu_{\rm spect} - |\Delta \nu_{\rm MOT}| \approx 155$  MHz (e.g. for a -5 MHz MOT detuning). This is then coupled into a fiber and sent to the experiment table to be split into six beams for MOT operation. Figure 3.2b shows this portion of the 583 nm optics.

The MOT beams are distributed by way of a homemade fiber cluster, assembled from various off-the-shelf beamsplitters and waveplates,<sup>21</sup> that takes one fiber input and couples it into six fiber outputs with tunable power ratios. Each of these fibers is sent near the MOT chamber, circularly polarized with a quarter-waveplate, collimated with a 150 mm lens to a 28.5mm diameter, and aligned through the chamber with 2" mirrors, as shown in figure 3.3.

#### 3.3.3 Spectroscopy and Frequency Stabilization

For both cooling lasers we choose to frequency lock directly to an erbium reference using spectroscopic methods. This has the obvious advantage of guaranteeing a laser

<sup>&</sup>lt;sup>16</sup>Toptica TA-SHG Pro

<sup>&</sup>lt;sup>17</sup>The laser is designed to output over 1 W at higher TA current.

<sup>&</sup>lt;sup>18</sup>The result is a  $1/e^2$ -diameter of approximately 1mm, sufficiently below the AOM active aperture size of  $1.5 \times 2.5$  mm.

 $<sup>^{19}</sup>$ Gooch & Housego, AOMO 3080-125,  $f_0 = 80$  MHz, driven by VCOs and  $\sim 24$  dB amplifiers (Mini-Circuits ZHL-3A-S+).

 $<sup>^{20}\</sup>mathrm{All}$  583 nm fibers are from Thorlabs: PM460-HP Panda-style polarization maintaining fiber with FC/APC end connections.

<sup>&</sup>lt;sup>21</sup>This was designed and built from Thorlabs components by Milan Krstajic.

frequency directly on resonance with the cooling transition, though it also typically precludes the sort of ultra-stable (< 1 kHz) lock achievable with techniques involving ultra-low-expansion (ULE) cavities. In comparable experiments with Er and Dy, the narrow MOT transition is almost universally locked to a cavity, and it was unclear *a priori* whether we would be able to achieve locking stability below the linewidth of the 583 nm transition.

Due to the low vapor pressure of Er at room temperature, we are unable to use a conventional vapor cell and instead use a see-through hollow cathode lamp  $(HCL)^{22}$  as the atomic reference. The lamp consists of a sealed glass tube filled with argon that contains a hollow cylindrical cathode, covered with Er metal, as well two anodes, spatially separated from the cathode at either end. We apply a potential of ~ 120 V between the cathode and one of the anodes; the positively charged anode ionizes the background Ar gas nearby, and the ions accelerate towards the cathode, sputtering Er atoms off the surface at impact. This produces a gas of Er atoms in the middle of the cathode tube, from which we can obtain laser absorption signals.

We opt for a Doppler-free spectroscopy technique known as modulation transfer spectroscopy (MTS) [96, 97]; its principal advantages over standard saturated absorption spectroscopy are (1) the production of dispersive lineshapes with zerocrossings on resonance (useful for locking) and (2) zero background offset.<sup>23</sup> It begins as a standard pump-probe technique, as in saturated absorption, but with the counter-propagating pump beam phase-modulated with an electro-optic modulator (EOM) at a frequency  $\omega_m$  of order  $\Gamma$  (figure 3.4). We detect the unmodulated probe beam and demodulate the signal at  $\omega_m$ . Away from resonance the probe is unaffected by the modulated pump and so there is zero demodulated signal; near a resonance,<sup>24</sup> however, the nonlinearity of the atomic medium induces a four-wave mixing process in which the pump modulation is transferred (hence the name) onto the probe [96, 97]. The result is a zero-background signal flat everywhere except within the homogeneous linewidth of a transition, where a dispersively shaped feature with zero crossing on resonance appears.

Figure 3.5 shows a representative modulation transfer signal collected on the 401 nm transition. By locking the Ti:Sapph cavity, via a piezoelectric transducer (PZT)

<sup>&</sup>lt;sup>22</sup>From Heraeus.

<sup>&</sup>lt;sup>23</sup>We note that the zero background also improves upon alternative techniques with dispersive error signals, including frequency modulation spectroscopy (FMS) and polarization spectroscopy. Nonetheless FMS, particularly when combined with MTS as in [98], can yield better short-term linewidth stability, while polarization spectroscopy does not require an EOM and has been successfully demonstrated in Er [99].

<sup>&</sup>lt;sup>24</sup>i.e. within the homogeneous linewidth of the transition



**Figure 3.4:** Modulation transfer spectroscopy setup for 401 nm and 583 nm cooling light using the same hollow cathode lamp (HCL). Here EOM stands for electro-optic modulator, PD is photodetector, LPD is long-pass dichroic mirror (which transmits the longer wavelength), and SPD is short-pass dichroic mirror.

on one of the cavity mirrors, to an MTS feature with a simple PI loop we are able to achieve short-term locking stability of a few hundred kHz RMS<sup>25</sup> and no clear long-term drifts. Locking of the MOT laser to the much weaker 583 nm transition was demonstrated by a Part III student this year, and self-heterodyne linewidth measurements [101] suggested a stability well below the 190 kHz transition linewidth [100].<sup>26</sup> Because of an HCL supply issue, we lock both lasers to the same lamp as shown in figure 3.4.

## **3.4** Zeeman Slower

We use a Zeeman slower (ZS) to cool and slow Er in one dimension, with the aim of taking atoms leaving the Er oven with velocities around 500 m/s and decelerating them to the MOT capture velocity of <10 m/s. The design roughly follows that of the Innsbruck Er apparatus [41] and is further motivated by detailed discussions of ZS design and construction in, e.g., [103–107].

 $<sup>^{25}</sup>$ This is measured by looking at the RMS voltage fluctuations on the locked MTS signal; though this method is susceptible to inaccuracy due to electronic noise, measurements on the 583 nm lock using an independent method [100] suggest that this does not underestimate the laser stability.

<sup>&</sup>lt;sup>26</sup>Measurements from the locked error signal, however, had the locked linewidth closer to 300 kHz compared with < 50 kHz via the self-heterodyne technique.

<sup>&</sup>lt;sup>27</sup>Note that these features *do not* correspond to the isotopic abundances – most notably, Er-170 is in reality half as abundant as <sup>168</sup>Er and <sup>166</sup>Er. We attribute this in part to the more abundant isotopes saturating at lower vapor densities than the less abundant ones and have verified that at lower currents the amplitudes begin to reflect natural isotopic abundances. A similar observation is discussed in [99].



Figure 3.5: Modulation transfer spectroscopy and laser frequency stabilization for the 401 nm transition. (a) Modulation transfer signal for ~ 1.4 mW in each counterpropagating beam, an HCL current of 6 mA, RF frequency  $f_0 = 25.9$  MHz, and RF modulation depth  $\beta \approx 2$  rad. The bosonic isotopes give clear, isolated features,<sup>27</sup> and the hyperfine structure of the fermionic <sup>167</sup>Er isotope leads to smaller features labeled with arrows above (these all originate from  $F \rightarrow F + 1$  transitions and are discussed in [102]). (b) Free-running and locked error signals for the <sup>168</sup>Er MTS lock. The locked RMS linewidth is under 200 kHz for <sup>168</sup>Er and under 100 kHz for <sup>166</sup>Er over timescales of several seconds as measured by fluctuations in the electronic signal.

#### 3.4.1 Design of the Slower

Designing a ZS requires that decisions be made regarding a large number of parameters (some of which are interdependent), including slower type (increasing-field, decreasing-field, spin-flip), capture velocity, deceleration parameter  $\eta$ , slower length, laser detuning, etc. The first decision we made, as noted in section 2.1.1 above, was to use a spin-flip slower because of its advantageous properties regarding laser detuning and *B*-field behavior near the MOT.

Many of the remaining parameters can be chosen based on the desired capture velocity of the ZS. This depends on a knowledge of the axial velocity profile of atoms exiting the oven, which is given by the speed distribution for an atomic beam,  $f(v) \propto v^3 \exp(-mv^2/2k_BT)$ . The peak of this distribution is at  $v_{\text{max}} = \left(\frac{3k_BT}{m}\right)^{1/2}$ , which is around 470 m/s for a typical temperature of 1200°C.<sup>28</sup> If we hope to capture all atoms up to this most probable velocity and choose a reasonable deceleration parameter of  $\eta = 0.5$ , which gives good design security without horribly compromising capture velocity,<sup>29</sup> this implies a slower length of (from equation

<sup>&</sup>lt;sup>28</sup>Meanwhile, at  $T = 1100^{\circ}$ C , integrating the velocity distribution up to  $v_c = 470$  m/s gives 50% atom capture; at 1200°C , 45% of atoms are captured.

<sup>&</sup>lt;sup>29</sup>Assuming a saturation parameter s = 5 (corresponding to about 85 mW of power for a 6mm-diameter beam),  $\delta_{sec}$  (equation 2.6) is  $\Gamma$ , allowing a reasonable amount of headroom for

2.9)  $x_0 = 470^2/(2\eta a_{\text{max}}) = 40$  cm. Further increasing the length of the slower yields diminishing returns, as higher velocities become less and less probable per the Maxwell-Boltzmannian velocity distribution. More importantly, increasing the slower length diminishes the range of atom solid angles out of the oven which pass through the tube without hitting the walls; this is particularly apparent in our slower, with its differential-pumping-constrained ZS tube diameter of 8 mm.

We choose the laser detuning  $\Delta$  and bias field offset  $B_b$  primarily with the intention of limiting the magnitude of the magnetic field at either end of the slower. This is minimized for  $B_b = -B_0/2$ , which for  $\eta = 0.5$  gives  $\Delta = -2\pi \times 580$  MHz if we take  $\delta_{sec} = 0$  (eqn. 2.11). In practice we operate with a detuning near  $\Delta/2\pi = -540$ MHz. This is achievable with a double-pass AOM at 270 MHz and decreases the field at the end of the slower as compared with the exact  $B_b = -B_0/2$  case, which has the advantage of reducing the residual field at the MOT.

### 3.4.2 Physical Implementation

We use five independently controllable coils to implement the magnetic field profile of equation 2.10 for our design values of  $x_0 = 400$  mm and  $\Delta/2\pi \approx -540$  MHz (figure 3.6). These are meant, as best as possible, to afford direct control over the field amplitude and bias,  $B_0$  and  $B_b$  (coils 1-3), as well as to reduce the residual field and gradient to zero at the MOT (coils 4-5). Coils 1 and 2 (the "profile" coils) contain a variable number of windings meant to capture the square-root-like, position-varying part of the field; coil 3 (the "bias" coil) is a solenoid that provides constant field offsets; and coils 4-5 are rectangular with 6 × 4 and 6 × 2 windings, respectively, placed symmetrically on opposite sides of the MOT chamber.

Coil 1 captures the negative-B section of the slower, and coil 2 captures the positive-B section after the spin flip. Under ordinary circumstances these are run at the same current, though it is possible to, for instance, run a lower current through coil 2 so as to increase the final velocity of the atoms (see section A.2 of the appendix). The coils are designed so that the maximum field amplitude is the same at either end of the slower – this allows for lower currents to be run, but also leads to a portion of the bias term in equation 2.10 being intrinsically built into the profile coils.<sup>30</sup> Accordingly, we have to first set the current through the profile coils

imperfections in the slower (see appendix A for further discussion of security detunings).

<sup>&</sup>lt;sup>30</sup>In order for  $B_b$  and  $B_0$  to be truly independent, the profile coils would need to run from  $-B_0$  to 0, and the bias field be large and positive. This would result in the inefficient circumstance of large positive and negative currents cancelling to produce small fields.



Figure 3.6: Geometrical coil configuration (a) and simulated magnetic field (b) for the Zeeman slower used in our experiment. Coils 1 and 2 are the profile coils, used to create the negative-*B* and positive-*B* sections, respectively, of the spatially varying field component. Coil 3 is a single-layer solenoid used to create the required bias field offset. Coils 4 and 5 are used to compensate the field and gradient to approximately zero at the location of the MOT (black dot above); the field from coil 5 is small enough not to be visible on this scale. The plot above is for ZS parameters  $\eta = 0.5$  and  $\Delta/2\pi = -540$  MHz, and the small ripples on the total field, caused by discrete jumps in winding number, are on the level of  $\pm 2$  G.

to determine  $\eta$  and only then adjust the bias current to match the laser detuning. Any subsequent change to the profile currents will alter the  $B_b$  term of the field (as well as  $B_0$ ), necessitating further bias current adjustments.

The physical ZS possesses cylindrical symmetry and is wound on a custom-built vacuum part made of 316L stainless steel. It contains an inner, 8mm-diameter tube under vacuum through which the atoms travel and whose size is constrained by the differential-pumping requirement between the HV source chamber and UHV MOT chamber. Surrounding this is a 20mm outer-diameter tube upon which the coils are wound; water flows through the space between the two tubes (which is sealed off, with 1/4" Swagelok connectors on either end) to cool the slower from the inside.<sup>31</sup> The outside of the cooling tube is covered with Kapton tape to electrically insulate the vacuum chamber in the event of wire insulation failure, and we machined endcaps fitted on either end of the tube to hold the winding wires in place at the start and end of the ZS.

We wound the ZS with 1mm-diameter enameled copper wire,<sup>32</sup> chosen because it is small enough to achieve the necessary spatial resolution of the magnetic field profile without dissipating too much heat (the resistance scales inversely with the wire's cross-sectional area). While the nominal ZS length is  $x_0 = 40$  cm, the physical slower spans 42 cm; the extra 2 cm at the end allow us to more closely match the rapidly increasing field profile there. The bias coil (coil 3) consists of a single layer (approximately 420 turns) wound directly on the cooling tube and spanning the entire length of the slower. On top of this, we wind 12 (9) layers of wire for coil 1 (2) to achieve the spatially varying profile; for each layer, we start from the beginning (end) of the slower and wind inwards a predetermined distance. Following [105], we terminate each winding in the middle of the slower rather than winding a subsequent layer back in the opposite direction; while this requires cumbersome connections between the end of one layer and the start of the next, winding backwards leads to messy coils because of the opposite handedness of successive layers (meanwhile, when all windings have the same pitch, the wires in one layer nest in the inter-wire

<sup>&</sup>lt;sup>31</sup>We choose water cooling via the tube rather than with hollow wires because such hollow magnet wire tends to be large ( $\sim 4 \text{ mm}$  diameter – or width in the case of rectangular wires), which doesn't allow enough spatial resolution to achieve a satisfactorily smooth profile over our (relatively) short ZS. Admittedly, the "center of mass" of our design is at a larger radius than that of a wire-cooled slower (in that case the water cooling is evenly distributed rather than concentrated at the center), though this larger average radius only modestly increases the current density through (and power consumption of) the slower.

 $<sup>^{32}</sup>$ Coated with Polyester 200/polyamide-imide insulation for a high temperature index of 212°C ; this is motivated by the need to bake the chamber with the ZS already in place.

	T	/ -		\ \								
	Layer/Length (mm)											
	1	2	3	4	5	6	7	8	9	10	11	12
Bias	420											
Profile 1 (negative)	282	257	230	200	168	134	98	61	20	20	20	20
Profile 2 (positive)	114	94	75	62	37	37	37	37	24			

**Table 3.1:** Zeeman slower winding configuration, expressed in terms of lengths of successive layers, measured from the start (bias and profile 1) or end (profile 2) of the slower. Higher numbered layers are radially farther away from the center of the ZS and are shorter, since the largest fields (and hence 'tallest' coils) required are at either end.

'valleys' from the previous layer). After each layer is wound, it is coated in a hightemperature encapsulating resin<sup>33</sup> and allowed to cure; this mechanically sets the layer and fills in the gaps between wires, allowing better thermal conductivity for cooling. The exact winding configuration was optimized with a MATLAB program and is given in table 3.1. The coils are run with a three-channel power supply<sup>34</sup> capable of supplying 10A and up to 32V to each coil.

### 3.4.3 Characterization of the Slower

The small diameter (8 mm) of the differential pumping tube at the center of the Zeeman slower makes measuring its magnetic field impossible with readily available equipment. Because of this (with further motivation from delays in the arrival of the vacuum components), we decided to wind a prototype ZS coil in order to iron out technical issues with the winding process and characterize the field produced by our coil geometry. The prototype was formed on a 3/4-inch (19 mm) outer diameter aluminum tube<sup>35</sup> but otherwise possessed identical geometry to the slower implemented on the experiment.<sup>36</sup>

Figure 3.7 shows the measured field from this prototype slower, for a bias current of 3.26 A and profile currents of 3.20 A – these values are predicted from our coil simulations to yield slowing parameters  $\eta = 0.5$  and  $\Delta = 2\pi \times -540$  MHz. As shown in 3.7a, we find that the measured profile is a good match (up to discrepancies in

 $<sup>^{33}\</sup>rm{Electrolube}$  ER2223, rated for temperatures of 180°C long-term (and 210°C for 30 min); we choose this epoxy to withstand baking.

<sup>&</sup>lt;sup>34</sup>Hameg HMP4030, with 384W maximum total power output, from Rohde & Schwarz.

<sup>&</sup>lt;sup>35</sup>The final ZS was wound on a 20 mm OD tube; this small discrepancy was due to the discrete sizes of stock tubing available. We find that this 1mm difference in diameter makes little difference in the field profile; nonetheless we alter our field simulations for purposes of comparison to the prototype.

<sup>&</sup>lt;sup>36</sup>Up to some messiness in the wires at either end, which was remedied in the final version with a more elegant means of fastening wires in place at the end caps.



Figure 3.7: Measured (a) magnetic field and (b) field residuals for the prototype Zeeman slower at currents of  $I_{\text{bias, meas}} = 3.26$  A and  $I_{\text{prof1, meas}} = I_{\text{prof2, meas}} \equiv I_{\text{prof, meas}} = 3.20$  A, corresponding to the currents expected to give  $\eta = 0.5$  and  $\Delta = 2\pi \times -540$  MHz. The curve is the simulated field plotted for  $I_{\text{bias}} = 0.88I_{\text{bias, meas}}$  and  $I_{\text{prof}} = 0.96I_{\text{prof, meas}}$ , suggesting that the density of windings for the bias (profile) coil is roughly 88% (96%) of the ideal value.

the tails which we attribute to messy winding at either end) to the simulated field for a bias (profile) current 88% (96%) of the experimental value.<sup>37</sup> This accounts for lower winding densities in the physical slower – apparently we were able to wind roughly 9 turns per 10 mm, rather than the ideal 1 turn/mm. As shown in the inset to figure 3.7b, for the middle 30 cm of the slower the field deviates by typically under 5 G, and no more than 10 G, from the simulated value. The behavior at the tails diverges more strongly, but here the effect is only to slightly diminish the capture velocity (start of the slower) and increase the final velocity (end of slower – this behavior is tolerable since we actually *want* finite final velocities so the atoms can traverse the ~ 15 cm between the ZS and MOT).

We next fit the central portion<sup>38</sup> of the measured profile to the analytical curve of equation 2.10 in order to extract slowing parameters (figure 3.8a). This yields

<sup>&</sup>lt;sup>37</sup>i.e. the plotted curve is for  $I_{\text{bias}} = 0.88 \cdot 3.26$  A, and  $I_{\text{prof}} = 0.96 \cdot 3.20$  A.

<sup>&</sup>lt;sup>38</sup>The region of positive slope.



Figure 3.8: Extraction of predicted Zeeman slowing parameters from measured *B* field data. (a) Fit to the square-root field profile of equation 2.10 revealing slowing parameters for this configuration of  $\eta = 0.47(1)$  and a -534(4) MHz laser detuning. We also plot numerically-calculated (b) atom trajectories and (c) time-varying  $\eta$  values for initial velocities ranging from 280 to 460 m/s, and for laser parameters of  $\Delta/2\pi = -533$  MHz, 100 mW of power, and a beam diameter converging from 8mm to 4mm at the start of the ZS. All captured atom trajectories converge onto the same curve, explicitly demonstrating a narrowing of the velocity distribution that leads to cooling as well as slowing, and here we achieve final velocities of 34 m/s at the end of the slower (region II) and 6 m/s at the MOT (end of region III).

values of  $\eta = 0.47(1)$  and  $\Delta = 2\pi \times -534(4)$  MHz, both of which are reasonably close to the nominal values of 0.5 and -540 MHz, respectively. Once again, the discrepancies here are largely attributable to a winding density below unity.

While this fit tells us something of the steady-state behavior of the slower (in particular,  $a_{\rm ZS} = \eta a_{\rm max}$  over the middle 30 cm), it does not capture the effect of the tails and therefore tells us little of the capture velocity or final velocity.<sup>39</sup> In order to learn about these quantities we turn to numerical methods and solve the one-dimensional equation of motion for an atom through the slower,<sup>40</sup>

$$\ddot{x} = -\eta a_{\max} = -\frac{\hbar k\Gamma}{2m} \frac{s(x)}{1 + s(x) + 4/\Gamma^2 (\Delta + k\dot{x} - \mu' B(x)/\hbar)^2}$$
(3.1)

 $<sup>^{39}\</sup>eta = 0.47$  does suggest a capture velocity around 455 m/s (equation 2.9), but because its extraction ignores the field at the start of the slower, it cannot imply anything precise.

<sup>&</sup>lt;sup>40</sup>In order to do this we assume the field smoothly varies between measured data points and craft a simple interpolation function. Though non-negligible ripples at length scales below 1 cm are physically unlikely, they could lead to dramatically different slowing results.

where for generality we take the saturation parameter s(x) to vary with position. Because of a divergence of the atom beam over the length of the slower (appendix B), we choose to focus the laser from the full ZS tube diameter of 8mm at the end of the slower to approximately 4mm at the start – this somewhat improves the capture velocity without contributing to atom losses.<sup>41</sup> Accordingly, for 100 mW of laser power, this leads to saturation parameters of  $s \approx 13$  at the start of the ZS and  $s \approx 3.5$  by the end.<sup>42</sup>

Figures 3.8b-c show atom trajectories and instantaneous  $\eta$  values calculated for various initial velocities, the above laser parameters, and a detuning of -533 MHz. Here we label three regions, reflecting (I) atom travel before the ZS, (II) the Zeeman slowing proper, and (III) post-slowing by off-resonant light between the ZS and the MOT. We find a capture velocity of approximately  $v_c = 465$  m/s for this field/laser configuration, meaning that atoms moving at below this speed in region I will be captured (note that 540 MHz is the Doppler shift for 215 m/s atoms, which is why the 280 m/s trajectory is so significantly slowed in this region). By the end of region II the atoms still have a velocity of 34 m/s, and the remaining slowing happens during their long interaction time (~ 10 ms) with off-resonant light between the ZS and MOT. For these simulation values we find a final velocity of 6 m/s at the MOT, which is within its capture range. The capture and final velocities can be tuned by adjusting the laser detuning and/or bias field, adjusting the laser power and beam divergence, and changing  $\eta$  by tuning the profile coils.

# 3.5 Magnetic Fields

The final, critical component of the experiment is the precise control of magnetic fields used to introduce gradients and homogeneous fields at the location of the atoms (for, e.g., trapping and the control of scattering properties) as well as to compensate external fields which could interfere with the physical processes we hope to study. Section 3.5.1 describes the coils used to apply field gradients and offsets at the center

 $<sup>^{41} \</sup>rm Since$  those atoms with radii greater than  $\sim 2 \rm mm$  at the start of the slower will typically hit the tube walls before reaching the MOT chamber.

<sup>&</sup>lt;sup>42</sup>We assume a uniform beam to calculate these saturation parameters, which effectively reflect a spatially averaged intensity value  $s_{\text{avg}}$ . In reality, for a Gaussian beam the maximum intensity gives  $s_{\text{max}} = 2s_{\text{avg}}$ , but it also drops to  $2/e^2 = 0.27$  of  $s_{\text{avg}}$  at the  $1/e^2$  waist. The Gaussian intensity is equal to the averaged value at  $r = 0.59\omega_0$  (where r is the radial position and  $\omega_0$  the  $1/e^2$  waist), suggesting that this average value actually is a decent reflection of the intensity a 'typical' atom would experience. Nonetheless, for a full treatment of the Gaussian profile we would need a two-dimensional simulation.

of the MOT chamber. Section 3.5.2 describes a large cage used for external field compensation, and bias field introduction, at the MOT chamber as well as over the optical transport axis between the MOT and (planned) science chamber.

#### 3.5.1 MOT Chamber

The most obvious requirement for the MOT chamber is the introduction of a quadrupole field, with field gradients along the vertical and radial directions, needed for the production of the MOT itself. This is achieved with two circular coils, radius R, in anti-Helmholtz configuration, separated by a vertical (z) distance 2D and situated symmetrically with respect to the origin, defined as the center of the MOT (figure 3.9a). We run currents through each coil of equal magnitude and opposite direction. For a suitable separation between the coils  $D = \sqrt{3/4R}$ ,<sup>43</sup> this produces a gradient along the z axis  $\frac{dB_z}{dz} \approx \text{const}$  for  $z \ll D$ . In particular, for this coil arrangement

$$B_z(z) = \frac{48}{49} \sqrt{\frac{3}{7}} \frac{\mu_0 NI}{R^2} z + \mathcal{O}(z^5) , \qquad (3.2)$$

where  $\mu_0$  is the vacuum permeability, I is the current, and N is the number of turns of wire. By  $\nabla \cdot \mathbf{B} = 0$ , the gradient along the radial directions is negative with half the magnitude,  $2\frac{dB_x}{dx} = 2\frac{dB_y}{dy} = -\frac{dB_z}{dz}$ . In practice, when designing our fields we sum the individual contributions from each turn of wire separately to account for the non-negligible size of the wires with respect to R and D; nonetheless, equation 3.2 captures the essential behavior of the coils.

It is clearly beneficial to minimize the radius of our coils in order to achieve large gradients with modest currents; in practice the limiting constraint is the 70 mm diameter of the vertical MOT windows. We ultimately converge on a solution involving 6 (radial) × 4 (vertical) turns of  $1.25 \times 2.5$  mm rectangular copper wire, with inner radius  $R_{\rm in} = 44$  mm and minimum half-separation  $D_{\rm in} = 36.6$  mm; this yields a calculated vertical (radial) gradient of  $830(-415) \frac{{\rm mG/cm}}{{\rm A}}$ . More complete properties of these and the other MOT coils are given in table 3.2. The rectangular wire is chosen to maximize physical contact, and therefore thermal conductivity, between adjacent windings, and the relatively large cross-sectional area is chosen to reduce the ohmic power dissipation for a given number of windings.<sup>44</sup> At the currents

<sup>&</sup>lt;sup>43</sup>One can show that when  $R = \sqrt{4/3D}$ , the third-order term in a Taylor expansion of the field vanishes, rendering it maximally linear (see, for instance, [12]). We operate in this configuration, though it is also possible to use R = 2D, in which case the gradient is maximized.

<sup>&</sup>lt;sup>44</sup>The dissipated power is  $I^2R$ , where R here is the wire resistance, and  $R \propto 1/A$ , where A is the



Figure 3.9: Coil setup at the MOT chamber. (a) Schematic representation of the three sets of coils mounted at the chamber, along with definitions of the geometrical parameters R and D used in the text. (b) Section view of the chamber with 3D printed coil mount(s) and coils attached.

run for MOT operation the coils dissipate under 5W of power and experience only negligible heating.

In addition to a quadrupole field for MOT operation, we also need the ability to implement homogeneous magnetic fields along the vertical axis. These are used (1) to establish a quantization axis and preserve spin polarization and (2) to control the scattering properties of the atoms. In particular, as discussed in section 2.2.1, the competition between gravity and the scattering force in the 583 nm MOT gives us spin polarization into the  $|F, m_F\rangle = |6, -6\rangle$  ground state 'for free.' By turning on a small homogeneous bias field when we extinguish the quadrupole field we will be able to preserve this polarization. The other important use for homogeneous fields is to tune the magnitude of two-body contact interactions, characterized by the *s*-wave scattering length  $a_s$ , through the use of Feshbach resonances [58].<sup>45</sup> Er has a particularly dense Feshbach spectrum [41, 72] – including resonances as low as 900 mG for <sup>168</sup>Er and 50 mG for <sup>166</sup>Er [71] – so precise control of magnetic fields is especially

cross-sectional area. Clearly the more important concern is to minimize I when designing coils for low power dissipation, which can be achieved both by reducing the coil radius and by increasing the number of windings N. Nonetheless, practical concerns (e.g. the difficulty of actually winding the wire) and the necessity of decent tuning resolution (so that small changes in I don't lead to large changes in B) dictate that N not be made too large. For fixed N we therefore opt to increase A.

 $<sup>^{45}</sup>$ We need control over  $a_s$  for scientific purposes – namely to tune the relationship between the contact and dipole-dipole interactions – but this will mostly happen in the science cell rather than at the MOT. At the same time, however, for efficient evaporative cooling – which we might initiate in the MOT chamber before transport – it is important to have a reasonably large scattering length.

	Quadrupole	Feshbach	Low Field
Windings	$6 \times 4$	$6 \times 2$	$2 \times 2$
$R_{ m in}-R_{ m out}$	44mm – $52.1$ mm	49.3mm – $57.4$ mm	65.8 mm - 68.5 mm
$D_{ m in} - D_{ m out}$	$36.6 \mathrm{mm} - 47 \mathrm{mm}$	$24 \mathrm{mm} - 29.2 \mathrm{mm}$	$31 \mathrm{mm} - 36.2 \mathrm{mm}$
Magnetic Field	$833 (-416) \frac{mG/cm}{A}$	$2.03 \mathrm{~G/A}$	535  mG/A
Resistance	$78 \text{ m}\Omega$	$46.3 \text{ m}\Omega$	$18.4 \text{ m}\Omega$
Inductance	$150 \ \mu \mathrm{H}$	$50 \ \mu H$	$10 \ \mu H$

**Table 3.2:** Nominal winding dimensions and calculated properties for the various coils mounted on the MOT chamber. The wire is  $1.5 \times 2.5$ mm rectangular-cross-section, enameled copper wire, and we take it to have a 50  $\mu$ m enamel coating on each surface for purposes of dimensional calculations. In determining the optimal coil geometries we calculated the fields from each current loop separately, without assuming a specific (anti-)Helmholtz condition on R and D. These values were then varied in order to minimize the (third) second derivative of the field at the MOT center; the results actually differ somewhat from the single-winding (anti-)Helmholtz configurations presented in the text because of the finite size (and asymmetry) of the wires.

important in this regard. While this suggests a need for high-resolution tuning, we also might require higher fields – e.g. to access broader Feshbach resonances in Er, or Feshbach resonances in an alkali species to be added to the experiment at a later date<sup>46</sup> – which runs counter to the need of high tuning resolution.<sup>47</sup> Accordingly, we choose to implement two sets of homogeneous coils at the MOT chamber – one that accesses low fields under ~ 5 G with high resolution, and another that will achieve higher fields of at least 30-40 G (depending on the power supply) with less resolution.

For both coils we operate in Helmholtz configuration, consisting of two circular coils radius R running equal currents of the same sign (in contrast to anti-Helmholtz configuration) and placed symmetrically about the origin with half-separation D = R/2 (figure 3.9a). For this separation the second-order term vanishes and we are left with a maximally homogeneous field

$$B_z(z) = \frac{8\mu_0 NI}{5\sqrt{5}R} + \mathcal{O}(z^4)$$
(3.3)

As with the quadrupole coils, we wind both homogeneous coils using  $1.25 \times 2.5$ mm rectangular copper wire. The MOT chamber is recessed about the vertical windows,

 $<sup>^{46}{\</sup>rm The}$  lowest resonances in  $^{39}{\rm K},$  for instance, are at 26 and 33 G, depending on the atomic state and desired resonance width.

<sup>&</sup>lt;sup>47</sup>By tuning resolution, we mean the precision with which we can tune the field as limited by current stability and tunability. For a 10G/A field it is difficult to control the field below 10 mG as most power supplies will have current noise and tuning precision around the 1 mA level.

allowing us to achieve a small vertical separation (and therefore larger field) for the higher-field Feshbach coils (consisting of  $6 \times 2$  windings) of  $D_{\rm in} = 24$  mm. The inner radius of this set of coils is  $R_{\rm in} = 49.3$  mm, and the calculated field produced is slightly above 2 G/A. The low-field coils, meanwhile, consist of 4 windings of rectangular copper wire and produce approximately 530 mG/A. See table 3.2 for details on the coil properties.

We operate all three coils at present with 10A power supplies,<sup>48</sup> which caps their fields at 8.3 G/cm, 20.2 G, and 5.3 G, respectively. These values should be sufficient for the quadrupole and low-field coils, though we will most likely in the future require more current through the higher-field Feshbach coils.

All three are mounted on the MOT chamber using a custom-designed, 3D printed support that enforces the inner coil dimensions given above (figure 3.9b). The material is an ABS-like plastic with low thermal conductivity, presenting problems for coil cooling in principle, though we introduce holes in the mount – ostensibly for the purpose of diminishing warping of the 3D printed part – that also seem to improve heat dissipation. While switching to a metal mount would improve thermal dissipation further still, the major advantage of plastic is the elimination of unwanted eddy currents. Pockets for the quadrupole and Feshbach coils in the mount allow room for water cooling tubes should that later prove necessary.

#### 3.5.2 Compensation Cage

Due to its highly magnetic nature<sup>49</sup> as well as its unusually dense Feshbach spectrum, compensation of stray magnetic fields is particularly important for erbium. While we leave magnetic compensation at the science chamber to the future,<sup>50</sup> during this work we have designed and built<sup>51</sup> a large cage for three-axis compensation both at the MOT chamber and over the optical transport axis between the MOT and science chambers. It consists of three pairs of rectangular coils in maximally-uniform Helmholtz configuration<sup>52</sup> about the MOT center, with half-separations  $D_{x,y,z}$  and side lengths  $a_{x,y,z}$  and  $b_{x,y,z}$ . A third coil exists a further distance  $2D_y$  along the

 $<sup>^{48}</sup>$ Delta Elektronika, ES 015-10.

<sup>&</sup>lt;sup>49</sup>Erbium's  $m_J = -6$  ground state means a magnetic force roughly 6 times stronger than for many alkalis with  $m_F = \pm 1$ .

<sup>&</sup>lt;sup>50</sup>Because compensation will be most important in the science cell, we plan to use smaller coils with a better response frequency than the cage, with its high inductances.

<sup>&</sup>lt;sup>51</sup>While I designed the cage, it was built primarily by Milan Krstajic. I contributed some to the wire winding (which, it turns out, was a *much* more tedious process than expected).

<sup>&</sup>lt;sup>52</sup>Which for square coils exists at a half-separation  $D \approx 0.54a$ , where a is the side length. Note that our coils are rectangular and so this relationship differs somewhat.



**Figure 3.10:** Cage used for compensation of external fields at the MOT and over the optical transport axis, as well as for the application of bias fields to preserve spin polarization during transport.

transport (y) axis towards the science cell for use during optical transport. Figure 3.10 shows the cage geometry. We use equations for the field from rectangular coils given in [108] to numerically optimize coil positions for maximal uniformity at the MOT center.

The cage is approximately 1000mm long along the transport axis and 550mm in width and height, centered at the MOT chamber. This means that it comfortably covers the entire transport axis between the MOT chamber and science cell  $\sim 35$  cm away. Each coil consists of 24 windings of 1mm-diameter enameled copper wire, mounted in a structure made of aluminum U-shaped channeling 10mm wide  $\times 15$ mm high. All coils are wound in two sets of 12 windings, each of which generates half the full field. This allows us, if necessary, to wire them separately and produce both positive and negative fields without need for an expensive bipolar power supply. Specific properties of the coils are given in table 3.3.

During MOT operation we run equal currents through each symmetric pair of coils to achieve homogeneous fields with magnitudes given in table 3.3. During transport, which breaks the symmetry of the cage, such uniform fields are impossible to maintain and instead we aim simply to minimize spatial variations and higher-order terms (gradients and harmonic potentials).

Along the y axis, by adding a third coil  $(y_3)$  we are able to produce a maximally

	x	y	z
Windings	$4 \times 6$	$6 \times 4$	$6 \times 4$
Center Dim. $a \times b$	$977 \text{mm} \times 482 \text{mm}$	$540 \operatorname{mm}(x) \times 515 \operatorname{mm}(z)$	$560\mathrm{mm} \times 1010\mathrm{mm}$
Center Sep. $2D$	288mm	286mm	$336 \mathrm{mm}$
Field at MOT	664  mG/A	744  mG/A	583  mG/A
Resistance	$3.3 \ \Omega$	$2.4 \ \Omega$	$3.6 \ \Omega$
Inductance	$3.5 \mathrm{~mH}$	$2.3 \mathrm{mH}$	$3.8 \mathrm{mH}$

**Table 3.3:** Dimensional, magnetic, and electrical properties of the compensation cage. Winding numbers are given as (radial)  $\times$  (axial), and dimensions reflect the center of the wire bundle (i.e. winding "3.5  $\times$  2.5"). The messy-looking dimensions are chosen because they correspond to clean dimensions for the support structure (x-coil channels 970mm and 475mm long, for instance).

homogeneous field during transport by running equal currents  $I_{y_1} = I_{y_3}$  through the extremal coils and a smaller current  $I_{y_2} = 0.53I_{y_1}$  through the middle coil (figure 3.11b). This leads to a field at the midpoint of the transport axis of 648 mG/A, and the field varies by a tolerable 3.7% between the midpoint and the science cell.

Figure 3.11a, meanwhile, shows that the bias field  $B_z$  increases on-axis towards the science cell and varies by at most 4.3% during transport. Again, this variation is reasonable, particularly if we choose to sit at a field far from any Feshbach resonances, and should allow for the preservation of spin polarization by virtue of always being larger than residual x and y fields (which we can use the x and y coils to compensate).

We further consider field inhomogeneities that could lead to trap instability during transport. The most notable of these is a quadratic term in  $B_z$  along z(i.e.  $d^2B_z/dz^2 = B_z''$ ) which continually increases along the transport axis (figure 3.11c). Because the field increases off axis this leads to an anti-confining potential for our high-field seeking atoms ( $m_J = -6$ ). The anti-trapping frequency can be calculated from

$$U_{\rm harm}(z) = \mu_B g_J m_J B_z'' z^2 = \frac{1}{2} m \omega^2 z^2 \implies \omega = \sqrt{\frac{2\mu_B g_J m_J B_z''}{m}} \tag{3.4}$$

and, as seen in figure 3.11c for a nominal bias field of 1.5 G, rises to a maximum value of  $i \times 0.15$  Hz at the science cell. This is negligible compared to the expected radial ODT trapping frequencies of order 1 kHz [91].<sup>53</sup>

<sup>&</sup>lt;sup>53</sup>Of course, the axial trapping frequency is closer to 20 Hz, so inhomogeneities along the y axis are more significant. Nonetheless, we find that the linear and quadratic terms here (i.e.  $dB_z/dy$ and  $d^2B_z/dy^2$ ) lead, respectively, to displacements of the trap center of under 1  $\mu$ m (c.f. the Rayleigh range of order 1 mm) and to trapping frequencies of order 0.1 Hz (as above, though



Figure 3.11: Magnetic field behavior of the compensation cage over the MOT-science cell transport axis. (a) The bias field  $B_z$ , here plotted for  $I_{z_1} = I_{z_2} = 2.57$ A corresponding to 1.5 G, varies by 4.3% over the full transport axis, which should allow for reasonable adiabatic transport to maintain spin polarization. We choose 1.5 G because it sits far away from Feshbach resonances for both <sup>166</sup>Er and <sup>168</sup>Er (see [41, 72]). (b) By adding a third y coil we produce a reasonably homogeneous field  $B_y$  along the transport axis, here shown for currents  $I_{y_1} = I_{y_3} = 1$ A and  $I_{y_2} = 0.53I_{y_1}$ , that varies by at most 3.7%. Without the extra coil the field would drop towards zero away from the MOT. (c) Vertical anti-trapping frequencies, caused by nonzero quadratic terms  $d^2B_z/dz^2$  in the bias field, as a function of distance along the transport axis. Insets are cuts of the field along the z axis at y = 0 and y = 20 cm, demonstrating the appearance of inhomogeneities.

We finally note that we have purchased a sensitive magnetometer<sup>54</sup> which will eventually be used for active compensation of external fields in the experiment. The large inductances of the compensation cage mean that active compensation can only be achieved for DC and time-varying fields slower than a few kHz, though with smaller coils at the science cell we will not have this limitation.

here at times the potential is confining rather than anti-confining). We have similarly confirmed that other diagonal and off-diagonal, 1st- and 2nd-order terms are negligible compared with the expected ODT potential.

<sup>&</sup>lt;sup>54</sup>Honeywell HMC2003 three-axis magnetometer,  $\pm 2$  G range and 40  $\mu$ G resolution.

# Chapter 4

# **Ongoing and Future Work**

Though an overwhelming majority of the work encompassing this thesis involved the experimental design and construction detailed in the previous chapter, we have very recently begun the transition towards system switch-on and characterization. The first significant milestone in this process was the detection of an erbium beam in our vacuum chamber, discussed in section 4.1. This represents a necessary and important starting point from which to begin optimizing the laser cooling and searching for a MOT. Section 4.2 looks ahead and discusses two scientific projects upon which we hope to embark after achieving a fully functional apparatus. In section 4.3 we offer final remarks and conclude the thesis.

# 4.1 Initial Observation of the Erbium Beam

For initial work with the erbium atomic beam,<sup>1</sup> we decided to postpone vacuum chamber bakeout and use only a turbo pump<sup>2</sup> backed by a scroll pump,<sup>3</sup> which together allow us to maintain pressures between  $10^{-6}$  and  $10^{-7}$  mbar in the HV section without turning on the ion pumps.<sup>4</sup> After pumping down, we began by heating up the erbium oven to a temperature of 1100°C in the main effusion cell and 1200°C in the separately controllable 'hot lip' section. The effusion cell holds

 $<sup>^1{\</sup>rm The}$  experimental work in this section was conducted jointly by myself, Milan Krstajic, and Rob Smith.

<sup>&</sup>lt;sup>2</sup>Leybold Oerlikon Turbovac TW70H.

<sup>&</sup>lt;sup>3</sup>Leybold Oerlikon Scrollvac SC 5D.

<sup>&</sup>lt;sup>4</sup>Turning on the ion pumps could reduce their lifetime, and atomic beam work does not require pressures lower than those achievable without them. The system will be baked, the ion pumps turned on, and the NEG elements activated, after the science cell is added to the system in fall 2018, at which point the pressures should be reduced to  $\sim 10^{-9}$  mbar in the HV section and  $\sim 10^{-11}$  mbar in the UHV section.

the Er source metal and is responsible for producing a significant vapor pressure, while the hot lip contains apertures which help to collimate the beam.<sup>5</sup> A slightly larger aperture, made from a custom copper gasket, is placed downstream from the oven and prevents those atoms which would not reach the MOT anyway from passing through TC and eating up laser power [92].

For initial study of the atomic beam, we passed a collimated, 5mm-diameter beam of 401 nm light vertically downwards through the characterization windows just after the TC cube. By collecting the atomic beam fluorescence (which is faintly visible by eye with ~ 10 mW of laser power) with a 25mm lens and a photodiode through another viewport at 90° to the pumping beam, we observed well-resolved peaks corresponding to the relatively abundant <sup>166</sup>Er, <sup>168</sup>Er, and <sup>170</sup>Er isotopes, as well as a noticeable <sup>164</sup>Er peak and one or two <sup>167</sup>Er hyperfine transitions (figure 4.1).<sup>6</sup> Fitting Lorentzian lineshapes to these peaks yields widths of roughly 60 MHz,<sup>7</sup> which is significantly larger than the expected power-broadened linewidth of ~ 40 MHz.<sup>8</sup> Unsurprisingly, this suggests that there is a non-negligible transverse velocity spread associated with our mechanically collimated – but not yet laser cooled – erbium beam, leading to a finite Doppler width.

To characterize this effect, we fit the observed spectrum<sup>9</sup> to a sum of Voigt profiles (figure 4.1a) in an effort to extract the Doppler width, which in this case is on the order of the homogeneous broadening. Because these widths are of the same order, even small imperfections in the signal – including unresolved hyperfine peaks or slight asymmetries due to misalignment through the chamber – can dramatically affect the extracted parameters. Indeed, while the individual width parameter errors are on the order of 1 MHz for each of the three most prominent peaks,<sup>10</sup> the scatter between the three is much larger. Taking a conservative approach to error estimation, by accounting for this scatter we arrive at values of  $\Gamma = 36(9)$  MHz for the homogeneous linewidth and  $\sigma = 23(3)$  MHz for the Gaussian width. The former is consistent with an expected power broadening of ~ 40 MHz.

We can use the fitted Doppler width to estimate the transverse temperature in our beam. In particular, comparing the standard Gaussian function  $f \propto \exp[-(\omega - \omega)]$ 

<sup>&</sup>lt;sup>5</sup>The higher temperature helps, in part, to prevent these apertures from clogging.

<sup>&</sup>lt;sup>6</sup>The most prominent of these hyperfine peaks (we include it in the fit in figure 4.1a) corresponds to the  $F = 19/2 \rightarrow F' = 21/2$  transition [102].

<sup>&</sup>lt;sup>7</sup>We calibrate the frequency axis using published isotope shift values from [102].

 $<sup>^8\</sup>mathrm{This}$  corresponds to a 5mm-diameter beam with about 10 mW of power.

 $<sup>^9\</sup>mathrm{Or}$  really the average of 16 such spectra.

 $<sup>^{10}\</sup>mathrm{It}$  is harder to trust the values for the  $^{164}\mathrm{Er}$  and hyperfine peaks, so we exclude them in this analysis.



**Figure 4.1:** Erbium atomic beam fluorescence spectrum at 401 nm. (a) Fit of the normalized signal to a sum of five Voigt profiles, corresponding to the four most abundant bosonic isotopes as well as the most prominent <sup>167</sup>Er hyperfine transition. Another feature due to the fermionic isotope is visible between the <sup>168</sup>Er and <sup>166</sup>Er peaks, and an unresolved shoulder on the right side of the <sup>168</sup>Er line, corresponding to yet another hyperfine transition, can be seen with the right amount of squinting. The relative peak heights observed here agree to within 1% with the expected isotopic abundances in table 1.2. (b) Fits to the <sup>170</sup>Er peak with both a Voigt profile and a Lorentzian. The Voigt profile captures the tails of the fluorescence feature much better than its purely Lorentzian counterpart.

 $\omega_0)^2/(2\sigma^2)$  with the Doppler broadening lineshape [42],

$$f_{\rm Dopp}(\omega) = \frac{c}{u\omega_0\sqrt{\pi}} \exp\left[-\frac{c^2}{u^2} \left(\frac{\omega-\omega_0}{\omega_0}\right)^2\right] , \qquad (4.1)$$

where  $u = \sqrt{2k_BT/m}$  is the most probable velocity, we find that the temperature is given by

$$T = \left(\frac{\sigma c}{\omega_0}\right)^2 \frac{m}{k_B} = \left(\frac{\sigma}{k}\right)^2 \frac{m}{k_B} \tag{4.2}$$

This gives a measured transverse atomic beam temperature of T = 1.7(5) K, which is 3 orders of magnitude higher than the Doppler limit. This corresponds to a most probable velocity u = 13 m/s and implies, for 500 m/s longitudinal atoms, a reasonable-sounding  $1.5^{\circ}$  beam divergence angle. With a decently well optimized transverse cooling stage we should be able to cut the divergence by roughly a factor of two and bring the transverse temperature down to around 400 mK prior to Zeeman slowing (see appendix B).<sup>11</sup>

<sup>&</sup>lt;sup>11</sup>The expected 5.5 m/s most-probable cooled velocity given here is in rough, but not precise, agreement with the number expected from a detailed numerical simulation of the atomic beam cooling process implemented by another student, Milan Krstajic [92]. The cooled transverse ve-

# 4.2 Future Scientific Projects

Even in the early stages of an experiment when day-to-day work requires immersion in technical minutiae and engineering details, it is exciting and worthwhile to consider proposed avenues of research suitable to the apparatus. Below we consider two such possibilities.

#### **Roton Physics in Dipolar Condensates**

First proposed by Landau in the 1940s [109], the concept of the roton relates to a peculiar feature predicted, and later observed [110], of the energy-vs.-momentum dispersion relation  $\omega(k)$  for superfluid <sup>4</sup>He. The characteristic excitation spectrum begins with a linear dispersion relation at low k, characteristic of collective phonon modes, or sound waves. The spectrum then reaches a local maximum (these excitations are sometimes referred to as maxons) and decreases into a quadratic energy well, characterized by a minimum wavenumber  $k_{\rm rot}$  and energy gap  $\Delta$ . These excitations were dubbed rotons by Landau.<sup>12</sup> For higher momenta still the spectrum becomes quadratic in k, characteristic of free particle excitations (figure 4.2a).

Though the microscopic interpretation of these rotons is still not agreed upon in the context of superfluid helium, one possible picture arises in light of Feynman's relation [111],

$$\omega(k) = \frac{\hbar k^2}{2m} \frac{1}{S(k)} , \qquad (4.3)$$

where m is the atom mass and S(k) is the static structure factor, which can be thought of as quantifying the structured-ness or crystallinity of the system for excitations of momentum  $\hbar k$ . In this picture, the roton minimum arises due to a maximum in the structure factor, suggesting that roton excitations have something to do with the particles' tendency to locally organize or "stay apart" (Feynman's words) [111] – this could also suggest that rotonization hints towards solidification or even supersolidity.<sup>13</sup>

More recently, it was shown that dipolar Bose-Einstein condensates (dBEC) are

locity distribution in that case is peaked with a width of about 1 m/s but has long tails giving RMS velocities between 3 and 5 m/s. This is due to a coupling between axial velocity and radial cooling efficiency, whereby faster atoms have less interaction time with the TC beams and remain hot, hence the wide tails. 500 m/s atoms presumably already fall within the widened region of this distribution, which could be one reason our simplistic model slightly overestimates the result from the detailed numerical one.

<sup>&</sup>lt;sup>12</sup>This may have been a misnomer, as it is not clear that these excitations are in any way rotational in nature.

<sup>&</sup>lt;sup>13</sup>Though the latter connection is now generally accepted to be incorrect [59].



Figure 4.2: Roton mode in dipolar condensates (dBEC). (a) Dispersion relation, normalized to the vertical trapping frequency  $\omega_z$  and harmonic oscillator length  $l_z$ , of a 2D dBEC for fixed dipolar length and variable *s*-wave scattering length, calculated from equations in [59, 112]. For  $a_{dd} \leq a_s$  (top curve), the dispersion relation exhibits only phonon-like (low *k*) and particle-like (high *k*) excitations. For lower  $a_s$ , the roton minimum appears at momentum  $k_{rot}$  and energy  $\Delta$ , and for large-enough  $\varepsilon_{dd}$  an instability occurs (bottom curve). (b) 'Pancake' trapping geometry for observation of roton physics in dBEC (figure from [113]).

expected to exhibit a roton-maxon spectrum similar to that found in superfluid helium [114]. This occurs in 2D polarized gases of erbium atoms, confined in a uniform in-plane potential<sup>14</sup> and a tightly-confining vertical (z) harmonic potential<sup>15</sup> and polarized out of the plane via a magnetic bias field (figure 4.2b). For low-k excitations the side-by-side orientation of the atoms leads to a repulsive potential, and these excitations exhibit a linear, phonon-like dispersion relation. For higher k and sufficiently large dipolar interactions ( $\varepsilon_{dd} > 1$ ), however, the excitations begin to become 3D in nature, insofar as the dipoles are pushed out of the plane, and the attractive character of the 'head-to-tail' portion of the dipole-dipole potential becomes apparent. This leads to a roton minimum, the strength of which depends on the s-wave scattering length (figure 4.2a). High-k excitations are once again particle-like.

For small enough contact interactions and/or large enough particle densities, the roton minimum turns into a roton instability, wherein the excitation frequencies become imaginary for certain values of k and are no longer stable (lowest curve in figure 4.2a). This phenomenon was recently observed in a 1D gas of Er [79]. As one of the first projects on the new experiment, we hope to go one step further and map out the roton-maxon dispersion relation using Bragg spectroscopy, a coherent

 $<sup>^{14}\</sup>mathrm{Perhaps}$  enforced by the projection of a cylindrical, blue-detuned beam for radial confinement in the dark central region.

<sup>&</sup>lt;sup>15</sup>Imposed with a red-detuned sheet beam.

process through which one can measure the excitation energy for a given wavenumber k [115–117]. Doing so could further lead to studies of the superfluid critical velocity and the Landau criterion [114].

#### Supersolidity in Lattices with Long-Range Interaction

A supersolid is a counterintuitive quantum phase which simultaneously exhibits characteristics of a superfluid (evidenced by frictionless flow or particle delocalization<sup>16</sup>) and a solid (evidenced by spatial order or a broken translational symmetry<sup>17</sup>) [118]. To date it has not been unequivocally demonstrated, though recent results in cavity-based and spin-orbit coupled cold-atom systems claim to have done so [119, 120]. In dipolar quantum gases, the phenomenon of quantum droplet formation has been linked to supersolidity, and active research is being pursued down that avenue [121].<sup>18</sup>

Another approach is to use optical lattices of dipolar atoms as a starting point. Perhaps the most promising strategy, considering experimental temperature limitations, is to work with large numbers of atoms per site [122]. As a potential project for the erbium experiment, we propose to project a 1D lattice onto the 2D pancake geometry described above in search of such a supersolid phase, which should appear at temperatures around 100 nK with roughly 200 atoms per lattice site.

# 4.3 Conclusion

The work undertaken over the course of this thesis has laid the foundation<sup>19</sup> for what looks to be a productive dipolar BEC experiment. The starting point was a nearly empty lab with one operational cooling laser, and at the time of writing the apparatus is near-fully assembled with all physical components necessary for a narrow-line MOT in place (figures 4.3 - 4.4). In particular, the present author has designed and set up a majority of the cooling optics; designed, wound, and tested the Zeeman slower<sup>20</sup>; designed and built the magnetic field coils for the MOT chamber (including a compensation cage and various gradient and bias coils); and more

<sup>&</sup>lt;sup>16</sup>Formally off-diagonal long-range order (ODLRO).

<sup>&</sup>lt;sup>17</sup>Formally long-range order (LRO).

<sup>&</sup>lt;sup>18</sup>The roton minimum also intuitively seems to lead to a breaking of translational symmetry through the introduction of a length scale associated with  $k_{\rm rot}$ , though resulting phases are likely unstable, meaning that quantum droplets represent a more plausible approach towards continuous-symmetry-breaking supersolidity in dBECs [121].

<sup>&</sup>lt;sup>19</sup>In collaboration with a PhD student [92] and (for a few months) two Part III students.

 $<sup>^{20}\</sup>mathrm{Or}$  really two such slowers.



Figure 4.3: Laser table on October 10, 2017 (left) and July 13, 2018 (right).

generally contributed to many aspects of the project at a broad level, including, notably, the specification and purchasing of the experiment's computer control system (not discussed in this thesis but briefly described in [92]).

The experiment has been set up on breadboards to facilitate a move, in early August 2018, to its new home at the University of Oxford. The science cell and ODT laser will arrive there early in the fall, and after installation of the cell, followed by chamber bakeout, it should be possible to push for BEC before the end of 2018. We already own the  $\sim 372$  nm laser that will be used to implement a repulsive box potential following the attainment of BEC.

First projects with the newly completed experiment will involve studies of roton physics, and then supersolidity, through the implementation of a uniform 2D 'pancake' trap followed by the introduction of a 1D lattice. More generally, the experiment will have the versatility to explore a variety of phenomena, including the effect of dipolar interactions on various quenched and driven behavior already observed in uniform alkali systems [51, 123]. If two of the more significant recent trends in ultracold-atom-based many-body physics have been the switch from harmonic to uniform potentials and the introduction of dipolar interactions, then our combination of both should represent an ideal platform from which to launch a novel series of quantum gas experiments.



Figure 4.4: Experiment table as of July 13, 2018.

# Appendix A Extended Laser Cooling Theory

In this appendix we discuss in more detail the theories of Zeeman slowing and narrow-line trapping presented in sections 2.1.1 and 2.2.1, respectively, of this thesis. Beginning with ZS, though the equations presented in 2.1.1 are correct for a *perfect* slower, they fail to accurately capture the variety of behavior that will happen if, among other things, the magnetic bias field or the laser power differs even slightly from its design value. In such circumstances, an atom's motion through the slower is no longer accurately described by constant-acceleration kinematics, and we must consider the differential nature of the slowing equations. We turn our attention to this problem in section A.1, while in section A.2 we consider the means through which one can achieve a nonzero final velocity (desirable because the atoms have to travel a finite distance before reaching the MOT) by tuning various slower parameters. In section A.3, meanwhile, we turn our attention to a simple discussion of the mechanisms underlying the capture behavior of the narrow-line MOT.

# A.1 Zeeman Slowing Out of Equilibrium

In order to develop a more intuitive picture of the Zeeman slower, it is helpful to visualize the relationship between the location of the atom and of the peak of the scattering curve,  $F_{\text{scatt}}(\delta)$  (equation 2.1), in position and velocity space. In particular, the maximum value of the scattering force is always found when the total detuning vanishes,

$$\delta(x,v) = \Delta + kv - \frac{\mu'}{\hbar}B(x) = 0.$$
 (A.1)

For a maximally efficient slower, in which the deceleration is as large as possible (i.e.  $a = a_{\max} = \hbar k \Gamma/2m$ ),  $\delta(x, v) = 0$  always and the deceleration parameter  $\eta$  is given by  $\eta_{\max} = \frac{s}{1+s}$ , where s is the laser saturation parameter. In the intuitive picture, the atom sits at the very top of the scattering curve as the atom and peak move simultaneously from 0 to  $x_0$  in position space, and from  $v_c$  down to 0 in velocity space (figure A.1(a)). While this allows us always to take advantage of maximal scattering, the obvious difficulty is in its stability: for any small fluctuation in the magnetic field or laser intensity (or even atom velocity due to quantum fluctuations in the scattering of photons), the atom necessarily drops off the peak of the scattering curve. While the curve continues to move at the same pace, the atom now experiences a decreased scattering force and is unable to keep up with the ideal trajectory – this problem compounds upon itself as the atom drifts farther away from the scattering resonance, and the result is a critically unstable trajectory.

Now, suppose we decrease the security parameter<sup>1</sup> below its optimal value, so that  $\eta = \frac{s}{1+s} < \eta_{\max} = \frac{s'}{1+s'}$ , where s' is the real laser saturation parameter and s < s'is the 'effective' saturation parameter corresponding to our choice of  $\eta$ .<sup>2</sup> Because the acceleration is now lower than that allowed by the laser radiation, there must be some finite total detuning that decreases  $\eta$  below  $\eta_{\max}$ ; in the intuitive picture, the atom drops off the peak of the scattering curve and 'surfs' along its side, following it at (to a decent approximation for the majority of the slowing) a near-constant offset related to<sup>3</sup>  $\delta$  (figure A.1(b)). This improves the stability of the system in that any perturbation which causes the atom to slow less than it should moves it closer to resonance, at which point the scattering force increases and corrects for the temporary perturbation.

The simplest way to implement this 'security detuning' is to adjust the magnetic bias field,  $B_b$ , as was done in section 2.1.1 by including a term proportional to the required total detuning,

$$\delta_{\text{sec}} = -\frac{\Gamma}{2}\sqrt{\frac{s'}{\eta} - s' - 1} = -\frac{\Gamma}{2}\left(\frac{s' - s}{s}\right)^{1/2} \tag{A.2}$$

<sup>&</sup>lt;sup>1</sup>We use the term 'deceleration parameter' and 'security parameter' interchangeably for  $\eta$ , which takes on both names in the literature. Decreasing  $\eta$  is accomplished by decreasing the magnitude of  $B_0$  (equation 2.10) on the physical apparatus.

<sup>&</sup>lt;sup>2</sup>Here we follow the notational convention of [104]. An alternative way of thinking about this situation is that we begin with a maximally efficient slower of  $\eta = \eta_{\text{max}}$ , then increase the saturation parameter from s to s' while leaving everything else unchanged.

<sup>&</sup>lt;sup>3</sup>In fact, the offset in velocity space is equal to  $\delta/k$ . In position space, the offset actually varies and is determined by the slope of the magnetic field, which maps frequency detuning onto position. Nonetheless, we call this offset near-constant in that the *vertical* displacement remains unchanged.



Figure A.1: Numerically calculated atom and scattering curve locations in position and velocity space at several times during Zeeman slowing for (a) slowing at  $\eta = \eta_{\text{max}}$  and ideal *B* field (i.e. the analytic expression of equation 2.10), (b) slowing at  $\eta < \eta_{\text{max}}$  and analytic *B* field with  $\delta_b \neq \delta_{\text{sec}}$ , and (c) slowing with the magnetic field simulated for our actual slower. Arrows indicate the motion of the atom and scattering peak in time, and dashed lines indicate  $\delta_{\text{sec}}$ , the design detuning corresponding to the chosen  $\eta$  value. For typical non-equilibrium slowing the detuning begins small (the atom is close to resonance) and eventually increases beyond  $\delta_{\text{sec}}$  at the end, though for the majority of slowing it is nearly equal to  $\delta_{\text{sec}}$ , thus justifying the 'surfing approximation' we make in the text. Note that for the real Zeeman slower, the atom actually begins on the wrong side of the scattering peak (lower right). The stability condition for the slower is that the atom not progress beyond  $\delta = +\delta_{\text{sec}} - \text{i.e.}$  the dashed line on the 'wrong' side of the scattering peak (see [104] for a more detailed treatment of this condition). Because this is not the case, the atom is still captured.

In this case,  $\delta = \delta_{\text{sec}}$  is exactly constant over the length of the slower and the velocity is equal to its design value of  $v(x) = v_c \sqrt{1 - x/x_0}$ .

If we even slightly vary  $B_b$ ,  $\Delta$ , or s' from the values assumed for this optimal configuration, though, the total detuning is no longer constant, and the velocity becomes offset from its design value in an effort to compensate for the discrepancy between  $\delta_{sec}$  and the actual detuning. Below, we present an analytical treatment of this scenario, following [104] but with our own modifications to make the treatment more general.

To begin, let us write the velocity of the atom as

$$v'(x) = v(x) - \epsilon(x) \; ,$$

where v'(x) is the actual atomic velocity, v(x) is the ideal square-root-like design velocity, and  $\epsilon(x)$  denotes the (position-dependent) velocity offset. The sign is chosen because the velocity is typically less than its design value to enforce a stable detuning  $\delta < 0.4$  Furthermore, we introduce an arbitrary 'bias offset' between the magnetic bias field and laser detuning,

$$\delta_b = \Delta - \frac{\mu'}{\hbar} B_b \tag{A.3}$$

We note that  $\delta_b = \delta_{\text{sec}}$  satisfies the steady-state scenario presented in section 2.1.1 and mentioned above, wherein  $\epsilon(x) = 0$ .

To solve for the velocity offset we begin, as always, with the scattering force equation, which gives an atomic acceleration

$$\frac{dv'}{dt} = -\frac{s'}{1+s' + \frac{4}{\Gamma^2} \left(\delta_b - k\epsilon\right)^2} a_{\max} .$$
 (A.4)

We now make the 'surfing approximation': the atom 'surfs' along the side of the scattering curve at a nearly constant offset,<sup>5</sup> and so mathematically,  $\frac{d\epsilon}{dx} \ll \frac{dv'}{dx}$  and

<sup>&</sup>lt;sup>4</sup>This is apparent from the detuning equation (eqn. 2.7, or A.1). Another way to see this: if the atom begins exactly on the scattering peak and s' > s, the increased scattering force immediately slows the atom below its design velocity.

<sup>&</sup>lt;sup>5</sup>Equivalently, the slope of the actual trajectory is nearly the same as that of the design trajectory.

 $\frac{dv}{dx} \approx \frac{dv'}{dx}$ . Equation A.4 then becomes

$$\frac{dv'}{dt} = v'\frac{dv'}{dx} \approx \frac{v'}{v}\left(v\frac{dv}{dx}\right) = -\frac{s'}{1+s'+\frac{4}{\Gamma^2}\left(\delta_b - k\epsilon\right)^2}a_{\max}$$
$$\implies \left(\frac{v'}{v'+\epsilon}\right)\left(\frac{s}{1+s}\right) = \frac{s'}{1+s'+\frac{4}{\Gamma^2}\left(\delta_b - k\epsilon\right)^2}$$
(A.5)

where in the second line we have used the fact that  $v \frac{dv}{dx} = \frac{dv}{dt} = -\eta a_{\text{max}}$ .

This is a quadratic equation for  $\epsilon(x)$ ; after a bit of algebra, we find

$$\epsilon(x) = \theta \pm \left(\theta^2 + \phi^2\right)^{1/2} \tag{A.6}$$

where

$$\theta = \frac{\Gamma^2}{8k^2} \frac{s'(1+s)}{sv'} + \frac{\delta_b}{k} \tag{A.7}$$

and

$$\phi = \left[\frac{\Gamma^2}{4k^2} \left(\frac{s'-s}{s}\right) - \left(\frac{\delta_b}{k}\right)^2\right]^{1/2} \tag{A.8}$$

From this we immediately get the result we saw previously: namely, for  $\delta_b = \delta_{\text{sec}}$ ,  $\phi = 0$  and thus  $\epsilon = 0$ . We further note that for large velocities  $v' \gg \Gamma/k$ ,  $\theta \sim \text{const}$  and therefore<sup>6</sup>

$$k\epsilon = \text{const} = \delta_b + \frac{\Gamma}{2} \left(\frac{s'-s}{s}\right)^{\frac{1}{2}} = \delta_b - \delta_{\text{sec}}$$
 (A.9)

This gives the constant velocity offset observed for the majority of slowing. Once again, it is equal to zero for the equilibrium case  $\delta_b = \delta_{\text{sec}}$ , and it only fails to capture the atom's behavior near the end of the slower when kv' becomes comparable to the width of the cooling transition. More generally, as can be seen in figure A.1, the surfing condition fails near the end of the ZS because the detuning moves rapidly from its equilibrium value, at which point the above result diverges from numerically calculated velocity profiles (see [104]).

# A.2 Zeeman Slowing to Nonzero Final Velocity

When we consider Zeeman slowing in theory (section 2.1.1), we construct a magnetic field profile by assuming that the atom is ultimately slowed to zero velocity. While for the narrow-line, low-capture velocity MOT in our experiment such behavior

<sup>&</sup>lt;sup>6</sup>Here we choose the sign that leads to stable trajectories.

would appear ideal, atoms that are brought to near-zero velocity at the end of the slower will in fact never have the opportunity to interact with the MOT. Indeed, over the  $\sim 15$ cm between the end of the slower and the center of the MOT chamber, atoms slowly scatter photons detuned some 540 MHz from the 401-nm transition. While this detuning is such that

$$\frac{a_{\rm scatt}}{a_{\rm max}} \approx 0.4\% \ ,$$

the magnitude of the acceleration is still  $a_{\text{scatt}} \sim 2000 \text{ m/s}$ . This means that an atom must have a velocity above roughly 25 m/s at the end of the Zeeman slower to successfully traverse 15 cm of off-resonant scattering and arrive at the MOT with zero velocity. Below, we develop two approaches towards achieving this goal. The first is to modify the magnetic field profile B(x); the second relies on tuning the constant bias field  $B_b$  and the laser detuning  $\delta \omega$  with respect to one another.

### A.2.1 Altered Magnetic Field Profile

Given a desired final velocity  $v_f$  and constant deceleration  $a_{\text{ZS}}$  (as before), solving for the atomic velocity as a function of position x along the slower (length  $x_0$ ) follows from simple kinematics. The result is

$$v(x) = v_0 \sqrt{1 + \left(\frac{v_f}{v_0}\right)^2 - \frac{x}{x_0}} \equiv v_c \sqrt{1 - \frac{x}{x_{\text{eff}}}},$$
 (A.10)

where  $v_0 = \sqrt{2a_{\text{ZS}}x_0}$  is the capture velocity of a slower designed to bring atoms to rest. In the second equality we have defined the capture velocity in this finite-slowing case,

$$v_c = v_0 \sqrt{1 + \left(\frac{v_f}{v_0}\right)^2} \;,$$

as well as the 'effective slower length'

$$x_{\text{eff}} = x_0 \left( 1 + \frac{v_f^2}{v_0^2} \right) = x_0 \left( \frac{v_c}{v_0} \right)^2$$

The velocity profile thus behaves like that of a slower of length  $x_{\text{eff}}$  that abruptly turns off at  $x = x_0 < x_{\text{eff}}$ .

Using equation 2.7 we can then solve for the magnetic field, which takes the same

form as in the case of slowing to rest:

$$B(x) = B_b + B_0 \sqrt{1 - \frac{x}{x_{\text{eff}}}} , \quad 0 \le x \le x_0 ,$$
 (A.11)

with

$$B_b = \frac{\hbar}{\mu'} \left( \Delta - \delta_{\text{sec}} \right) \quad \text{and} \quad B_0 = \frac{\hbar k}{\mu'} v_c$$

as before, but with  $v_c$  redefined above.

This approach is conceptually clean, but it requires that we build in a specific final velocity to our slower when constructing its field profile.<sup>7</sup> Because of the necessity of atomic-flux optimization *after* the physical construction of the slower, below we consider the effect of changing bias fields and laser detunings on the final velocity.

## A.2.2 Bias Field and Laser Frequency Tuning

In principle, if we adjust the magnetic bias field  $B_b$  and laser detuning  $\delta \omega$  such that

$$B_b = \frac{\hbar}{\mu'} \left( \Delta - \delta_{\rm sec} + \Delta \delta \right) \;,$$

the resonance condition (eqn. 2.7) reads:

$$\Delta + kv(x) - \frac{\mu'}{\hbar}B(x) = \delta(\eta)$$
  

$$\Longrightarrow \Delta + kv(x) - (\Delta - \delta_{\text{sec}} + \Delta\delta) - kv_c\sqrt{1 - \frac{x}{x_0}} = \delta_{\text{sec}}$$
  

$$\Longrightarrow v(x) = \frac{\Delta\delta}{k} + v_c\sqrt{1 - \frac{x}{x_0}} \equiv \Delta v + v_c\sqrt{1 - \frac{x}{x_0}} .$$
(A.12)

This seems to imply that the velocity at the end of the slower is  $v(x = x_0) = \Delta v$ , and so simply adjusting the relationship between the laser detuning and the bias field allows one to choose a final velocity.

Unfortunately, equation A.12 for v(x) no longer describes motion under constant acceleration, and so the total detuning varies as a function of x. Nonetheless, the principle still holds, and as long as  $B_b \neq \frac{\hbar}{\mu'} (\Delta - \delta_{sec})$  at a given position x, there will be a velocity offset (which also varies as a function of x) from the ideal square-root

<sup>&</sup>lt;sup>7</sup>Admittedly, in practice the discrepancy between this profile and that of the slowing-to-rest case is negligible. When  $v_f = 25$  m/s and  $\eta = 0.5$  (corresponding to  $v_0 = 470$  m/s),  $x_{\text{eff}} = 40.1$  cm for a 40-cm slower.



Figure A.2: Numerically calculated ZS final velocities as a function of the detuning offset,  $\Delta \delta$ .

profile. Figure A.2 shows the numerically calculated final velocities corresponding to small frequency offsets  $\Delta\delta$  and slowing parameters of  $\eta = 0.5$ , s' = 5,  $\Delta = -540$ MHz and an ideal magnetic field profile. The final velocity increases for detunings on either side of the optimal value, if not quite symmetrically.

# A.3 Narrow-Line MOT Behavior

Here we briefly present a simple model we developed to help explain some of the capture phenomena observed in numerical simulations of the narrow-line MOT behavior from section 2.2.1. Here we only concern ourselves with the atom's behavior between initial capture and its initial achievement of  $v_x = 0$  (i.e. where it first turns around), during which time it interacts almost exclusively with the laser in the -x direction, opposing its motion (initially in the +x direction). The total detuning with respect to this beam is given by

$$\Delta + kv_x + |\mu'| \frac{dB}{dx}x = \delta \tag{A.13}$$

where here  $\frac{dB}{dx}$  is positive given the negative vertical gradient in our MOT. Because of the small linewidth of the transition,  $\delta$  must be small for any appreciable slowing to occur, and approximating  $\delta \approx \text{const.} \approx 0$ , we see that

$$\frac{dv}{dx} = -\frac{|\mu'|}{k}\frac{dB}{dx} = \text{const},\tag{A.14}$$

i.e. the slowing follows a linear trajectory imposed by the B gradient in order to keep the detuning small. As with the Zeeman slower, this trajectory is stable for negative  $\delta$  insofar as perturbations in which  $v_x$  temporarily becomes larger than expected are greeted with smaller  $\delta$  and faster slowing. The linear-trajectory approximation, which we will take for the remainder of this appendix, appears good when compared with the numerically calculated results in figure 2.5a.

Atoms are captured by the MOT when they are slowed to zero velocity before 'running out of real estate' by moving beyond the spatial extent of the MOT beams. Accordingly,  $\frac{dv}{dx}$  has to be large enough to reach  $v_x = 0$  before this happens. In particular, the maximum allowed value of this derivative is

$$\left|\frac{dv}{dx}\right| \ge \frac{\Delta v}{\Delta x_{\max}} = \frac{v_0}{l/2 - x_0} \tag{A.15}$$

where  $v_0$  is the initial velocity, l/2 is half the atom's interaction length with the MOT beams, and  $x_0$  is the position at which the linear slowing trajectory  $v_x(x)$  is equal to  $v_0$ , taken to be the position at which the atom starts slowing (i.e. it initially follows  $v(x) = v_0$  for  $x < x_0$  and then  $v(x) = v_0 - \frac{dv}{dx}(x - x_0)$  at later positions). This is defined by the position at which  $\delta = 0$ , given by

$$x_0 = -\frac{(\Delta + kv_0)}{|\mu'| \, dB/dx} \tag{A.16}$$

Substituting equation A.14 for dv/dx and the above for  $x_0$ , eqn. A.15 becomes

$$\frac{dB}{dx} \ge \frac{2\left|\Delta\right|}{\left|\mu'\right| l} \tag{A.17}$$

which is eqn. 2.26 in the main text. We find decent quantitative agreement between this equation and the gradient cutoffs observed in 1D numerical simulations at s = 12for  $l/2 \approx 1.8w_0$ , where  $w_0$  is the MOT beam waist. This corresponds to the  $1/e^4$ value of the scattering force amplitude s(x)/(1 + s(x)), suggesting that the atoms experience a noticeable scattering force even down to very low laser intensities.<sup>8</sup>

We next present a straightforward (though imprecise) approach towards predicting the capture velocity. We begin by noting the equilibrium-kinematics result for  $v_c$ ,

$$v_c = \sqrt{2\eta a_{\max}l} \tag{A.18}$$

<sup>&</sup>lt;sup>8</sup>This might be explained by noting that as the intensity decreases the interaction time increases (corresponding to a slower atomic velocity), which helps offset the former.
where  $\eta = a/a_{\text{max}}$  as with the Zeeman slower. We next note that in order to achieve a particular linear slowing slope, the condition

$$\eta a_{\max} = -\frac{dv}{dt} = -v\frac{dv}{dx} = v\frac{|\mu'|}{k}\frac{dB}{dx}$$
(A.19)

must be satisfied – i.e. the acceleration needs to be large enough to allow the velocity to change as quickly or slowly as it does. We next make the approximation, for the fastest captured atoms, that on average over the slowing process  $v = v_c/2$ . Plugging this in above and solving equations A.18 and A.19 yields

$$v_c \approx l \frac{|\mu'|}{k} \frac{dB}{dx} \tag{A.20}$$

For the same values of l used above, this gives predicted capture velocities that differ by typically under 2 m/s from the 1D numerical simulations.

## Appendix B

# Atomic Beam Flux and Spatial Profile

In this appendix we consider the properties of the erbium atomic beam as it makes its way from the effusive oven source at one side of the apparatus, through transverse cooling and then Zeeman slowing, and finally towards the center of the MOT chamber. The small diameter of the Zeeman slowing/differential pumping tube places tight constraints on the range of solid angles atoms can possess out of the oven while successfully reaching the MOT chamber; here we present simple geometrical and mechanical arguments quantifying the nature of these constraints with an eye towards (a) determining the allowed atomic divergence over the beam path from oven to MOT (necessary, for instance, for matching the slowing light to the atomic beam geometry) and (b) determining the expected capturable atomic flux at the MOT chamber.

### **B.1** Spatial Profile

In an effort to roughly quantify the atomic beam size and divergence at various points in the setup, we adopt the crude model shown in figure B.1. The atoms exit the oven at an initial radius  $r = R_{out}$  from the central beam axis with a maximum divergence angle  $\theta_{out}$  and propagate along straight-line paths a distance  $z_{TC}$  until they reach the TC beams, at a maximum radius  $R_{TC}$ . During transverse cooling, the axial velocity remains unchanged while the radial velocity profile narrows, meaning that the atoms leave at an angle  $\theta_{ZS} < \theta_{out}$ . We next make a critical assumption regarding the beam profile during Zeeman slowing: if we match the convergence of the Zeeman slower beam to the divergence of the atoms, so that the ZS beam fills



**Figure B.1:** Geometry of the atomic beam divergence model discussed in the text. Atoms begin in the effusive oven on the left and propagate to the right through TC and the Zeeman slower while slowly spreading out.

the full differential pumping tube at the MOT end and focuses towards the oven, then the small radial component of the absorbed ZS photons leads to a small amount of transverse cooling. In particular, the ratio of the radial to axial photon momenta is precisely that of the radial to axial atom velocities for extremal atoms; thus both velocities are reduced in a way such that  $\theta_{\rm ZS}$  stays approximately constant [106]. Accordingly, in our model the atoms maintain a divergence angle  $\theta_{\rm ZS}$  through the entire Zeeman slower. The maximum allowed divergence angle out of the oven is determined by back-propagating the trajectory of an atom whose radius is precisely  $r = R_{\rm max}$  at the end of the slower, where  $R_{\rm max}$  is the inner radius of the ZS tube.

Onto this simple model we also add the effect of transverse heating due to spontaneous emission that arises during the Zeeman slowing [106], which gives a velocity random walk  $v_T(t) = \frac{\hbar k}{\sqrt{3m}} (\frac{\Gamma}{2} \eta t)^{1/2}$ , where the numerical prefactor reflects the average projection of an emitted photon onto the transverse axis,  $\langle \cos^2 \theta \rangle^{1/2}$ ,<sup>1</sup> and  $\eta = a/a_{\text{max}}$  is the ZS deceleration parameter. This results in a total expected radial diffusion over the full slower

$$\Delta r_{\rm diff} = \int_0^{t_{\rm slow}} v_T(t) dt = \frac{2\hbar k}{3\sqrt{3}m} \left(\frac{\Gamma}{2}\eta\right)^{1/2} t_{\rm slow}^{3/2} = \frac{2\hbar k}{3\sqrt{3}m} \left(\frac{\Gamma}{2}\eta\right)^{1/2} \left(\frac{v_i - v_f}{\eta a_{\rm max}}\right)^{3/2} \tag{B.1}$$

where  $v_i$  and  $v_f$  are the initial and final ZS velocities. We account for this diffusion by taking  $R_{\text{max}} \rightarrow R_{\text{eff}} = R_{\text{max}} - \Delta r_{\text{diff}}$  in the geometrical calculations. For a ZS with  $\eta = 0.5$  and  $v_i - v_f = 470$  m/s, we get  $\Delta r_{\text{diff}} \approx 1$  mm, so  $R_{\text{eff}} \approx 3$  mm for our geometry and slower.

We quantify the effect of the transverse cooling simply by considering the equations in section 2.1.2 of this thesis. For reasonable values of the laser power between 20 and 80mW, an optimal detuning of  $\Gamma/3 = -10$  MHz (this maximizes the damp-

<sup>&</sup>lt;sup>1</sup>The square root arises since the average velocity displacement goes as the square root of the number of scattered photons.

ing coefficient  $\alpha$ ), elliptical beam dimensions of 35 mm × 3.5 mm, and a 500 m/s atom, from equations 2.15 and 2.16 we see that the atom interacts with TC light for at most 1.5 damping time constants, meaning that the kinetic energy decreases by a factor of  $e^{-1.5}$  and the radial velocity decreases by a factor of  $\sqrt{E} \propto e^{-0.75} = 0.47$ during transverse cooling.<sup>2</sup>

The condition on the full radial displacement of the atoms is therefore

$$R_{\text{eff}} \ge R_{\text{out}} + \frac{v_i^{(T)}}{v_i^{(L)}} z_{\text{TC}} + \frac{v_i^{(T)} e^{-0.75}}{v_i^{(L)}} (z_{\text{MOT}} - z_{\text{TC}})$$
(B.2)

where  $v_i$  is the initial velocity out of the oven in the transverse (T) and longitudinal (L) direction,  $z_{\rm TC}$  is the axial position of the transverse cooling stage measured from the oven output at z = 0, and  $z_{\rm MOT}$  is the position of the *start* of the MOT chamber (i.e. the point at which atoms leave the 8mm ZS tube). Solving for the velocity ratio for extremal atoms gives

$$\frac{v_i^{(T)}}{v_i^{(L)}} = \frac{R_{\text{eff}} - R_{\text{out}}}{z_{\text{TC}} + (z_{\text{MOT}} - z_{\text{TC}})e^{-0.75}}$$
(B.3)

We can then get the beam divergences in this effective picture (under the small angle/paraxial approximation commonly seen in geometric optics) from

$$\theta_{\text{out, eff}} = \frac{v_i^{(T)}}{v_i^{(L)}} \quad \text{and} \quad \theta_{\text{ZS, eff}} = \frac{v_i^{(T)}e^{-0.75}}{v_i^{(L)}}$$

The final task is to switch from the effective picture to the real-space picture. Since all of the heating happens during Zeeman slowing, the atom propagates at  $\theta_{\text{out}} = \theta_{\text{out, eff}}$  from the oven to TC and at  $\theta_{ZS, \text{eff}}$  from TC to the start of the Zeeman slower. The difference between  $R_{\text{eff}}$  and  $R_{\text{max}}$  is exclusively made up in the Zeeman slower, then, leading to a larger divergence through the ZS given by

$$\theta_{\rm ZS} = \frac{R_{\rm max} - R_{\rm ZS}}{z_{\rm MOT} - z_{\rm ZS}} = \frac{R_{\rm max} - [R_{\rm out} + z_{\rm TC}\theta_{\rm out} + (z_{\rm ZS} - z_{\rm TC})\theta_{\rm ZS, \, eff}]}{z_{\rm MOT} - z_{\rm ZS}} \tag{B.4}$$

For  $R_{\rm out} = 0$  mm this gives  $\theta_{\rm ZS} \approx 0.30^{\circ}$  and  $\theta_{\rm out} \approx 0.38^{\circ}$ , while for  $R_{\rm out} = 1.5$  mm these numbers are reduced to  $0.21^{\circ}$  and  $0.19^{\circ}$ , respectively. In the latter case, the ZS divergence is larger than the output divergence, despite the TC stage, because

 $<sup>^{2}</sup>$ The atoms do not approach the Doppler limit in this short time interval and therefore we neglect heating effects here.

of the effect of transverse heating during ZS. In all cases it is clear that the narrow Zeeman slower tube limits the allowed atomic divergence to extremely low values.

#### **B.2** Atomic Flux

The axial velocity distribution of atoms emitted from the erbium oven is given, under the approximation that the oven aperture diameter is much smaller than the oven-bound-atoms' mean free path, by [41, 124, 125]

$$f(v,T) = \left(\frac{\pi D}{2}\right)^2 \frac{p_{\text{vap}}(T)}{k_B T} \left(\frac{m}{2\pi k_B T}\right)^{\frac{3}{2}} v^3 \exp\left(-\frac{mv^2}{2k_B T}\right) \sin^2\theta , \qquad (B.5)$$

where

$$p_{\rm vap}(T) = 10^{A - \frac{B}{C+T}}$$

is the Antoine equation for vapor pressure, v and T are the atoms' velocity and temperature, D is the oven aperture diameter,  $\theta$  is the angle under which atoms are emitted from the oven, and the empirical constants A, B, and C are given by  $A_{\rm Er} = 7.103(4), B_{\rm Er} = 12170(20), \text{ and } C_{\rm Er} = 100(2)$  for  $p_{\rm vap}$  in mbar and T in °C [41].

By integrating this equation up to the capture velocity of the Zeeman slower, using  $\theta = \theta_{out}$  from the discussion above, we can estimate the atomic flux expected at the MOT chamber. With  $v_c = 470$  m/s and a representative emission angle of  $\theta = 0.25^{\circ}$ , this gives a flux at  $T = 1200^{\circ}$ C of  $\Phi \approx 7 \times 10^{10}$  atoms per second. In practice, we have not accounted for other loss mechanisms, including decay to metastable states during Zeeman slowing (~ 50%), a broadened velocity profile after the Zeeman slower that could have a significant tail above the MOT capture velocity (this broadened profile might arise due to the Gaussian laser profile, so that there is spatial variation in the slower's effectiveness), and the fact that we only slow one isotope (at most ~ 30% of the atoms). These combined mechanisms could result in a loss of at least an order of magnitude in flux, so it is more reasonable to expect ~ 10<sup>9</sup> atoms per second at the MOT.

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