Quantum Behavior of an Atomic Fermi Gas

by

Brian DeMarco

B.A. Physics, SUNY Geneseo, 1996

A thesis submitted to the Faculty of the Graduate School of the University of Colorado in partial fulfillment of the requirements for the degree of Doctor of Philosophy Department of Physics

2001

This thesis entitled: Quantum Behavior of an Atomic Fermi Gas written by Brian DeMarco has been approved for the Department of Physics

Deborah S. Jin

Eric A. Cornell

Date _____

The final copy of this thesis has been examined by the signatories, and we find that both the content and the form meet acceptable presentation standards of scholarly work in the above mentioned discipline. DeMarco, Brian (Ph.D. Physics)

Quantum Behavior of an Atomic Fermi Gas

Thesis directed by Dr. Deborah S. Jin

Results from the production and study of the first degenerate Fermi gas of atoms are presented. By adapting the magnetic trapping and evaporative cooling techniques that were used to produce atomic Bose-Einstein condensation, a gas of fermionic 40 K atoms is cooled into the quantum regime. The fundamental difficulty in cooling a gas of fermionic, compared to bosonic, atoms is the lack of rethermalizing collisions in a spin polarized sample. This obstacle is explored in cold collision studies and then overcome by magnetically trapping two spin-states of 40 K and developing a technique for simultaneous evaporative cooling.

The ability to cool an interacting, two-component gas to one-quarter of the Fermi temperature is demonstrated. A spin-polarized, ideal gas can be cooled to as low as ~ 0.17 times the Fermi temperature. The emergence of quantum behavior at low temperature is observed both as "excess" energy in the gas and a distortion of the momentum distribution compared to the classical expectation. Furthermore, the effect of the Fermi-Dirac statistics of the gas on individual collisions via Pauli blocking is observed through measurements of the thermal relaxation time.

Acknowledgements

It is difficult to decide where, or when, to start thanking people. Perhaps it is always best to start from the beginning. I must thank my family, from my parents to my grandparents to all of my aunts and uncles, for always supporting my interests. Different family members have made a mission out of supporting both the scientific and artistic sides of my personality. This competition between the two halves of my brain has given me, I think, a unique perspective on many problems. Special thanks are deserved by my parents who were not too perturbed by the disassembly of the family Marantz stereo equipment at an early age, various fires, mixtures of random chemicals, co-opting of the patio for use as a studio, and many other exploits. My sister Nicole has been a constant source of support and a foil to my nefarious plans.

Many teachers from high school were responsible for a great deal of personal growth, and also gave me the opportunity to explore the various facets of my character. Nicki Corgel gave me free reign over the art department, overlooked the theft of a lot of paint, and always provided a safe haven. Dorothy Bahrenburg encouraged my development into a leadership position and boosted my self-esteem by allowing me to submit ideas and art into the public domain. Roger Thorpe was the source of a lot of "nerdy" and scientific fun. Most of all, these teachers and friends saved me from wilting from the overpowering boredom of high school.

In college, I received an amazing education from the Physics Department at SUNY Geneseo. Dedicated teachers like Steve Padalino, Kurt Fletcher, Jerry Reber, Ken Kinsey, Dave Meisel, and Savi Iyer invested great effort and time in developing young minds. Dave provided me with a bicycle helmet, which still protects my head today. Ken challenged me academically more than anyone had before. The image of an operasinging Kurt walking into my first quantum mechanics class is forever etched into my memory. Jerry showed me some of the in's and out's of politicking, and revealed tricks to me for dealing with people in the College Planning Council and College Senate.

Two of my college professors and friends, Steve Padalino and Stacey Edgar, deserve special praise. Steve gave me my first gig in experimental physics. Steve's love of physics was infectious. He was a close friend and taught me more about life and physics that can be summarized in a few sentences. Stacey, besides her support as a friend, made me think hard about important questions. Actually, Stacey just made me think a lot in general. And, she gave me an excuse every semester for some fresh air.

I also met Carol in college. She has been a true, steadfast, and trustworthy friend from the instant that we met. Carol has taught me much about understanding and communicating with other people. And, through Carol, I have learned more about myself than many people discover in a lifetime. I cannot imagine my life without her.

Graduate school has been the Spring of my life — a time for the growth of new things. My peers made my experience in graduate school rich. In no particular order, Dirk Mueller, Neil Claussen, Kristen Corwin, Mike Matthews, Jason Ensher, Brian Anderson, David Hall, Tim Gfroerer, Heather, Dwight, Kurt Vogel, Paul Halijan, Jonathan Goldwin, Jake Roberts, Stefan, Kurt Miller, Lu, Simon Cornish, and many others (I'm sure that I've forgotten almost everyone — don't be insulted!) have taught me volumes of physics and experimental tricks. Behind the uniquely excellent treatment of graduate students by JILA are the amazing and friendly JILA fellows, too numerous to mention here. John Bohn has been a good friend, drunk a lot of coffee and beer, and taught me all kinds of things about atoms. Murray Holland has been a source of many ideas and a constant reference. Eric Cornell and Carl Wieman have been like second advisors, and have given me great advice on physics and life from different perspectives. They surely deserve more thanks than I can give here.

I'm not exactly sure why Scott Papp has put up with me over the years. Scott has built circuits, investigated calculations, and machined parts necessary for my graduate work, all the while suffering me. Scott has also been a great friend, and has a bright future as a physicist.

The efforts of the support staff at JILA have been invaluable. The quick and efficient work of Ed Holliness and everyone in the purchasing office was absolutely necessary to my thesis work. Without equipment and tools, no science can be done. The advice and hard work of the electronics shop staff was also necessary to making the experiment function. In particular, James Fung-A-Fat has built more circuits for the experiment than I can count without more caffeine. Terry Brown has taught me almost everything that I know about electronics. Terry's careful advice and thought was crucial to making the experiment work well enough to satisfy our desires. Everyone in the machine and special techniques shop deserves my thanks as well. Seth Wieman and Hans Green have literally taught me everything that I know about machining, and they constructed essential components of the apparatus. Hans Rohner helped us to develop the enriched source that made producing the world's first degenerate Fermi gas of atoms possible, and he made our beautiful glass cells as well. The coffee brewed every day by the reception desk staff has kept me awake enough to write this thesis.

And then there is Debbie. Debbie, my graduate advisor, has been a mentor without equal. She has taught me how to write an excellent scientific paper, how to think in new ways, how to dissect and delve into the guts of a problem, and always treated me like an equal. I am sure that my life would not have turned out the same way without the opportunities given to me by Debbie. She was always generous generous with her time, her support, her thought, and the rewards of our hard work. I will treasure Debbie's friendship for all of the years that lie before me. Actually, as I trace out the pattern of my life, it seems that the wealth that I find connects primarily to a few incredible women. I will always owe Carol and the Boulder Airporter a debt of gratitude for bringing Ariana into my life. The unpredictability of life never ceases to amaze me. Ariana is a constant source of joy. Her love and support has helped me through some of the dark days and frustration of the last two years of my graduate work. Walking together with her, I fear no surprises held by the future.

Contents

Chapter

1	INT	RODUCTION	1
	1.1	In the Beginning	1
	1.2	There Were Fermions	2
	1.3	It Seemed Like a Good Idea at the Time	5
	1.4	Why the Delay for a DFG Compared to a BEC?	6
	1.5	Poor, Neglected ⁴⁰ K	7
	1.6	A Cunning Plan	10
	1.7	First Task: Find a Source	11
	1.8	Collisions — Another Magical JILA Atom	11
	1.9	Simultaneous Cooling	12
	1.10	Fun Experiments with DFG's	14
	1.11	My Crystal Ball — The Future	15
	1.12	Thesis Outline	15
2	THE	LEAN MEAN FERMION MACHINE	17
	2.1	Overview	17
	2.2	What <i>is</i> Unique to our Experiment?	18
	2.3	Double-MOT Setup	21
		2.3.1 Lasers	22

		2.3.1.1 Peak-locked Laser	22
		2.3.1.2 JILA MOPA design	26
		2.3.1.3 Trap and Repump MOPA's	42
	2.3.2	Optical setup	44
	2.3.3	K DAVLL lock	49
	2.3.4	Enriched Source	51
	2.3.5	Collection MOT	58
	2.3.6	Science MOT	62
	2.3.7	Push	72
	2.3.8	Vacuum System	77
2.4	Optica	l Pumping and Spin Composition	78
2.5	Magne	tic Trap	85
	2.5.1	Cloverleaf Trap	85
	2.5.2	Trap Depth	91
	2.5.3	Mechanical Design	91
	2.5.4	Current Servo and Switching Electronics	92
	2.5.5	Bias Field Stability	.06
		2.5.5.1 Jitter	.06
		2.5.5.2 Shot-to-Shot Stability	.09
2.6	Absorp	otion Imaging	12
2.7	Microv	vave Coil and Sources	19
2.8	Anti-G	Aravity Coil	.26
2.9	Loadin	g and Release Sequences	31
2.10	Compu	ıter Control	.33
2.11	Appen	dix to Chapter 2 1	.36
	2.11.1	MOPA mounts	.36
	2.11.2	PA Current Driver	.40

		2.11.3 Synthesizer Driver Circuit
		2.11.4 XTreme Diode Cooling
3	COI	$ISIONS - {}^{40}K IS MAGICAL! $ 146
	3.1	Overview
	3.2	Colliding Fermions
	3.3	Definition of Cross-Section
	3.4	Elastic Collision Measurement
	3.5	nelastic Collisions
	3.6	Appendix to Chapter 3
		6.1 Time dependence of the aspect ratio
4	ΕV	
4		ORATIVE COOLING 175
	4.1	Overview
		.1.1 Chapter Content
	4.2	General Cooling Considerations
	4.3	Evaporation Sequence and Performance
		.3.1 Single Frequency Stage 181
		a.3.2 $m_f = 5/2$ and $m_f = 3/2$ Cleanout $\dots \dots \dots$
		.3.3 Two-frequency Stage
		4.3.3.1 Original Two-Frequency Evaporation Results 189
		4.3.3.2 Evaporation Results circa Late 2000
		.3.4 $m_f = 7/2$ Removal Sweep
	4.4	Toy Model of Evaporation
		.4.2 Calculation Mechanics
		.4.3 Simulation Results
		.4.4 Evaporation Simulation Conclusions

	4.5	Heatir	ng Study	217
		4.5.1	(Lack of) Models	217
		4.5.2	Characterization	219
	4.6	Chapt	er Conclusions	224
	4.7	Apper	ndix	226
		4.7.1	Ramps	226
		4.7.2	Frequency Sweeps	226
			4.7.2.1 Single Frequency	226
			4.7.2.2 Two Frequency	228
		4.7.3	Landau-Zener Transitions	229
F	TUI		METDY AND THEDMODYNAMIC MEASIDEMENTS	<u></u>
Э	111		METRY AND THERMODYNAMIC MEASUREMENTS	232
	5.1	Overv	iew	232
	5.2	Every	thing You Wanted to Know About a Trapped Fermi Gas but were	
		Afraid	l to Ask	233
		5.2.1	Trapped Fermi Gases — Statistical Mechanics	233
		5.2.2	Density and Momentum Distribution	237
		5.2.3	Expansion from the Magnetic Trap	238
		5.2.4	Expanded Column Density	238
	5.3	Minim	nal Assumption Thermometry — The "Mixture" Fits	239
	5.4	TF TI	hermometry	240
	5.5	Imagin	ng and Thermometry Systematics	242
	5.6	Check	on the Expansion	248
	5.7	Energ	y Measurements	250
		5.7.1	Direct Method	251
		5.7.2	Fitting Method	252
	5.8	Source	es of Uncertainty	253

		5.8.1	Shot-to-Shot Reproducibility	253
		5.8.2	Systematic Uncertainties	254
	5.9	Experi	imental Results	255
		5.9.1	Deviation from Classical Momentum Profile — "Mixture" Fits $% \mathcal{A}$.	255
		5.9.2	Measurements of the Fugacity — Experimental Results with the	
			TF Fits	257
		5.9.3	Measured Momentum Profiles	260
		5.9.4	Measurements of Energy	262
	5.10	Appen	dix	265
		5.10.1	Image Subtraction	265
6	АТ	WO-CO	DMPONENT DEGENERATE FERMI GAS	267
	6.1	Introd	uction and Overview	267
	6.2	Two-C	Component Thermodynamics	268
	6.3	Two-C	Component Dynamics	272
	6.4	Chapt	er and Thesis Conclusion	276
	6.5	Appen	dix	277
		6.5.1	Rethermalization Model	277
			6.5.1.1 Overview	277
			6.5.1.2 Time Dependence of Δ	279
			6.5.1.3 Time Dependence of δ	284
			6.5.1.4 Functional Form for the Aspect Ratio vs. Time	286
			6.5.1.5 Determination of n and v	287
		6.5.2	Fast $m_f = 7/2$ Component Removal	288

Bibliography	280
Dibliography	203

Appendix

A LAB NOTEBOOK INDEX

295

xiii

Figures

Figure

1.1	Hyperfine structure of 40 K	8
2.1	Variation of the optical table temperature before installing the lab air	
	conditioner.	18
2.2	Light frequencies and detunings derived from the two MOPA's	23
2.3	Schematic of peak-locked laser cooling scheme.	24
2.4	Schematic of MOPA setup.	27
2.5	Schematic of interior of PA housing.	29
2.6	PA protection circuit.	29
2.7	MOPA cylindrical optics setup	32
2.8	Spatial profile of shaped and collimated MOPA beam	33
2.9	MOPA ASE power dependence on current.	34
2.10	Output power vs. current for two different MOPA's	35
2.11	Saturation of PA gain at different currents.	36
2.12	Sensitivity of the MOPA output power to the input light polarization	37
2.13	Rb saturated absorption spectrum of MOPA light	38
2.14	Comparison of MO and injected PA light on an OSA.	39
2.15	PA gain vs wavelength at 1.2 A.	40
2.16	Repump MO cooling scheme.	43

2.17	Optical layout.	45
2.18	DAVLL locking signal	50
2.19	Schematic of enriched source.	53
2.20	Yield measurement for enriched source.	55
2.21	Saturated absorption spectrum from enriched source	57
2.22	Detail of collection MOT.	59
2.23	Collection MOT fill dependence on repump power.	61
2.24	Detail of science MOT	62
2.25	Science MOT repump power dependence.	65
2.26	Science MOT lifetime.	67
2.27	Science MOT fill	68
2.28	Science MOT detuning dependence.	70
2.29	Science MOT temperature dependence on cold MOT stage detuning	71
2.30	Push duration.	73
2.31	Atom velocity from push beam	74
2.32	Collection to science MOT transfer dependence on detuning and polar-	
	ization.	76
2.33	Dependence of the fraction of atoms transferred in the magnetic trap on	
	optical pumping polarization and detuning.	80
2.34	Evaporation performance for optical pumping on different transitions.	81
2.35	Spin composition dependence on optical pumping power	83
2.36	Spin composition dependence on the cold MOT stage repump detuning.	84
2.37	Cloverleaf arrangement.	86
2.38	Determination of harmonic oscillator frequency.	88
2.39	Determination of the bias field (B_0)	89
2.40	Magnetic trap current control scheme	93
2.41	"Main" current servo.	95

2.42	High-frequency ("Pi") filter	96
2.43	Bias coil current shunt servo.	98
2.44	Control voltage circuit	100
2.45	Magnetic trap switching circuit.	102
2.46	Magnetic trap turn off	103
2.47	Magnetic trap electronics grounding scheme	105
2.48	Magnetic trap bias field stability on the evaporation timescale. $\ . \ . \ .$	108
2.49	Schematic of imaging optics	113
2.50	Magnification measurement	114
2.51	Camera focus	115
2.52	Optical transition lineshapes	117
2.53	Microwave coil.	121
2.54	Stub tuning impedence matching scheme.	122
2.55	Transmittance vs. frequency of the microwave coil	123
2.56	Position dependence of microwave field.	124
2.57	Microwave frequency synthesizer setup.	125
2.58	Absorption image of separated spin components at low temperature	127
2.59	Sample anti-gravity correction data.	128
2.60	Data corrected for the effect of the AG coil	129
2.61	Magnetic trap loading timeline.	132
2.62	Magnetic trap release sequence.	133
2.63	Computer control schematic.	135
2.64	PA double flexure mount	137
2.65	PA lens mount.	138
2.66	PA coupling mount.	139
2.67	PA current driver.	141
2.68	Synthesizer modulation input driver.	142

3.1	Example of thermalization data for cross-section measurement	157
3.2	Elastic cross sections vs. temperature	160
3.3	Dependence of σ on spin composition at $T = 9\mu K$	162
3.4	Spin exchange rate constant for $m_f = 5/2 + m_f = 7/2 \rightarrow 9/2 + 3/2$ at	
	different collision energies and 5 gauss	168
3.5	Spin exchange rate constant (calculated at 5 gauss and 20 $\mu {\rm K}$ to match	
	the experimental conditions) for $m_f = 5/2 + m_f = 7/2 \rightarrow 9/2 + 3/2$ for	
	different values of a_t	169
4.1	Microwave transitions used for evaporation.	177
4.2	Measurement of location of microwave transition frequencies in the "tight"	
	trap	179
4.3	Microwave lineshapes and evaporation regimes.	181
4.4	Single frequency evaporation trajectory.	183
4.5	End of single frequency evaporation trajectory.	184
4.6	Effect of microwave power on the end of the single frequency evaporation.	185
4.7	Fast $m_f = 5/2$ and $m_f = 3/2$ removal sweep timing	187
4.8	Microwave lineshapes of spin-polarized gases	189
4.9	First demonstration of cooling below T_F using two-frequency evaporation.	190
4.10	Evaporation data from 1999 taken in different traps	191
4.11	Evaporation trajectory from the quantum kinetic evaporation model by	
	Murray Holland.	193
4.12	Collision rate calculation from the quantum kinetic evaporation simula-	
	tion by Murray Holland	194
4.13	Evaporation trajectory after optimization of the magnetic trap current	
	servo's	196

4.14	Evaporation trajectory with different initial number, taken under the	
	same conditions as the data in figure 4.13	197
4.15	Two-frequency evaporation trajectory taken mid 2000 with a spin-mixed	
	gas	199
4.16	Evaporation performance vs. the sweep "bottom" frequency.	200
4.17	Two-frequency evaporation performance vs. the main FET ${f F4}$ drain-	
	source (DS) voltage drop.	201
4.18	Typical two-frequency evaporation trajectory from late 2000	203
4.19	Sensitivity to Δf , the difference between $\nu_{9/2}$ and $\nu_{7/2}$ during "evap2".	204
4.20	Optimized $m_f = 7/2$ removal sweep	206
4.21	Evaporation simulation dependence on microwave knife efficiency	211
4.22	Evaporation simulation dependence on microwave knife energy resolution.	212
4.23	Evaporation simulation without heating for two components.	213
4.24	Evaporation simulation dependence on Δf	214
4.25	Evaporation simulation with unequal removal probability for different	
	spin components.	215
4.26	Evaporation simulation dependence on the location of the sweep "bot-	
	tom" compared to ν_0	216
4.27	Heating of the trapped gas.	220
4.28	Heating rate dependence on properties of hypothetical Oort cloud and	
	the number of atoms	221
4.29	Heating rate dependence on magnetic trap power supply voltage	223
4.30	Heating rate dependence on spin composition at 190 nK in the tight trap.	224
51	Universal thermodynamic plate for a hermonically trapped Fermi gas	าวผ
0.1	oniversal thermodynamic plots for a narmonically trapped Fermi gas.	200
5.2	Effect of unabsorbable, uncounted light on the mixture fit thermometry.	243
5.3	The effect of saturation of the atomic resonance on imaging	244

5.4	Effect of detuning on the TF thermometry.	245
5.5	Bad effect of features in the probe beam on the mixture fit thermometry.	247
5.6	Measured TF fit widths and corresponding temperatures as a function of	
	the expansion time	249
5.7	Original data demonstrating the emergence of quantum degeneracy as	
	seen in the shape of the momentum distribution	256
5.8	Measured fugacity using the TF thermometry	258
5.9	Comparison between mixture and TF fits using the data in 5.7	259
5.10	Momentum profiles of a $m_f = 9/2$ gas at different values of T/T_F	261
5.11	Original data showing deviation in the energy of the gas from the classical	
	expectation	263
5.12	Energy in the gas as a function of T/T_F . The extra energy in the gas	
	compared to the classical case is plotted vs. T/T_F	264
6.1	Measured thermodynamics of the interacting gas.	269
6.2	Observed effect of Pauli blocking on the equilibrium thermodynamics of	200
0.2	the gas.	271
6.3	Sample rethermalization data used for the Pauli blocking measurement.	273
6.4	Collisional Pauli blocking observed in a measurement of the effective col-	
0.1	lision cross-section	975
		410

Tables

Table

2.1	PA lifetime for several experiments at JILA	41
2.2	Shutters	46
2.3	Half-wave plates	46
2.4	Miscellaneous optical elements	47
2.5	Lenses	48
2.6	Different magnetic trap parameters used in this thesis	90
2.7	Calibration of the effect of the AG coil on different dates for the $m_f = 9/2$	
	component	130
2.8	Different measurements of AG corrections for $m_f=7/2$ component	130
3.1	Triplet scattering lengths a_t in Bohr radii for collisions between potassium	
	isotopes	165
3.2	Possible spin exchange reactions for the measurement of K	166
4.1	Heating rate dependence on trap parameters before servo optimization.	222
4.2	Single frequency evaporation sweeps	227
4.3	Two-frequency evaporation sweeps.	229
5.1	TF fit parameters.	241

Chapter 1

INTRODUCTION

1.1 In the Beginning

This thesis will cover the first experiments with a degenerate Fermi gas (DFG) of atoms, specifically with a gas of ultracold ⁴⁰K atoms. There are several purposes behind this thesis, and different kinds of material presented herein. Published work on collisions and thermodynamic and dynamic measurements is highlighted. Experimental details, never published in refereed journals, are included in thesis, and will hopefully prove useful to people running the apparatus at JILA and others who may want to start new experiments. Cursory investigations, for example into inelastic collisions between ⁴⁰K atoms, that were not pursued far enough for publication are disclosed in this thesis. Also, a catalog of useful theoretical results can be found spread across several chapters. This thesis was written assuming that the reader is working in the field of atomic physics.

This introduction will give a historical perspective of our experiment and, in the process, establish the material covered by this thesis. I will also motivate our work and explain why we are excited by what we do. The journey that led to the production of the first Fermi gas of atoms is an engaging story.

Construction of the experiment began in the fall of 1997. Bose-Einstein Condensation (BEC) of ⁸⁷Rb [1], ²²Na [2], and ⁷Li [3] had been around for about two years, and the field of quantum degenerate atomic gases was starting to explode. Techniques for cooling alkali gases to ultra-low temperature (< 1 μ K) had become established and and were rapidly being perfected. A clear path to BEC involved collecting and precooling atoms from a room-temperature vapor or atomic beam using a magneto-optic trap (MOT) [4,5], transferring the atomic gas to a magnetic trap [6–11], and then evaporatively cooling [12–16] the gas by ejecting the highest energy atoms from the trap. The remaining atoms rethermalize to lower temperature via collisions.

At JILA, the double-MOT, a particularly attractive experimental implementation of this BEC recipe, had been invented by 1996 [11]. The double-MOT obviated the need for a physically large and complicated atomic beam machine, and achieved high numbers of magnetically trapped atoms with long trap lifetimes using simple MOT technology. In the early fall of 1997 there were still only three purely magnetic traps (which did not use permanent magnets) that had been proven for the production of BEC: the TOP trap [17], the "baseball" trap [18], and the "cloverleaf" trap [19].

BEC experiments were exploring exiciting physics in late 1997. However, these experiments were probing the physics of only one of the two classes of quantum particles that are found in nature. The next step in the evolution of the field of quantum degenerate atomic gases was clear: the production of a degenerate Fermi gas (DFG) of atoms.

1.2 There Were Fermions

Fermions are ubiquitous in nature. The constituent particles of matter — electrons, protons, and neutrons — are all fermions. Even protons and neutrons are the bound state of more fundamental fermions — quarks. Fermi-Dirac (FD) statistics governs the structure and behavior of diverse physical systems, from atoms and nuclei to metals to neutron and white dwarf stars. Fermi systems found in nature are generally dense and strongly interacting. The only realization of a low-density Fermi system up to the work presented in this thesis was a dilute solution of liquid ³He dissolved in superfluid ⁴He (see [20], for example).

An ultracold atomic Fermi gas constitutes a dilute system in which the interparticle interactions are weak and readily treated theoretically. Furthermore, fundamental control over the interactions is available, for example through magnetic field Feshbach resonances [21–23]. By adjusting the magnitude of an applied magnetic field near an atomic Feshbach resonance, the effective interaction between atoms can be tuned smoothly from strongly attractive to strongly repulsive (and to zero in between). This can even be done in some cases while maintaining low inelastic losses [24]. Not only is precise control over the interactions possible, but the internal (see [25], for example) and external (see [26], for example) states of the atoms can be carefully manipulated as well. An ultracold atom gas of fermions is an ideal system for quantitative study of quantum statistical effects in a controlled environment.

For a gas of bosons, the emergence of quantum behavior as the gas is cooled is sudden and marked by a phase transition at a critical temperature T_c . In contrast, the quantum behavior of an ideal (non-interacting) Fermi gas emerges gradually as the temperature T is lowered below the Fermi temperature, T_F . T_c and T_F are comparable for trapped atom gases, and depend only on the number of atoms and the spring constant of the trap. For our experimental parameters, T_F is on the order of one micro-Kelvin. This is strongly contrasted with the high density electron "gas" in a metal where $T_F \sim$ 10000 K. The phase space density, $n_{pk}\Lambda_{dB}^3$ (where n_{pk} is the peak density of the gas and Λ_{db} the thermal deBroglie wavelength), is a measure of the importance of quantum statistics in the gas. For $n_{pk}\Lambda_{dB} \sim 1$ or $T/T_F \sim 0.6$, the average inter-atomic distance in the gas is comparable to the average atomic thermal deBroglie wavelength and quantum mechanical effects become important.

As the temperature of the gas is lowered below T_F , the Pauli exclusion principle starts to dominate the thermodynamics and the dynamics of the gas. The Pauli exclusion principle is an expression of the requirements of exchange symmetry for fermions, and stipulates that no two fermions can occupy the same state of a quantum mechanical system. The atoms in the gas are forced (by the Pauli exclusion principle) to form a Fermi sea arrangement in the energy levels of the harmonic trapping potential at low T/T_F . Unlike uniform systems such as electrons in a metal, the inhomogeneous trapping potential then leads to a Fermi surface that is manifest in momentum and space [27].

Because the atoms "stack up" in the energy levels of the trapping potential, the gas has more energy at low T/T_F than one would expect classically. Classically, the equipartition theorem predicts that the average kinetic energy per particle is $\frac{3}{2}k_bT$. In fact, the kinetic energy of the Fermi gas does not approach zero as the temperature approaches zero. Rather, each atom has $\frac{3}{8}E_F$ of kinetic energy on average at T = 0(this is the result for atoms trapped in a harmonic potential — see [28] for example). The Fermi energy E_F is the energy of the highest occupied level at T = 0, and defines the Fermi temperature through $T_F = E_F/k_b$. Dynamic processes, such as collisions [29–32] or light scattering [33,34,31,35,36], that move atoms between the energy levels of the trapping potential are suppressed at low T/T_F by Pauli blocking. For example, a collision that would result in an atom state at low energy is suppressed at low T/T_F because the low energy states are already highly occupied. Photon scattering can be suppressed because the recoil energy may not be high enough compared to E_F in order to transfer an atom at low energy to a free state (found only within k_bT of E_F).

In interacting Fermi systems, a sharp phase transition can occur in addition to the smooth emergence of quantum behavior below T_F . A phase transition to a paired state for any Fermi system with attractive interactions between the constituent particles must exist (see [37], for example) at some temperature. In metals, the effective attraction between electrons is provided by phonon coupling and pairing produces superconductivity. For low T_c superconductors, the phase transition can be described by the Bardeen-Cooper-Schrieffer (BCS) theory (see [37], for example). The possibility for investigating the analogous physics in an atomic Fermi gas is a major motivator for our work. More detail on the possibility of "Cooper pairing" in a gas of ⁴⁰K atoms will be given at the end of this chapter and in the thesis conclusion at the end of chapter 6.

1.3 It Seemed Like a Good Idea at the Time

The magneto-optical trapping techniques that were crucial for creating atomic BEC work best with alkali atoms. There are only two stable fermionic alkali atoms: ⁶Li and ⁴⁰K (actually ²H and metastable ³He are also experimental candidates, although more difficult to work with).¹ Many people were interested in a DFG of ⁶Li because it was known to have a large, attractive interaction [38]. For this reason, ⁶Li seemed like a good candidate for BCS studies [39,40] in a dilute gas. However, ⁶Li was not without problems — the internal states with a large attractive interaction also suffer from enormous inelastic loss rates [38].

My graduate advisor in the summer of 1997, Eric Cornell, suggested that I quit working for him and help a new fellow at JILA, Deborah Jin, build a 40 K DFG experiment from scratch. While 40 K had some advantages compared to 6 Li, this seemed like a risky proposition. Not much was known at the time about the collisional properties of 40 K [41]. We did not know if the collision cross-section would be large enough to allow evaporative cooling to progress efficiently. Also, no one had succeeded in trapping large numbers of 40 K atoms because the natural abundance of 40 K is only 0.012%. There had not been a good method proposed for even detecting the onset of Fermi degeneracy in a trapped atomic gas. Further, we were already a few years behind the competition.

But, it seemed like a good idea at the time.

¹ Actually, ⁴⁰K β^- decays with a 1.3 billion year half-life to ⁴⁰Ca in an excited nuclear state. In a funny twist of fate, as an undergraduate student I had used the ~ 1.4 MeV emitted from the relaxation of the ⁴⁰Ca nucleus to calibrate gamma-ray detectors. These detectors were then used in an experiment to determine the absolute neutron flux from laser fusion experiments. The abundance of ⁴⁰K is high enough in concrete so that a sensitive photon detector (such as a cooled Ge:Li detector) can be saturated by ~ 1.4 MeV gamma-rays if placed next to a concrete wall or floor without shielding. Note that the lifetime of ⁴⁰K is comparable to the lifetime of ⁸⁷Rb, which is not stable either.

1.4 Why the Delay for a DFG Compared to a BEC?

Evaporative cooling is fundametally more difficult for a gas of fermionic atoms compared to bosonic atoms. Elastic collisions between identical fermionic atoms are prevented at low temperature ($T < 100 \ \mu$ K) [42] (see Chapter 3). At ultralow T, elastic collisions between atoms must occur in an s-wave channel. Not enough energy is present to overcome the centrifugal barrier for higher partial waves. However, atoms colliding in an s-wave channel have a spatial wavefunction that is symmetric under exchange of the colliding partners. The overall symmetry under exchange of the two particle wavefunction is determined by the product of the symmetries of the spatial and spin parts, and must be anti-symmetric for colliding fermions. Since completely identical fermionic atoms must have a spin wavefunction that is also symmetric under exchange, s-wave collisions are prevented.

Cold atom experiments using magnetic traps are normally done with spin polarized gases in order to avoid atom loss due to spin-changing collisions that populate untrapped states. For a spin polarized gas of fermionic atoms, evaporative cooling would fail because there are no rethermalizing collisions possible in the gas.

There were several proposals for defeating this problem. The first involved simultaneously trapping bosonic and fermionic atoms, for example ⁶Li and ⁷Li [43]. In this scheme, collisions between the bosons and fermions are allowed at low T and provide the rethermalization mechanism for evaporative cooling. In another approach, Geist and co-workers [44] had proposed enhancing the p-wave collision cross-section at low T by applying a static electric field in order to permit collisions between identical fermionic atoms. Our choice was to magnetically trap multiple spin states of ⁴⁰K [31]. At low T, s-wave collisions are allowed between atoms in different Zeeman levels because the two-atom spin wavefunction can be anti-symmetric under exchange.

Aside from complications caused by the collisional properties of fermionic atoms,

the thermodynamic properties cause trouble for evaporative cooling too. Efficient, or "runaway", evaporative cooling relies on an increase in the collision rate in the gas even as atoms are removed in order to cool to lower temperature. Normally, this increase in collision rate comes from an increase in density caused by compression from the harmonic trapping potential. However, the density for a trapped Fermi gas starts to "freeze" as the gas is cooled into the quantum regime [27]. Even worse, Pauli blocking of collisions drives the collision rate to zero at T = 0. So as a Fermi gas of atoms is evaporatively cooled into the degenerate regime, the collision rate in the gas must decrease and evaporative cooling must become more difficult.

1.5 Poor, Neglected ⁴⁰K

Potassium was the last alkali atom to be trapped in a MOT [45]. Some efforts had been made to trap 40 K [46,47], but the number of atoms were severely limited by the low natural abundance (0.012%). The hyperfine structure of 40 K is a little strange, with both the excited states and ground states inverted (see figure 1.1) [46,48]. This complete inversion is caused by a large nuclear spin (I=4) that points in the opposite direction from the nuclear magnetic moment. As shown in figure 1.1, the separation between excited states is relatively small and the ground states are only separated by ~ 1 GHz.



Figure 1.1: Hyperfine structure of 40 K.

The hyperfine structure of ⁴⁰K is its biggest advantage over ⁶Li for a DFG experiment. The large angular momentum of the lower ground state, f = 9/2, means that there are many magnetically trappable Zeeman levels (nominally 5). In the lower ground state, the positive m_f levels are the weak magnetic field seekers and can be magnetically trapped. Furthermore, a mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms in the f = 9/2 manifold is stable against spin changing collisions. Collisions between an $m_f = 9/2$ and an $m_f = 7/2$ atom cannot conserve the projection of angular momentum without the final state being identical in spin. Inelastic (or any) collisions between two $m_f = 7/2$ atoms are prevented by the FD statistics at low T. Further, hyperfine changing collisions are prevented since not enough energy is available below 10 mK to drive $f = 9/2 \rightarrow f = 7/2$. For comparison, there is no combination of Zeeman states in ⁶Li that is stable against spin changing collisions at reasonable values (less than many hundreds of gauss) of magnetic field.

So, our plan was to use a stable mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms to provide the rethermalizing collisions needed for evaporative cooling. Using multiple spin states is a great advantage compared to trapping multiple isotopes, since the experimental setup is considerably simplified. Further, the optical transition for potassium is 767 nm, which is very close to 780 nm in rubidium. This allowed us to use all of the optics and diode laser technology that had been developed for experiments with rubidium.

The total angular momentum, f = 9/2 for the states that we use, defines 40 K as a fermion. The total angular momentum is the sum of the intrinsic nuclear and electron spin and the electron orbital angular momentum. For the regime of temperature (or energy) in which we work, the quantum statistics of the internal components (the electrons, protons, etc.) of 40 K is not probed and is irrelevant.

The astute reader will also note that the equilibrium state of potassium is a solid at room temperature, and is most definitely a solid at 1 μ K! Experiments with quantum atomic gases rely on the fact that the process of forming a solid (via molecule and cluster formation) takes a long time at typical densities ($\sim 10^{13} \text{ cm}^{-3}$). This separation of solidification and experimental timescales allows us to perform experiments that can take many minutes.

Potassium also has two stable bosonic isotopes, ³⁹K and ⁴¹K, and so naturally lends itself to future studies of Bose-Fermi mixtures. These bosonic isotopes were also our collision backup plan. Although the collisional properties of potassium were not well known in 1997, it seemed like some combination of different spin states and/or isotopes had to have a large enough elastic collision cross-section for evaporative cooling to work well.

1.6 A Cunning Plan

Our attitude was to treat the technology for producing ultra-cold atom gases and atomic BEC as established. We planned to borrow as much as we could and invent as little as possible. We chose the double-MOT technique and a cloverleaf magnetic trap for our experimental apparatus. Potassium MOT's tend to be hot (~ 100 μ K, although one group has observed sub-Doppler cooling in ⁴⁰K [49]), and a Ioffe-Pritchard (IP) type magnetic trap [50] (compared to the TOP trap) had enough trap depth to efficiently capture and compress the gas from the MOT. Out of the two IP traps available at the time, we chose the cloverleaf configuration because it seemed to have the greatest amount of optical access.

With the help of the excellent support staff at JILA, and the members of Eric Cornell's and Carl Wieman's groups, we were able to quickly progress to the point of magnetically trapping and evaporatively cooling gases of 40 K atoms. The apparatus that we built is described in Chapter 2 of this thesis.

1.7 First Task: Find a Source

Our first task was to develop an enriched source of 40 K atoms for the MOT. High numbers of atoms are required so that the elastic collision rate in the gas is large enough for efficient cooling and so that T_F is at an experimentally attainable temperature. Eric Cornell's ⁸⁷Rb BEC experiment at JILA was using alkali dispensers from SAES, Inc. as a ⁸⁷Rb vapor source, which was an attractive option for us. These dispensers produced an alkali vapor in the vacuum system via *in situ* chemistry. Not only were the dispensers free enough of contaminants to be suitable for use in ultra-high vacuum, but they could be turned off when not in use (thereby extending the useful lifetime of the device).

Unfortunately, commercial vendors were unwilling to create potassium alkali dispensers that were enriched in ⁴⁰K. With the help of Hans Rohner at JILA, in 1998 we were able to make our own enriched alkali dispensers [51] (details in chapter 2). These sources consisted of a fine powdered mixture of KCl and Ca held in a small Nichrome foil "boat". The KCl is enriched to ~ 5% in ⁴⁰K. Electrical current is used to heat the "boat" and drive the reaction KCl+Ca→K+CaCl. These sources allowed us to trap 10000 times more atoms (~ 10^9 atoms) than previous efforts.

1.8 Collisions — Another Magical JILA Atom

Next, we turned to investigating the collisional properties of the states $(m_f = 9/2)$ and $m_f = 7/2$ in the f = 9/2 ground state) that we planned to use for cooling. We were able to transfer atoms in these Zeeman levels into the magnetic trap, and evaporatively cool (this was not yet optimized) to 10 μ K in order to sample a range in temperature. Sampling a range in temperature is equivalent to exploring a range of collision energies.

Using cross-dimensional rethermalization techniques, we were able to measure the s-wave and p-wave collision cross-sections from 5 to 100 μ K (Chapter 3) [42] in 1999. The simultaneous s-wave and p-wave measurement allowed us, in conjunction with John Bohn, Jim Burke, and Murray Holland, to accurately measure the triplet scattering length $a_t = 157(20)$ a₀, where a₀ is the Bohr radius. Not only was the triplet scattering length large (even compared to ⁸⁷Rb) but we found a p-wave shape resonance at ~ 280 μ K. This was great news — the elastic collision cross-section was high enough, considering the initial density and collision rate in the magnetic trap, such that evaporative cooling should be very efficient. P-wave collisions could also give the initial stages of cooling at high temperature a significant boost in collision rate. The measurement of the p-wave collision cross-section dependence on temperature was also one of the first measurements of the Wigner threshold law for neutral scatterers [42,52].

We also did a cursory investigation of inelastic loss processes (Chapter 3). By holding a combination of $m_f = 7/2$ and $m_f = 5/2$ atoms in the magnetic trap and looking for the production of $m_f = 9/2$ atoms, an upper limit on the spin-exchange rate constant was inferred. We found that the rate constant was unusually low ($< 10^{-14} \text{ cm}^3/\text{s}$). This was also exciting news. It was *not* crucial to load a very pure mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms into the magnetic trap. Nor was it important to avoid creating a population of other spin states during evaporative cooling. ⁴⁰K turned out to be a serendipitous choice of atoms from the point of view of collisional properties.

1.9 Simultaneous Cooling

Evaporative cooling of two components requires both m_f and energy selectivity in removing atoms. Maintaining a nearly equal mixture of components is important to ensuring a high overall collision rate at low T, since thermal equilibrium in one component is only maintained by collisional contact with the other. Removing energy from the gas in a balanced way (from both components equally) is necessary to achieve a high cooling rate. If one component was used to sympathetically cool the other, eventually the cooling would be spoiled by the heat capacity carried by the other component.

The radio-frequency (rf) transitions normally used for evaporative cooling of mag-

netically trapped gases could not give us component selectivity except at *very* high magnetic fields. We therefore chose to remove atoms via microwave transitions between hyperfine ground states. Selectivity in energy is possible because of the dependence of the transition frequency on magnetic field — the microwave frequency is set to be resonant with atoms that have enough energy to "roll" up the harmonic potential to high field. However, working at these higher frequencies (~ 1 GHz instead of \sim MHz) brought many complications (see Chapter 2). Among other technical issues, delivering high power at ~ 1 GHz to the atoms was tricky.

By mid 1999, we had developed an evaporation scheme that we dubbed "simultaneous cooling" [53] (see Chapter 4 for the details). By adjusting the magnetic trap parameters during the evaporation and using multiple microwave frequencies, we were able to always remove atoms equally from both components. We discovered a rather startling "feature" of the cooling — the evaporation became inefficient at some critical T/T_F below which we could not cool. In our first experiments [53], one component was removed in an evaporative way at the end of the bulk of the cooling, leaving a spinpolarized gas. This final sympathetic cooling, which was necessary to reach degeneracy, allowed us to cool a spin polarized gas to $T/T_F = 0.5$. At that time, we attributed the cooling failure to the properties of a trapped Fermi gas (namely, that the collision rate in a trapped Fermi gas must eventually decrease in the degenerate regime).

In late 1999, Murray Holland completed an simulation [54] which indicated that there was no fundamental limit to our cooling scheme. Spurred on by this knowledge, we began to weed out technical limits and improve the cooling. A toy-model of the evaporation that included technical limits not used in Murray's model proved useful for this task (see Chapter 4). Improvements to the magnetic trap stability and microwave instrumentation along with several small enhancements allowed us to cool a two-component, 50/50 mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms to $T/T_F \sim 0.2$ by late 2000. Now, in early 2001, we believe that the evaporation is still technically limited. Chapter 4 thesis provides many details on our evaporative cooling scheme.

1.10 Fun Experiments with DFG's

We published the first experiments with spin polarized DFG's in *Science* magazine [53]; that work was later selected as one of the top ten scientific breakthroughs of 1999 by *Science*. We were able to detect a deviation (from the classical expectation) in the momentum profile of a spin polarized ⁴⁰K gas released from the magnetic trap that was consistent with FD statistics (Chapter 5).² In this article we also published a measurement of the mean energy per particle in the gas. We measured "excess" energy in the gas at low T/T_F compared to the classical case (Chapter 5). The measurements published in [53] covered the classical regime $(T/T_F \sim 200)$ down to $T/T_F = 0.5$. Finally, both classes of quantum particles were experimentally available to workers in the field of quantum degenerate gases.

By 2000, we had perfected a thermometric technique that directly probed the Thomas-Fermi (TF) shape of the gas [55] (see Chapter 5). With the improvements in our cooling, we moved on to doing experiments with two-component degenerate gases. We also added an "anti-gravity" coil to the experiment which allowed to us spatially separate and simultaneously image the two-components (Chapter 2).

In late 2000, we published work [32] on the thermodynamics and dynamics of spin mixed degenerate gases (Chapter 6) at temperatures as low as $T/T_F \sim 0.2$. By tuning the Fermi energy of one component relative to the other (via controlling the spin mixture), we measured an imbalance in energy between the two components in the degenerate regime. We were also able to measure Pauli blocking directly by measuring the effective collision cross-section in the gas, again using cross-dimensional rethermalization. At $T/T_F \sim 0.25$, we were able to measure a factor of two reduction in the

 $^{^2}$ We had half hoped to discover that $^{40}{\rm K}$ was a boson, which would have been a *really* big scientific breakthrough.

effective collision cross-section compared to the classical regime.

1.11 My Crystal Ball — The Future

There are many experiments left to do with DFG's. Predictions exist for the emergence of a zero sound mode [56], suppression of light scattering [33,31,35,36], component separation [39,57], changes in the frequency and damping rates of collective excitations [58–63], and shell structure [64,65]. Also, the outlook for cooling further into the degenerate regime in our system is still bright (see Chapter 4).

Perhaps the most exciting possibility for future experiments is the prospect of a paired state at low temperature, similar to BCS superconductivity. The attractive interaction that is required for a phase transition to a paired state can be via a direct, collisional interaction (see [39,40] among *many* others), induced via photo-associative methods [66], or phonon mediated [67]. In early 2001, a DFG in ⁶Li was achieved by two groups (see [68] and C. Salomon, unpublished). With its large, attractive collisional interaction, ⁶Li seems like a good candidate for BCS studies. However, the inelastic losses for the relevant states are so large that one of the groups is already considering (see [68]) using other states and modifying the interaction with a Feshbach resonance.

There is a predicted [69] (and experimentally accessible) Feshbach resonance in 40 K for the $m_f = -9/2$ and $m_f = -7/2$ Zeeman levels in the f = 9/2 ground state. Unpublished work by Murray Holland predicts a phase-transition at $T/T_F \sim 0.5$ for these states using the Feshbach resonance to make the interaction strongly attractive. The 40 K group at JILA is currently setting up an optical trap with the hopes of trapping these states and accessing the Feshbach resonance.

1.12 Thesis Outline

Chapter 2 will give many details on the experimental apparatus, including information on stabilization of the magnetic trap, microwave equipment, lasers, optical pumping, and the "anti-gravity" coil. Chapter 3 covers the properties of collisions between fermionic atoms, a measurement of the s-wave and p-wave collision cross-sections for 40 K as a function of temperature, and a cursory measurement the of spin-exchange collision rate constant for 40 K. Next, in Chapter 4, many details of our cooling scheme are explained, evaporation results are highlighted, a toy-model of the evaporation is developed, and the results from a study of heating of the atoms held in the magnetic trap will be given. Chapter 5 contains a laundry list of useful thermodynamic equations, details on the techniques that we use for measuring thermodynamic properties, as well as the results of the first thermodynamic measurements of a degenerate Fermi gas of atoms. Chapter 5 will also address the sources of random and systematic uncertainty that are important in the experiment. Finally, Chapter 6 contains the results of twocomponent degenerate Fermi gases. The table of contents for this thesis is extensive, and should be used to find specific information.
Chapter 2

THE LEAN MEAN FERMION MACHINE

2.1 Overview

This chapter will describe the apparatus that was used to execute the experiments covered by this thesis. We built the apparatus, from the ground up, in the vacant space left after evicting some theorists. We chose to construct a double-magneto optic trap (double-MOT) apparatus for producing ultra-cold atom gases, since it was a proven system for Bose-Einstein condensation (BEC). Careful attention will be paid in this chapter to explaining the modifications that we had to make to the standard double-MOT scheme in order to make it compatible with cooling fermionic ⁴⁰K atoms. This chapter will also serve as a guide (hopefully useful to new graduate students and post-docs) to each piece of the experimental apparatus, highlighting any unique or essential features to producing a degenerate Fermi gas (DFG). For more information, the reader may want to refer to Jason Ensher's thesis [70] which describes an ⁸⁷Rb BEC apparatus very similar to our ⁴⁰K apparatus in form and operation.

Special mention should be made of the air conditioner (KoldWave 2K260812) that JILA supplied for the lab. The experiment could not produce a DFG reliably without active stabilization of the ambient lab temperature. Variation of the optical table temperature (before installation of the air conditioner) over a one and a half month period is shown in figure 2.1. The magnetic trap bias field and the optical alignment were affected significantly by changes on the order of 3 °F. A HEPA air filter (Kenmore

51500) has also proven useful for slowing the accumulation of dust on lenses and mirrors, which is a significant problem in Boulder's dry climate. On the other hand, Boulder's aridity is essential to our simple technique of cooling laser diodes that uses enough duct tape to make Red Green proud.



Figure 2.1: Daily variation of the optical table temperature, measured with a thermocouple, is plotted over 1.5 month period.

2.2 What *is* Unique to our Experiment?

Areas where our experiment deviate from the standard BEC apparatus include the laser system, the atomic vapor source, the optical pumping parameters, the attention to the magnetic trap stability, use of an "anti-gravity" coil, and our evaporation technique. Our experiment was the first all semiconductor laser trapping and cooling apparatus for potassium. MOT's for potassium require high power in two laser frequencies because of the relatively small excited state splittings [46,45,47]. To provide this light, we developed Master Oscillator Power Amplifier (MOPA) laser systems which are still a new technology to atom trappers. Standard locking saturated absorption spectroscopy locking techniques produce a small error signal for 40 K because of the low natural abundance (0.012%). We therefore implemented the first dichroic atomic vapor laser lock (DAVLL) [71] for potassium so that we could have a large locking signal with a large locking bandwidth.

We also had to develop an enriched source in order to trap large numbers of 40 K atoms. Efficient evaporation requires high numbers of atoms in order to have a sufficient collision rate in the magnetic trap. Further, reaching degeneracy at a reasonable absolute temperature (~ 1 μ K) relies on trapping many atoms since T_F depends on N. Our source allowed us to trap 10000 times more 40 K atoms than previous efforts, which were limited to using natural abundance sources [46,47].

Evaporative cooling requires a high rate of rethermalizing collisions in the gas. Typically, cold atom experiments that use magnetic traps employ optical pumping to prepare spin polarized samples. The atoms are all pumped into one Zeeman level in order to avoid loss due to spin changing collisions that populate magnetically untrapped states. However, collisions between ⁴⁰K atoms at low temperature occur only between atoms in different internal states. We therefore had to develop a method for reproducible, "imperfect" optical pumping in order to prepare the optimal spin polarization for the experiment.

Because potassium MOT's are relatively hot, a magnetic trap with a high trap depth is required to efficiently capture and hold the gas for evaporative cooling. To this end we chose a Ioffe-Pritchard, cloverleaf style magnetic trap [19] for our experiment. The Ioffe-Pritchard type of trap has a great disadvantage, though — the small (~ 1 gauss) bias field is created by the subtraction of two large (~ 250 gauss) fields and is therefore difficult to stabilize. The magnitude of the bias field sets the final temperature for a fixed frequency evaporative cut. Our requirements on temperature stability therefore necessitate very low noise on the bias field. We invested a great effort in stabilizing the bias field to reduce fluctuations to the order of 1 mG.

A method is required to separately image the different spin states in order to extract information about the both spin components in the trapped gas. Optical methods [72] that are used in other experiments to resolve atoms in different hyperfine states are not suitable for resolving atoms in different Zeeman levels. We therefore implemented an "anti-gravity" coil that spatially separates the spin components during the expansion [73] by exerting spin dependent forces via a magnetic field gradient.

Unlike other experiments, we cannot use transitions between Zeeman levels for evaporation because these transitions cannot distinguish between the different spin components except at *very* high magnetic field. We therefore use transitions between the hyperfine ground states to remove atoms from the trap. Using this type of transition is difficult because of the relatively high frequency and the need for two distinct frequencies. We developed techniques to deliver high power in two frequencies to the atoms in the magnetic trap.

There are some aspects of the apparatus covered in this chapter that are not unique to our experiment, for example the details of the MOT's and the transfer between MOT's. These sections of the experiment are covered for completeness. There are many figures in this chapter that use arbitrary or relative units. The purpose of these figures is to illustrate the general dependence of the behavior of the system on adjustable parameters, and to make clear which "knobs" in the experiment can be used to fix problems. The quantitative behavior of the apparatus tends to change over time, while the qualitative behavior does not. The lab book index (in the Appendix to this thesis) can be used to find specific information in more useful experimental units.

2.3 Double-MOT Setup

The double-MOT technique [11], developed at JILA, has now been employed to produce ultra-cold atom gases and atomic BEC around the world. The defining feature of a double-MOT apparatus is the use of two MOT's, one in a high vacuum section where large numbers of atoms are collected from a room-temperature vapor and the other in an ultra-high vacuum section where experiments are carried out. The separation of the "collection" MOT from the "science" MOT permits experiments with high numbers of atoms and long magnetic trap lifetimes. A schematic of our apparatus is shown in figure 2.17, with details appearing in figure 2.22, and figure 2.24. These figures will appear in sections 2.3.2, 2.3.5, and 2.3.6.

Atoms are gathered and cooled in the collection MOT from a room-temperature vapor [4,5]. These atoms are then pushed using a pulse of resonant light (the push beam — green in figure 2.17) down a small-diameter transfer tube and caught in the science MOT. The low conductance transfer tube creates the pressure gradient between the collection and science MOT's. Permanent magnets in a hexapole arrangement surrounding the tube confine the atoms and support the atoms against gravity. Once enough atoms are transferred into the science MOT, the gas is optically pumped (brown beam in figure 2.17) and loaded into the magnetic trap. The gas is then cooled evaporatively by forcing the highest energy atoms to leave the trap via microwave driven transitions between the hyperfine ground states. After cooling the gas and performing any experiments, the gas is destructively imaged by quickly turning off the magnetic trap and allowing the gas to expand (typically for 1-20 ms). Information, such as temperature, number, and energy about the gas is extracted from absorption images. These images are generated by illuminating the gas with a pulse of resonant light (the probe beam — violet in figure 2.17). The shadow cast by the gas into the probe beam is imaged onto a CCD camera, and a picture of the column density, or optical depth, is generated. The whole experimental cycle, taking ~ 3 minutes, then begins again.

2.3.1 Lasers

Three semiconductor lasers are used in the experiment, one ("peak-locked" laser) as a frequency reference and the others ("trap" and "repump" MOPA's) for manipulating atoms. The trap and repump MOPA's are used to produce the five light frequencies shown in figure 2.2. Significant power for the MOT repumper [47,46,45] is required because of the small excited state hyperfine splittings of 40 K.¹ The power and detuning in the MOT beams are optimized to ultimately maximize the number of atoms transferred into the magnetic trap. Discussion of this optimization and the optimization of the push and optical pumping parameters can be found in sections 2.3.6, 2.3.7, and 2.4 of this chapter.

Using diode lasers with ⁴⁰K is complicated by the lack of commercially available diodes at 767 nm. We therefore use wavelength selected², actively cooled 780 nm diodes. The best results have been obtained with 30 mW Sharp diodes (LT024MD); Sanyo (DL7140-201) and Mitsubishi (ML60125R) diodes have produced mixed results and generally have a very restricted external cavity tuning range.

2.3.1.1 Peak-locked Laser

The peak-locked laser is a JILA design external cavity diode laser (ECDL) [74] that was modified for diode cooling. The light from a commercial laser diode is collimated with a lens (Navitar Industries DO-818), and adjustable feedback is provided by a Littrow configuration cavity formed by a diffraction grating (Edmund Scientific NT43-215) and the rear facet of the diode. The grating provides wavelength control

¹ A ⁴⁰K MOT behaves somewhere in between a rubidium and a ³⁹K MOT — not quite a MOT with separate trap and repump frequencies, and not quite a MOT with two trap frequencies. We run the collection MOT, for example, in the regime where the calculated excited state population in the f' = 11/2 state is only 1.5 times higher than in the f' = 9/2 state.

 $^{^{2}}$ Optima Precision and Power Technology are two companies that are happy to wavelength select diodes for a small fee.



Figure 2.2: Light frequencies and detunings derived from the two MOPA's. The color scheme matches that used in the schematic in figure 2.17. More details on the hyperfine structure of 40 K (shown here) can be found in chapter 1.

and is glued onto a JILA optical mount so that it can be adjusted by hand with a ball-driver or electronically via a piezo-electric transducer (PZT). The laser has two stages of thermo-electric cooling ("TEC" will be used for "thermo-electric cooler" in this chapter) (figure 2.3) that are capable of cooling to diode to -20 °C, although the laser typically runs slightly below 0 °C. Coarse temperature control is used to tune the center wavelength of the diode within the range of the external cavity (typically ± 3 to 4 nm). Fine manipulation of the temperature is used to center a cavity mode on the atomic transition. For a more complete discussion of the cooling scheme, see the appendix to this chapter.

The peak-locked laser is frequency locked to the ³⁹K f = 2 ground state to f' excited state manifold transition.³ The laser current is dithered at 100 kHz, and the

 $^{^3}$ The excited state hyperfine structure of $^{39}{\rm K}$ cannot be resolved using standard saturated absorption spectroscopy.



Figure 2.3: Schematic of peak-locked laser cooling scheme. One inch thick Armaflex insulation that covers the laser housing is not shown. Also missing from the diagram is a small, interior box of dessicant (T.H.E. brand) that prevents condensation. The laser diode and diode mount are electrically grounded through the diode current driver.

resulting locking error signal is derived from a saturated absorption spectrum analyzed by a lock-in amplifier. Standard JILA electronics are used to frequency-lock to the zero crossing of the lock-in amplifier signal via feedback to the laser PZT. The light from the laser is combined with the light from either the trap or repump laser and focused onto a fast photodiode (New Focus 1601). The beat frequency between the lasers is measured by a frequency counter (Hewlett-Packard 53181) that is connected to the fast photodiode, and is used to set and monitor the frequency of the trap and repump MOPA's. The detunings ($\Delta_o, \Delta_t, \Delta_r, \Delta_r^{jump}, \Delta_t^{cold}$) of the trap and repump lasers will either be specified by the beat frequency in this thesis or by the actual detuning from the atomic transitions (figure 2.2). The relevant laser beams are shifted relative to the beat frequency by the acousto-optic modulators (**AO1** and **AO2**) shown in figure 2.17.

2.3.1.2 JILA MOPA design

Semiconductor lasers bring many benefits to atom trapping experiments. Although diode lasers are easy to frequency lock and to align and require little maintainence, they are typically low power (< 50 mW). While power at this level is sufficient for trapping large numbers of Rb atoms (for example), trapping potassium requires significantly higher power. At JILA, we have developed a Master-Oscillator (MO) Power-Amplifier (PA) (MOPA) laser design capable of providing high power ($\sim 500 \text{ mW}$) along with the standard benefits of diode lasers.

The JILA MOPA setup was designed under the guidance of Eric Cornell, and has now been integrated into multiple experiments involving Rb and K. A schematic of the setup is displayed in figure 2.4 and figure 2.5. The MOPA is a two-part laser [75,76] with a tapered power amplifier (PA) that amplifies light from an ECDL master oscillator (MO). The PA itself is a commercial semiconductor device that we integrate into a housing and the MOPA setup. Commercially available PA's are able to produce 500 mW of light from 5 mW of input light, and almost perfectly reproduce the frequency characteristics of the MO. Technical drawings for the PA mounts can be found in the appendix.

The PA is a semiconductor AlGaAs heterostructure device with a gain region that is tapered in order to permit uniform power density amplification. Light is injected into the input facet, freely diffracts along the gain region, and exits from the output facet at the opposite end of the device. The taper is defined by a (typically) 5×1 micron input facet and a 100×1 micron output facet, with a 1 mm long gain region. Production of PA's has been halted by SDL Inc., but they are now available for purchase from TUI Inc. Our experiment only used PA's produced by SDL. The only apparent difference between devices from SDL and TUI is the orientation of the polarization relative to the beam geometry.⁴ SDL devices have the same polarization orientation as a diode laser

⁴ This has been confirmed in preliminary tests by Dirk Müller, at JILA.



Figure 2.4: Schematic of MOPA setup.

(parallel to short beam direction), while the polarization of the TUI device is rotated by 90°. The PA is rather delicate as it is a bare semiconductor chip mounted in an SDL C-mount. Precautions must be taken to avoid physical damage to the PA. In order to prevent dust from attaching to the PA facets, the PA is housed in a sealed aluminum box with anti-reflection coated windows covering the input and output holes (figure 2.5).

The manufacturer stipulates standard handling procedures for the PA. Physically handling the device should always be done while taking anti-static precautions. Under no circumstances should cleaning of the device be attempted. Special AR coatings are used on the input and output facets (< 0.1 % reflectivity), and any damage to or accumulation of dirt on these facets results in hampered performance or device death.

The PA is electrically similar to a laser diode, and must be protected from static discharge. The PA is driven by a stable current source that is a modified JILA diode driver (see the appendix to this chapter for a circuit diagram). A protection circuit (figure 2.6)⁵ is connected between the current driver and PA as close as possible to the PA. A typical PA drops ~ 2.5 V while operating at 1-2 A with roughly a factor of 100 optical power gain. Electrically, the C-mount is the device anode and a separate pin connects to the device cathode.

⁵ This circuit was designed by Leo Hollberg's group at NIST Boulder.



Figure 2.5: Schematic of interior of PA housing.



Figure 2.6: PA protection circuit. The diodes prevent transient forward and backward voltages across the PA. The linear elements form a high frequency filter.

Temperature stabilization and heatsinking of the PA is required to prevent overheating, and can be used to tune the device characteristics. A schematic of the PA housing (figure 2.5) shows the TEC's used for cooling the device. The device should be kept below 25 °C (recommended by the manufacturer), and we measure weak tuning of the output power vs. temperature. The center wavelength of the PA tunes with temperature much like a diode laser, with a measured ~ 0.3 nm shift per degree Celsius. A proportional-integral-differential (PID) servo is used to stabilize the PA housing baseplate temperature via feedback to the TEC's. The temperature is measured using a small thermistor that is glued into a small hole in the baseplate.

The simple optical setup of a MOPA (figure 2.4) is certainly one of its advantages over other high power laser systems. We have had the best results using JILA optical mounts on large (1 inch diameter) posts. All of the MOPA optical mount posts are glued to a 0.5 inch thick sanded and lapped aluminum plate, which rests on sorbathane pads on top of the optical table. Two mirrors are used for the MOPA injection alignment. Optical isolators (40-60 dB isolation total) are needed to prevent optical feedback into the MO from PA input facet amplified spontaneous emission (ASE). A beam wedge is useful for picking off MO light for a frequency lock and to beat against the peaklocked laser. The alignment of the MO onto the PA input facet can be accomplished by following a standard procedure. The coupling lens inside the PA housing is adjusted to collimate the ASE from the PA input facet. Two mirrors are then used to overlap the MO beam and this ASE output over an extended path — if this is done well enough, amplified light should be observed from the PA output. The amplified power can then be optimized by adjusted the two mirrors and the distance of the coupling lens from the PA input facet (one half of the double-flexure mount is used for this). The PA output shows no steering with changes in the alignment of the input beam from the MO.⁶

⁶ There have been two PA's at JILA that have showed weak steering of the PA output beam with changes in the MOPA alignment. These two PA's had a very short lifetime, and are considered abnormal.

Reflection back into the PA output facet must be prevented, so all the optics following the PA up to the PA isolator should be slightly tilted. Variation in the optical properties of the amplified light among different devices necessitates implementing slightly different beam shaping optics on a case by case basis. The beam emitted from the PA after the spherical collimating lens has a large aspect ratio ($\sim 10:1$) and is quite astigmatic. The spherical collimating lens inside the PA housing is used to collimate the long direction of the beam. Cylindrical optics outside of the housing then collimate the short direction and shape the beam into a roughly 1:1 aspect ratio. The spherical collimating lens introduces a focus into the short direction 10-13 cm from the lens. Figure 2.7 shows the most commonly used (in JILA) cylindrical lens scheme, which is the scheme we use for the trap and repump MOPA's. In general, trial and error is used to find a collimation scheme that works for a particular device. Note that the output facet ASE has a very different spatial mode compared to the amplified MO light, and the collimating optics should be set up using an injected device.

All PA's are similar, with the gain, current dependence, input power dependence, and spatial mode varying somewhat from device to device. A compilation of the characteristics of several devices follows. Cross-sections of the spatial profile of a shaped and collimated beam is shown in figure 2.8; note that this device had a particularly clean mode. Some PA's have much more high-frequency spatial variation, and can show clear, high contrast stripes in the far-field spatial pattern. The dependence of optical ASE power on current is shown in figure 2.9 for both the input and output facets. The input facet ASE power can be equal to the injection power ($\sim 5 \text{ mW}$) at high current, explaining the need for the high degree of optical isolation between the MO and PA. Most of the output facet ASE is presumably pulled into the MO mode when the PA is injected. However, there are some indications from other experiments at JILA that residual output facet ASE can cause problems when light present in a several nanometer bandwidth is important.



Figure 2.7: MOPA cylindrical lens setup. The rays indicate the behavior of the beam in the short direction. The spherical collimating lens collimates the beam in the long direction, and introduces a focus into the short direction 10-13 cm from the lens. All distances in this figure are only accurate to within a centimeter.



Figure 2.8: Spatial profile of shaped and collimated MOPA beam. This profile was acquired by measuring the power transmitted through a small pinhole that was translated across the beam.



Figure 2.9: MOPA ASE power dependence on current.

The output power dependence on current from two devices is shown in figure 2.10. All PA's are imperfect amplifiers, and exhibit threshold behavior at low current because the facet coatings have finite reflectivity. The dependence of output power on input power at different currents is shown in figure 2.11. Saturation of the PA gain is observed for input power above 5 mW.



Figure 2.10: Output power vs current for two different MOPA's, each with a different injection power.



Figure 2.11: Saturation of PA gain at different currents.

The sensitivity of the MOPA output power to the polarization of the MO light is shown in figure 2.12. A half-wave plate was used to control the MO polarization relative to the PA input facet ASE polarization (which has the same polarization as the output light). The highest output power is measured for parallel or anti-parallel polarization.



Figure 2.12: Sensitivity of the MOPA output power to the polarization of the MO light. The relative polarization is the angle between the polarization of the MO light and the PA input facet ASE and amplified output light (the polarization of the output does not depend on the input). This data was taken with the PA operating at 1.2 A, and $\sim 5 \text{ mW}$ of injection light.

As a check on the ability of the PA to reproduce the frequency characteristics of the MO, a Rb saturated absorption spectrum was taken with the injected PA light (figure 2.13). The sub-Doppler spectroscopic features show no broadening over the same spectrum taken with the MO. The nano-meter scale frequency characteristics of an injected PA are shown in optical spectrum analyzer (OSA) spectra in figure 2.14. Over several nanometers, the injected PA spectrum is identical to the MO spectrum within the 0.1 nm resolution of the OSA. The total light outside of the MO mode is suppressed by at least a factor of 1000.



Figure 2.13: Rb saturated absorption spectrum of MOPA light.



Figure 2.14: Comparison of MO and injected PA light on an OSA. Although only 1 nm wide spectra are shown, other measurements indicate that the spectra are identical across many nm.

The gain of the PA at fixed current for different MO wavelengths figure 2.15 was measured by replacing the ECDL MO with a Ti:Sapphire laser. Light from the Ti:Sapphire laser was coupled into a fiber in order to avoid changes in the MOPA alignment as the Ti:Sapphire was tuned. The collimated light from the output of the fiber was injected into the PA input facet. The wavelength of the Ti:Sapphire laser was swept using a computer controlled and wavelength calibrated bi-refringent tuner. Changes in the injection power were monitored with a pick-off and the measured gain was normalized to 12 mW of input power.⁷ The gain curve of the PA has a measured FWHM of 13 nm, making the device usable across a wide wavelength range. In particular, 780 nm devices that are meant for use with Rb are still useful for experiments with K at 767 nm.



Figure 2.15: PA gain vs wavelength at 1.2 A.

 $^{^7}$ The light from the Ti:Sapphire laser was not coupled well into the PA, so that the gain appears low.

A table of PA lifetimes from various JILA experiments is given in table 2.1. Two different manifestations of PA death and damage have been observed: rapidly decaying output power preceded by hysteretic current or temperature tuning behavior, and damage indicated by a sudden decrease in the gain (which can then be stable over months or years).

Experiment	Status	Lifetime
Cornell/Wieman/Anderson atom guiding #1	dead	1 year
Cornell/Wieman/Anderson atom guiding $#2$	live	1 year
Cornell BEC	live	4.5 years
Jin K #1	live	3.5 years
Jin K #2	dead	2.5 years
Jin K #3	live	1 year
MOPA built by N. Claussen	live	2 years

Table 2.1: PA lifetime for several experiments at JILA.

2.3.1.3 Trap and Repump MOPA's

The MOPA's used in this experiment follow the standard JILA design. The trap MO is a standard JILA ECDL, while the repump MO has been modified with a 3 stage cooling system (figure 2.16). The trap MO can only use diodes with a center wavelength within 5 nm of 767 nm since the diode cannot be cooled very far below room temperature. Three stages of TEC cooling are capable of chilling the repump MO diode to -40 °C with active convective cooling of the heatsink. This extreme cooling is necessary only when sufficiently short wavelength diodes are unavailable. The diode mount, flexure mount, and baseplate are fastened together with stainless steel screws that are thermally insulated with fiber and nylon washers. All electrical feedthroughs and countersunk holes are sealed with epoxy or epoxy covered duct tape. The boxbaseplate interface is sealed with duct tape, and dessicant (not shown) held in a wire mesh box slows the accumulation of frost. This dessicant only needs to be replaced every 2-3 months. Flexible latex feedthroughs allow the grating to be adjusted and the hex wrenches then mechanically disengaged. Recent work has improved this design by using a commercial hermetically sealed box, and replacing the latex feedthroughs with ultra-torr feedthroughs. More details on the cooling scheme and temperature control can be found in the appendix to this chapter.



Figure 2.16: Repump MO cooling scheme. One inch thick Armaflex insulation covering the laser housing is not shown.

2.3.2 Optical setup

The optical layout is shown in figure 2.17, and the function of each optic is detailed in tables 2.2, 2.3, 2.4, and 2.5. The individual MOT optics and beams do not appear in this figure. The color scheme is: green — push beam; violet — probe beam; brown optical pumping beam; red — MOT trap light; blue — MOT repump light; light blue — frequency reference beams; black — laser lock beams.



Figure 2.17: Optical layout.

Key	Function
s1	science MOT repump light shutter
s2	science MOT trap light shutter
s3	science MOT bright repump beam shutter
s4	collection MOT shutter
s5	push beam shutter
s6	optical pumping shutter
s7	probe beam shutter

Table 2.2: Shutters (Uniblitz LS6T2).

Key	Function
HW1	sets science/collection MOT repump light split
HW2	sets science/collection MOT trap light split
HW3	sets science MOT bright/dim repump beam split
HW4	sets polarization for AO1
HW5	sets polarization for polarization preserving fiber
HW6	sets dim repump beam polarization

Table 2.3: Half-wave plates (Meadowlark Optics commercial retarders).

Key	Part	Function
iso	optical isolator (OFR IO-5C- LP)	MOPA output isolator
AO1	acousto-optic modulator (NEOS N15260)	frequency offsets trap light by 240 MHz
AO2	acousto-optic modulator (NEOS N15260)	change push, probe frequency
IA1	iris aperture (~ 1 mm)	cleans up repump spatial mode (85% transimission)
IA2	iris aperture (~ 1 mm)	creates flat intensity profile for prob- ing
QW1	quarter wave plate (Mead- owlark RQM-100-780)	sets optical pumping polarization
QW2	quarter wave plate (Mead- owlark NQM-100-780)	sets push and probe polarization
pol	polarizer	polarizes optical pumping beam be- fore QW1
\mathbf{SC}	security camera	images collection MOT
PD	photo-diode	measures collection MOT number
RM	gold mirror	optical pumping beam retro mirror
LCD	polarization rotator	sets push polarization
PBS	polarizing beam splitter cube	splits/combines various beams
FC	fiber-coupler (New Focus 9091)	couples probe beam into optical fiber
FH	fiber-holder (New Focus 9095)	couples probe beam out of optical fiber
fiber	polarization preserving single mode fiber (Wave-Optics WF 733)	cleans up probe beam spatial mode and pointing stability
BS	microscope slide cover slip	picks off light for probe beam
DS	3 mm bronze circle glued to pyrex window	science MOT dark spot
SF	25 micron spatial filter	cleans up science MOT trap light spa- tial mode
ND	12% transmission neutral density filter	used for low magnification imaging

Table 2.4: Miscellaneous optical elements.

Key	Function
$\ell 1, \ell 2, \ell 3$	repump MOPA collimating and shaping op- tics
$\ell 6, \ell 7$	trap MOPA collimating and shaping optics
$\ell 4, \ell 5; \ell 8, \ell 9$	telescope to get through isolator aperture
$\ell 10, \ell 18 - \ell 17 - \ell 20$	telescope to couple into AOM — focus occurs inside AO1
$\ell 11, \ell 12$	telescope to expand and spatially filter the sci- ence MOT trap beam — focus at SF
$\ell 19, \ell 23$	telescope to expand beams for collection MOT
$\ell 13, \ell 14$	telescope for science MOT bright repump beam
$\ell 15, \ell 16$	telescope for science MOT dim repump beam
$\ell 21, \ell 22$	telescope to match numerical aperture of fiber
$\ell 24, \ell 25$	telescope for push beam
$\ell 29$	focus push beam at science MOT
$\ell 26$	collimate output of fiber
$\ell 27, \ell 28$	telescope to expand probe beam
<i>ℓ</i> 30	image IA2 onto the atoms for flat intensity profile probe beam
$\ell 31, \ell 32$	high-magnification imaging optics
$\ell 33, \ell 34$	low magnification optics (used with $\ell 31, \ell 32$)

Table 2.5: Lenses.

2.3.3 K DAVLL lock

We use the DAVLL scheme [71] for frequency stabilization of the trap and repump MOPA's. A DAVLL lock for potassium offers the advantage of being able to *easily* frequency lock directly to the 40 K transitions. Further, the laser frequency can be quickly (faster than 1 ms) and reproducibly jumped by many 100's of MHz.

The K DAVLL locking signal is shown in figure 2.18 with the 40 K MOT transitions and other isotope cycling transitions marked for reference. The K absorption cell used in the DAVLL setup is enclosed in an oven and temperature stabilized to 50 °C. A 120 gauss bias field is created by permanent magnets that are located inside of the oven. The lasers are stabilized to the locking error signal using standard JILA electronics via feedback to the PZT driven diffraction grating and to the laser current. At best, the DAVLL locks are stable to 0.75 MHz over a day, with 300 kHz short term stability. The short term stability was measured by beating the all three lasers together, two at a time. The locked linewidth of each laser was extracted from the three beat spectra acquired on a spectrum analyzer. We assume that the laser spectrum is gaussian and that the beat spectrum represents the Fourier convolution of the spectra of the two lasers. We have found that using temperature insensitive polarization optics, particularly calcite polarizers, is important to long term stability. We have successfully manufactured our own calcite polarizers by having the JILA shop polish pieces of calcite that are hand selected from shops in Estes Park, CO.



Figure 2.18: DAVLL locking signal.

2.3.4 Enriched Source

With the usual type of alkali source, the low natural abundance (0.012%) of ⁴⁰K would severely limit the number of atoms that could be collected in the vapor cell MOT. Efficient evaporative cooling relies on trapping enough atoms to have a high collision rate in the gas. Furthermore, reaching T_F at a experimentally achievable temperature ($\sim \mu K$) also requires large initial numbers of atoms. Our experiment would not have been successful without the development of an enriched source that allowed us to trap 10000 times more atoms than previous efforts [47,46].

Potassium enriched in ⁴⁰K is available commercially⁸ in the form of a salt, KCl, rather than as a pure metal. The cost of the enriched material is quite high (roughly \$3000 per 100 mg of K from Trace Sciences), so that the ideal source should deliver potassium vapor for a MOT efficiently using a small amount of material. Large loss due to continuous exposure of the source to the vacuum pumps or due to adsorption of the alkali atoms onto the surface of the cell is unacceptable.

Our enriched source [51] is based on the design of commercially available alkali metal dispensers by SAES Getters [77]. The dispenser contains an alkali salt as well as a reducing agent and delivers small amounts of pure alkali metal through chemistry that occurs inside the vacuum chamber. The metal is released, or evaporated, simply by ohmic heating of the dispenser. The vapor pressure can then be controlled with the current applied to the device; the source is essentially turned off completely if there is no applied current. Unfortunately we were unable to find a company willing to make dispensers for enriched isotopes upon request, presumably because the commercial dispensers are generally mass-produced.

Our 40 K source contains enriched KCl plus calcium [78] for the reduction reaction.

⁸ For the experiments in this thesis, enriched material was purchased from Isotec, Inc. Enriched KCl is no longer available from Isotec, but can be purchased from Trace Sciences, International.

The KCl contains K that is 4.5% ⁴⁰K, 29.1% ⁴¹K, and 66.4% ³⁹K ⁹ (natural abundances are 0.012% ⁴⁰K, 6.73% ⁴¹K, and 93.26% ³⁹K). Upon heating the source, the enriched K is released while the Cl as well as other contaminants are captured by the Ca.

The Ca must be very pure, especially since its dominant contaminant tends to be alkali metals. If proper care is not taken, these alkali contaminants will be released when the source is activated, leading to higher background pressures as well as a reduced relative abundance of the desired isotope. The Ca used in our source was baked at 400 °C under vacuum for 4 days in order to drive out any impurities. As a check on the purity of the Ca, we performed a yield measurement (described below) and verified that a source containing only the clean Ca did not release a noticeable amount of alkali metal.

We prepared a 5:1 molar mixture of Ca and enriched KCl, with both chemicals in a powdered form. Since the reaction depends on adequate fresh calcium surface area, we used a powder of Ca prepared using a jeweler's file and sieved through a woven wire mesh (0.07 mm wire with 0.15 mm apertures).¹⁰ The mixture was put into a small "boat" made from 0.125 mm thick Nichrome (80%-20% nickel-chromium alloy) foil that had been flame annealed, then mechanically cleaned and electro-polished (figure 2.19). Electrical leads, 1 mm nickel wires, were spot-welded to the foil tabs on both sides of the boat. Several sources were made, each containing approximately 2.1 mg of KCl. These were incorporated into a glass arm for use in the collection MOT. It is crucial to avoid heating the dispensers while the pin-press to which they are attached is fastened to the glass arm.

To characterize our source, we have examined the level of contaminants, the total potassium yield, and the relative abundance of 40 K. A major concern for ultrahigh vacuum studies is that the level of contaminants released from the source be sufficiently

⁹ This data is from Isotec.

¹⁰ Calcium is not normally available in a powdered form for obvious reasons. In Boulder's arid climate we were able to prepare calcium filings without using a glove box; the power was stored in a desiccant jar.


Figure 2.19: Schematic of enriched source.

low. We fired one of our sources, as well as a commercial SAES K dispenser, into an uncalibrated residual gas analyzer. With our source, we detected levels of released contaminants, such as Cl, water, CO_2 , and other alkalis, that were no higher than that seen for the commercial dispenser.

Another concern is that the useful lifetime of the source, which is directly related to the total yield, is reasonable. We performed a yield measurement of our source and compared this to a commercial dispenser. Using a triode arrangement, we measured the amount of released potassium by ionizing the gas and counting the collected ion current. The source was mounted along the axis of a helical tungsten filament inside a water-cooled Kovar tube. The filament was heated to roughly 1400 °C and maintained at 30 V relative to the source, while the Kovar was held at -400 V. After baking the apparatus under vacuum and degassing the tungsten filament, the yield measurement was performed by activating the source and monitoring the ion current (using a picoammeter) as a function of time. For calibration, we performed a similar test using a purchased SAES potassium dispenser having a nominal yield of 4.5 mg. Assuming this nominal yield is accurate, we found that our dispenser, containing 2.1 ± 0.3 mg of KCl, released a total of 0.39 mg of alkali metal, giving an efficiency of $19 \pm 3\%$ (figure 2.20). Upon inspection of a spent source we found that the efficiency appears to be limited by the availability of fresh Ca surface area for the required reduction reaction. A single enriched dispenser has been found to last for about three years of daily use in our apparatus.



Figure 2.20: Yield measurement for enriched source.

While the above tests show that our source efficiently delivers potassium, they do not provide a measure of the relative abundance of the isotope of interest 40 K. To make this determination, we fired one of our sources into a small glass cell. Then, using this cell and a 767 nm diode laser, we performed saturated absorption spectroscopy. The results of this test, as well as a saturated absorption spectrum for a cell containing unenriched potassium, is shown in figure 2.21. The enriched source clearly delivers potassium with a higher abundance of both 40 K and 41 K. In particular, the 40 K lines, which are imperceptible for the unenriched source, clearly appear in the spectrum for the enriched source. The frequency axis in figure 2.21 was calibrated using the known hyperfine splittings for K.



Figure 2.21: Saturated absorption spectrum from enriched source. Spectrum (a) was taken using an absorption cell made from an unenriched source, and spectrum (b) from an enriched source.

2.3.5 Collection MOT

Both the Science and Collection MOT's are fairly standard compared to other experiments in the ultra-cold trapped atom gas business. Characterization of the MOT behavior is presented here for completeness and so that the importance of adjustable parameters is clear.

The collection cell (figure 2.22), constructed by Hans Rohner at JILA, is a sixway cross design made from pyrex tubes with optical quality windows. Three glass arms extend from the bottom of the cell and contain enriched K, SAES K, and SAES Rb alkali dispensers. The MOT windows are 1.5 inch in diameter, and there are three small (0.5 inch diameter) ports used for the push and probe beam and for monitoring the MOT position and number. Because of the low vapor pressure of K ($\sim 2 \times 10^{-8}$ torr for K compared to $\sim 2 \times 10^{-7}$ torr for Rb at room temperature [79]), the cell is heated to 40 °C using thermofoil resistive heaters (gray in figure 2.22) in order to prevent buildup of potassium on the glass surface.

The collection MOT uses a three beam, retro-reflected arrangement for the MOT light. The trap and repump light are nearly collinear in all directions. The quadrupole magnetic field (calculated 12 G/cm) for the MOT is provided by a pair of coils (black in figure 2.22). A security camera (**SC** in figure 2.17) can be positioned on one side port or behind the push/probe combining optic in order to align the MOT with the transfer tube. The MOT position can be altered with three orthogonal shim coils (red in figure 2.22) that are used to shift the location of the quadrupole magnetic field zero.



Figure 2.22: Detail of collection MOT.

Typically, the collection MOT runs with 100 mW of trap light and 30 mW of repump light.¹¹ The trap light is detuned by $\Delta_t \sim 25$ MHz from the f' = 9/2 ground state to f' = 11/2 excited state transition, and the repump light is detuned by $\Delta_r \sim 30$ MHz from the f = 7/2 ground state to f = 9/2 excited state transition.¹² The MOT number is calculated using a 6-level model [46] from the photodiode (**PD** in figure 2.17) signal. With an enriched source running at 3.2 A, the MOT fills (at best) up to 2×10^9 ⁴⁰K atoms with an exponential time constant of 0.5 sec. The dependence of the MOT fill on repump power is shown in figure 2.23. The maximum fill rate (atoms/sec) is obtained at the highest values of repump power.

¹¹ The power in each beam is measured directly after the shutter s4. About 90% of this light makes it into in a 1.3 inch diameter beam with a very non-uniform spatial profile.

 $^{^{12}}$ These detunings correspond to beat frequencies of ${\sim}559$ MHz and ${\sim}445$ MHz between the trap and repump lasers and the peak-locked laser.



Figure 2.23: Collection MOT fill dependence on repump power. The number of atoms N in the MOT, exponential loading time constant τ , and loading rate (which is just calculated from the data as N/τ) are shown vs the total repump power. The alkali dispensers were run at 3.0 A for this data, the beat frequencies for the trap and repump lasers were 560 MHz and 443 MHz, and there was 86 mW (measured after **s4**) of trap light.

2.3.6 Science MOT

The science cell (figure 2.24), also manufactured by Hans Rohner, is a modified six-way cross design. The cell is only 1 inch across in one direction so that the magnetic trap coils can be positioned as close together as possible. The MOT windows are 1 inch in diameter, and there are three 0.5 inch diameter ports for probing and miscellaneous optical access. A security camera or a photodiode and lens monitor the MOT flourescence through the top port.



Figure 2.24: Detail of science MOT.

The science MOT must serve multiple purposes: efficiently catching atoms from

the collection MOT push, filling up with large numbers of atoms, and then preparing the sample to load into the magnetic trap. There are strongly different requirements for each role. For example, an efficient catch requires a MOT with strong forces, while filling up with large numbers of atoms requires low collisional losses and therefore low excited state fraction and large trap volume. In order to make the science MOT multi-purpose, the MOT position, repump power, and laser detunings are adjustable and jump between values for different stages of the experimental cycle.

Four coils (red in figure 2.24) are used to position the MOT for the best catch from the collection MOT push and for optimized loading into the magnetic trap. These coils are also used to provide the quantization axes for probing and optical pumping. We use the following conventions for referring to spatial directions: East/West (E/W) for along the probing and push direction, vertical or up/down (U/D), and axial or North/South (N/S) for along the magnetic trap bias field. The calculated field at the center of the science cell from the shim coils is: U/D 2.63 gauss/A, N/S 1.86 gauss/A, and eastern E/W 2.66 gauss/A. Two coils (orange in figure 2.24) provide the quadrupole field (calculated 7 G/cm) for the MOT. The science MOT uses a six beam configuration, with a complicated repump light scheme that is designed to minimize collisional losses [80] in the MOT while allowing for strong viscous forces for the catch. The repump light is collinear with the trap light in all six directions. About 30 mW total of trap light in a 1 inch $1/e^2$ beam diameter is used in the MOT. Data showing the relative science MOT catch efficiency and density dependent loss rate vs. repump power are shown in figure 2.25. The best results are obtained with high repump power for the catch and low repump power for holding the atoms between catches. In order to take advantage of this, a bright repump beam (20 mW) is turned on for the catch, while a dim (5 mW) repump beam is used to hold the atoms. Both repump beams also have a ~ 1 inch $1/e^2$ beam diameter. A dark spot (**DS** in figure 2.17) [81] is put into the dim beam using a solid brass circle (diameter 4 mm) mounted on a 1.5 in diameter pyrex window. The dark spot is aligned along the optical axis to minimize bright Fresnel fringes at the MOT and along the orthogonal directions to minimize the MOT fluorescence. The dark spot allows for low MOT excited state fraction while avoiding loss due to "leaks".



Figure 2.25: Science MOT repump power dependence. Data taken under slightly different conditions is shown in the inset — at very low power, the MOT becomes "leaky". All of this data was taken without the dark spot. The data in the main plots was taken with 39 mW of trap light in the science MOT (measured after the spatial filter), and with beat frequencies of 560 MHz and 443 MHz for the trap and repump lasers.

The science MOT fill is characterized by two parameters: the number of atoms caught per second F and the inelastic loss rate β . Excited state-ground state collisional losses dominate compared to loss due to collisions with room temperature atoms (figure 2.26). The number in the MOT is measured using the CCD camera and fluorescence imaging. An extra set of optics (see table 2.5 and section 2.6) can be combined with the normal imaging optics to decrease the magnification by a factor of 5.8 in order to image the MOT (see the section on absorption imaging for more details). The CCD flourescence signal was calibrated against the signal from a photo-diode that imaged the MOT through the top port. The number is calculated using the CCD signal and a 6-level model [46]. The parameters F and β are measured by fitting the time dependence of the number of atoms in the MOT to a model [82] that assumes no loss except due to inelastic collisions:

$$\frac{dN}{dt} = F - \beta N^2 \tag{2.1}$$

$$N(t) = \sqrt{\frac{F}{B}} \tanh\left(t\sqrt{F\beta}\right) \tag{2.2}$$

and constant trap volume. Under optimized conditions, $F=4 \times 10^7$ atoms/sec and $\beta = 2 \times 10^{-10}$ 1/atoms/sec, so that 4 to 5×10^8 atoms fill into the MOT in 20 sec (figure 2.27). In published work, we have specified a 50% uncertainty in atom number, even when measured with absorption imaging. We specify this uncertainty because absorption and flourescence imaging only agree to this level. The primary source of the uncertainty is the lack of a direct measurement of the light intensity experienced by the atoms for flourescence imaging. See Chapter 5 for more details.



Figure 2.26: Science MOT lifetime. The lifetime is characterized by two exponential time constants (measured with a double-exponential fit): a fast ~ 18 second decay due to density dependent loss and a slower ~ 110 sec decay caused by collisions with room temperature atoms. This data was taken very early in the construction of the experiment, with the trap and repump beat frequencies at 565 and 373 MHz, respectively, and trap and repump powers measured at 55 and 8.5 mW. There was no dark spot in the repump beams for the data in this figure. This data should not be used to compare to the modern performance of the apparatus.



Figure 2.27: Science MOT fill. The fit (solid line) is used to extract the fill and loss rates, which are shown on the plot in experimental units for this particular data. For the data in this figure, the alkali dispensers were set at 3.2 A, the science MOT trap and dim/bright repump powers were 39 mW and 4.5/17.3 mW, and the laser beat frequencies were 559 and 443 MHz for the trap and repump lasers, respectively.

The MOT detunings for the transfer are optimized to load the highest number of atoms into the magnetic trap (figure 2.28); the number of atoms is measured after some evaporation. Before loading, the MOT trap detuning (Δ_t^{cold}) is jumped closer to resonance in order to Doppler cool the gas. Another group has seen evidence for strong sub-Doppler cooling [49], but we have been unable to reproduce these results. The temperature of the MOT vs detuning is shown in figure 2.29; the MOT temperature is cooled from 220 μ K to 150 μ K before loading into the magnetic trap by jumping the trap detuning to $\Delta_t^{cold} \sim 575$ MHz (referred to by the beat frequency) for 20 ms. The repump laser detuning is also jumped at the same time, and, although the effect on temperature is unknown, the detuning (Δ_r^{jump}) is optimized by measuring number after some evaporation.



Figure 2.28: Science MOT detuning dependence. The MOT number is optimized by measuring the number left after some evaporation. For the data in this plot, the total peak optical depth (OD) in the gas is plotted at a 15 ms expansion time after evaporating to 1281.5 MHz in the "evap1" stage (see Chapter 4). The data is taken at fixed temperature, and the peak OD reflects the total number of $m_f = 9/2$ and $m_f = 7/2$ atoms in the gas. Since the number is measured after some evaporation, effects on the evaporation performance due to changes in the MOT temperature are included in this plot. For this figure and the next (figure 2.29) the detunings are measured by the beat frequency with the peak-locked laser at the ³⁹K f = 2 to f' manifold. Note that **AO1** is set to 240 MHz, which adds to the frequency difference (which is the measured beat frequency) between the trap light and the ³⁹K f = 2 to f' manifold. For the data in this figure, the trap and repump dim/bright repump powers were 34 mW and 3.4/13.4 mW. For this figure and the next, the trap light power is measured after the spatial filter, and the repump power after the repump shutters.



Figure 2.29: Science MOT temperature dependence on cold MOT stage detuning. The temperature is inferred by measuring the MOT expansion energy. The r.m.s. size of the MOT is measured using absorption imaging for different expansion times after turning off the MOT quadrupole field and light. The peak optical depth (proportional to collision rate) is shown at a fixed expansion time. For the data in this figure, there was 34 mW of trap light and 1.5 mW of repump light in the science MOT.

2.3.7 Push

The technique we use to transfer atoms between MOT's is also standard compared to other double-MOT experiments. Again, the transfer behavior is characterized here in predominantly a qualitative way so that people running the experiment know which parameters affect the transfer efficiency.

The push beam (green in figure 2.17) is used to transfer atoms through the transfer tube from the collection MOT to the science MOT. Only a single frequency of laser light is used in the push beam. The collection MOT trap and repump light and the collection MOT quadrupole magnetic field is left on during the push. The push beam empties all of the atoms in the collection MOT during a 100 ms pulse every 870 ms (figure 2.30). Note that during the transfer sequence, the collection MOT only fills up to one-half of the maximum fill. The atoms travel down the transfer tube, guided by hexapole magnets, at 20 m/sec (figure 2.31) and are caught with *at best* 35% efficiency by the science MOT. We define the transfer efficiency as the number caught in the science MOT divided by the number that leave the collection MOT for one pulse of push light.



Figure 2.30: Push duration. The number of atoms caught by the science MOT is measured vs. the duration of the push pulse. Most of the atoms that can be caught are pushed in the first 100 ms. The data in this figure and the next were taken very early on in the construction of the apparatus, and may be difficult to compare to current performance.



Figure 2.31: Atom velocity from push beam. The push shutter opens at 0 ms and closes at 100 ms. The number of atoms caught in the science MOT is measured for different delay times after the beginning of the push. The atom arrival time implies a longitudinal velocity spread of 19 to 23 m/s given the 57 cm center-to-center MOT distance.

We do not fully understand the transfer efficiency dependence on detuning and polarization. We hypothesize that the dependence comes from the fact that the MOT quadrupole field is on during the push, and that the magnetic fields that the atoms experience as they enter the transfer tube are complicated. The polarization that produces the best push efficiency for particular push beam detuning is shown in figure 2.32. The best transfer is always produced with linear polarization, changing toward slightly elliptical with changes in detuning (figure 2.32). This complicates the optical setup since the probe beam, which shares polarization optics with the push beam, must be circularly polarized. For this reason an LCD polarization optic (LCD in figure 2.17) is used to modify the push beam polarization before going through a quarter-wave plate (QW2 in figure 2.17). The voltage on this LCD is optimized for the best transfer efficiency. The push beam power is also optimized for best transfer, typically running with 20 mW at ~ -40 MHz detuning (AO2 is set to 234 MHz and the trap laser to a 559 MHz beat frequency) from the f = 9/2 to f' = 11/2 transition. Note that measurements indicate that the velocity of the atoms does not seem to depend on the push detuning over a 120 MHz range.



Figure 2.32: Collection to science MOT transfer dependence on detuning and polarization. The detuning is measured from the f = 9/2 to f' = 11/2 transition. The light is circularly polarized (accordingly to the markings on the quarter-wave plate) when the angle is 325° . The data in this figure was taken with ~ 560 and ~ 440 MHz beat frequencies for the trap and repump lasers, and with 27 mW of trap light and 35 mW of repump light in the science MOT. The trap power is measured after the spatial filter, and the repump power after the repump shutter.

2.3.8 Vacuum System

The vacuum system consists of the science cell, collection cell, transfer tube, and connections to ion pumps. The entire system was baked at 300 °C for three days while connected to a turbo pumped roughing station. During this time the alkali dispensers were flashed to 1.5 A for 5 minutes in order to drive off impurities. After the bake, the system was closed off using an all-metal valve. The collection cell is pumped by a 20 l/s ion pump (**P2** in figure 2.17), and the science cell by a 40 l/s pump(**P1** in figure 2.17). The distance along the transfer tube to the connection to **P2** from the center of the collection cell is ~ 14 cm. The science cell is connected to the four-way cross through a 1 inch diameter glass tube that is ~ 15 cm long.

The limited conductance of the transfer tube isolates the science cell from the vapor in the collection cell. The tube itself is 30 cm long and has an 11 cm inner diameter. A bellows at one end of the transfer tube and glass arms at both ends extend the total length to ~ 45 cm. There is a short arm (9 long) (note that this arm is not shown on any figures in this thesis) that is connected to the transfer tube at a right angle 22 cm from the center of the collection cell. Inside this arm is a small glass block with an embedded magnet that can be moved in and out of the transfer tube. This valve is useful for performing tests where the conductance between the science and collection cells needs to be restricted even further.

The vacuum in the science cell can be characterized by the exponential lifetime for atoms held in the magnetic trap, which is typically 350 sec for our apparatus. Occasionally, for example during power outages, the vacuum may be compromised and the lifetime will drop. In this case, a titanium sublimation pump (**TSP** in figure 2.17) filament is fired to 48 A for 1 minute in order to adsorb contaminants onto a titanium layer on the four-way cross. Normally we have to fire the TSP only once a year.

The pressure in the science cell is not limited by potassium atoms from the col-

lection cell that wander down the transfer tube. This was confirmed with several measurements of the trap lifetime and heating rates for trapped gases. All of the following measurements were done with $\sim 50/50$ mixtures of $m_f = 9/2$ and $m_f = 7/2$ atoms, with initial temperatures between 1 and 2 μ K, and with a total of a few million atoms. The heating rate and trap lifetime were compared for: (i) the normal case (valve open), (ii) with the transfer tube valve closed during the measurement, (iii) with the enriched sources only on for the transfer of atoms between MOT's (43 sec out of a 3 minute cycle), and (iv) with the transfer tube valve closed overnight but open during the measurement (which was done immediately after opening the valve in the morning). Note that the lifetime measurement takes place over 300 to 400 seconds at the end of the experimental cycle. Within the typical 30 to 50 second uncertainty in the trap lifetime and few nK/sec uncertainty in the heating rate there was no difference between any of these four measurements.

2.4 Optical Pumping and Spin Composition

Atoms in multiple internal states are required for rethermalizing collisions to be present during evaporative cooling. Normally in cold atom experiments, the optical pumping is optimized very carefully to load atoms in only one Zeeman level into the magnetic trap. This is done in order to suppress loss caused by inelastic collisions that populate magnetically untrapped states. We found that it is not hard to load multiple Zeeman levels into the trap in a controlled and reproducible way. There are two primary parameters that control the spin composition of the gas that is loaded into the trap: (i) the optical pumping power and (ii) the repump laser detuning used during the cold MOT stage and during optical pumping.

The gas is optically pumped before transfer from the science MOT into the magnetic trap. The optical pumping is optimized to prepare the best spin mixture for evaporation and to improve the fraction of atoms transferred into the magnetic trap. The optical pumping sequence begins by turning off the science MOT trap light and quadrupole magnetic field, leaving the bright repump beam present for hyperfine pumping. During the 2.3 ms before the magnetic trap is turned on, a few gauss field collinear with the magnetic trap bias field is turned on as a quantization field for the optical pumping. The trap laser is set to the optical pumping frequency (Δ_o), and a 100 μ sec pulse of circularly polarized light (brown in figure 2.17) then optically pumps the atoms into predominantly the $m_f = 9/2$ and $m_f = 7/2$ states in the f = 9/2 ground state. The angle, ~ 13 °, that the optical pumping beam makes with the quantization field may be important to determining the spin composition. The light is retro-ed by a gold mirror (**RM** in figure 2.17) in order to avoid large recoil forces on the atoms.

The fraction of atoms transferred into the magnetic trap is measured using MOT recapture and flourescence imaging. The dependence on the quarter-wave plate (**QW1** in figure 2.17) angle used to circularly polarize the light and the laser detuning is shown in figure 2.33. The quarter wave plate is set to produce circular light. We pump on the $f = 9/2 \rightarrow f' = 9/2$ transition as this adds less heat to the gas because the atoms are shelved in the f = 9/2, $m_f = 9/2$ dark state. Measurements of the MOT temperature¹³ indicate that optical pumping adds less than 10 μ K to the temperature of the gas. The total number of atoms after some evaporation is shown in figure 2.34 for optical pumping on each transition.

¹³ Again, the MOT temperature is inferred from a measurement of the MOT expansion energy. The size of the MOT is measured using absorption imaging for different expansion times after turning off the MOT quadrupole field and MOT light. The temperature is compared with and without an optical pumping pulse.



Figure 2.33: Dependence of the fraction of atoms transferred in the magnetic trap on optical pumping polarization and detuning. The detuning is measured from the f = 9/2 to f' = 11/2 optical transition. In the experiment, we pump with $\Delta_o \sim 49$ MHz, which corresponds to a 634 MHz beat frequency between the trap and peak-locked laser. The inset shows the dependence on polarization over a larger range; this data was taken under different conditions when the overall transfer was not optimal. Note that the light is circularly polarized (according to the markings on the waveplate) at 228°. The fact that less than 7% transfer is never observed probably indicates that the light polarization relative to the quantization field is not perfect. The polarization data is taken when pumping on the dark state transition.



Figure 2.34: Evaporation performance for optical pumping on different transitions. The total number is measured after some evaporation. Effects caused by changes in both the initial temperature and number of atoms loaded into the magnetic trap are included into this plot.

Chapter 4 will explain why using a partially spin polarized gas is advantageous for evaporation. The optical pumping parameters are optimized to produce a 65-70% $f = 9/2, m_f = 9/2$ gas, with the other atoms predominantly in the $f = 9/2, m_f = 7/2$ state. The spin composition is measured after evaporating to low enough temperature so that the AG coil is effective at separating the spin components. For this measurement, we check that the spin composition has not changed from the initial magnetic trap load. The spin composition of the gas that is loaded into the magnetic trap is altered by the optical pumping power as well as the repump laser cold MOT stage detuning (Δ_r^{jump}), as shown in figure 2.35 and figure 2.36.



Figure 2.35: Spin composition dependence on optical pumping power. The peak optical depth (OD) is shown at a 10 ms expansion time for a "evap1" final frequency of 1281.85 MHz (see Chapter 4). Because the temperature of the gas is fixed, the peak OD in the gas is proportional to number. The "anti-gravity" coil is used in this measurement to spatially separate the two components. The x-axis in this and the previous figure is measured in the experimental units ("qvolts") used to set the optical pumping power.



Figure 2.36: Spin composition dependence on the cold MOT stage repump detuning. The repump laser detuning is specified by the beat frequency with the peak-locked laser. The data here and in figure 2.35 (above) represents only relative spin composition. The number is measured after evaporation into the regime where the spin composition changes (see Chapter 4). Typically, we choose to run the experiment with $\Delta_r^{jump} = 432$ MHz.

2.5 Magnetic Trap

The magnetic trap is truly the heart and soul of our experiment. We have put tremendous effort into creating a very stable magnetic trap. Working with fermionic atoms is more difficult than bosons for many reasons. For example, the ability of the evaporation to cool in the degenerate regime is very sensitive to the energy resolution with which atoms are removed [54]. The magnetic trap can broaden this resolution as we use field dependent transitions to remove atoms. Also, experiments with a degenerate Fermi gas require tighter constraints on temperature reproducibility. The goal is not just to cool below T_c and produce a condensate. Rather, shot-to-shot reproducibility in T/T_F is critical to making measurements, since all phenomena that depend on the Fermi statistics are sensitive to the degree of degeneracy. The magnetic trap bias field sets the temperature for a fixed evaporative cut. For this reason, we require an ultrastable bias field. Ioffe-Pritchard magnetic traps do not make this easy, as the relatively small bias field is produced by the subtraction of two large fields. On the other hand, a Ioffe-Pritchard trap can attain the large trap depth that is necessary to capture, hold, and compress potassium atoms that are relatively hot. After a serious investment in stabilizing the bias field we now have one of the most stable traps that we know of.¹⁴

2.5.1 Cloverleaf Trap

We use a cloverleaf [19] design Ioffee-Pritchard [83,50] magnetic trap. The cloverleaf coils produce a harmonic trap with cylindrical symmetry, so that there are two orthogonal trap directions: radial $(r = \sqrt{x^2 + y^2})$ and axial (z). The layout of the coils is shown in figure 2.37. The center coils are a Helmholtz pair designed to produce a large field with large curvature:

$$B_z(r,z) \approx B_c + \gamma \left(z^2 - \frac{r^2}{2}\right) \tag{2.3}$$

¹⁴ One outstanding exception is Ted Hänsch's QUIC trap, which is enclosed in "mu" metal shielding.

where B_c is the bias field produced by these coils, and 2γ is the axial curvature (measured in G/cm²). The bias coils are a Helmholtz pair wound in the opposite direction that are designed to cancel some the bias field from the center coils while adding very little curvature. For our system, the calculated bias coil curvature is $\sim 1/20^{th}$ of the center coil curvature. The trap bias field B_0 is then the difference between the fields produced by these two coils.



Figure 2.37: Cloverleaf arrangement. The arrows indicate the direction of current flow. The bias coils are black, the center coils blue, and the petal coils red.

The "petal" coils add radial gradient to the field while adding very little bias field:

$$B_r(r,z) = \beta r \tag{2.4}$$

where β is the petal gradient (measured in G/cm). The magnitude of the field then provides the trapping potential:

$$B(r,z) = \sqrt{\left[B_0 + \gamma \left(z^2 - \frac{r^2}{2}\right)\right]^2 + \beta^2 r^2}$$
(2.5)

which can be expanded to first order assuming that the bias field is the dominant

contribution to the total field:

$$B(r,z) \approx B_0 + \gamma z^2 + \left(\frac{\beta^2}{2B_0} - \frac{\gamma}{2}\right) r^2$$
(2.6)

The harmonic trap frequencies can then be determined from the magnetic field potential energy $U = -\vec{\mu} \cdot \vec{B}$ assuming that the Larmor frequency is sufficiently high so that the magnetic moment μ follows the direction of the field:

$$f_z = \frac{1}{2\pi} \sqrt{\frac{\mu}{m} 2\gamma} \tag{2.7}$$

$$f_r = \frac{1}{2\pi} \sqrt{\frac{\mu}{m} \left(\frac{\beta^2}{B_0} - \gamma\right)} \tag{2.8}$$

We normally operate with fixed center and petal current $I_c=190$ A, and vary the bias current I_b to adjust the radial frequency f_r and bias field. Note that the current through the petal and center coils is always equal.

We have measured the curvature, gradient, and bias field constants for our trap by measuring the trap frequencies and bias field for fixed center coil current and different bias coil currents. The trap frequencies are measured by exciting sloshing motion of the cloud. A shim coil is turned on for a few oscillation periods, so that the trap center is displaced. This shim is then quickly turned off, and the ensuing center-of-mass (COM) motion of the gas (figure 2.38) is measured in time. By fitting the COM coordinate to a sine wave, the trap frequency is determined. The bias field is measured by extrapolating to zero temperature in evaporation trajectory data (see Chapter 4). For a harmonic trap, the cloud temperature scales linearly with the final microwave frequency in evaporation. ¹⁵ Careful to avoid regimes where the spin composition is changing, the beginning of the cooling, and regimes where evaporation works only marginally, the intercept of a linear fit of temperature vs. microwave frequency is used to determine B_0 (figure 2.39).

 $^{^{15}}$ Note that this is *not* true for N, particularly in the degenerate regime.



Figure 2.38: Determination of harmonic oscillator frequency. The center of the cloud is plotted after exciting vertical slosh of a spin polarized $m_f = 7/2$ gas in the "tight" trap. A fit (solid line) is used to determine that the radial frequency is 117 Hz.


Figure 2.39: Determination of B_0 . Two-frequency evaporation is used to cool the gas, and the temperature at different final microwave frequencies is plotted (see chapter 4 for more details). A linear fit (solid line) is used to determine the frequency where the temperature would go to zero. This frequency corresponds to 5.452(1) gauss for the $f = 9/2, m_f = 9/2 \rightarrow f = 7/2, m_f = 7/2$ transition (this data taken in the tight trap).

We find that $\gamma = 0.295 \text{ G/cm}^2/\text{A} \times I_c$, $\beta = 0.885 \text{ G/cm}/\text{A} \times I_b$, and $B_0 = 1.27(7)I_c - 1.38(8)I_b$. The curvature γ is determined by the axial frequency at fixed I_c , while β is determined from B_0 and f_r at different bias coil currents. The dependence of B_0 on the center and bias coil currents is determined from the measurements of B_0 at different values of I_b .

All of the work done in this thesis used three different traps: the "evap", "tight", and "load" traps. Parameters for these traps are detailed in table 2.6. Some tests were also done with a low bias field trap — the "single frequency" trap. Variation from these parameters will be noted in this thesis where appropriate. Note that the frequencies for $m_f = 7/2$ atoms are reduced from the $m_f = 9/2$ frequencies by the square-root of the ratio of the magnetic moments ($\sqrt{7/9}$).

Trap	I_b	B_0	$m_f = 9/2 f_r$	$m_f = 9/2 f_z$
evap	174.9 A	1.3 G	234 Hz	19.5 Hz^*
tight	172.0 A	$5.34~\mathrm{G}$	$135 \mathrm{~Hz}$	$19.5~\mathrm{Hz}$
load	142 A	45 G	$44 \mathrm{~Hz}$	$19.5 \ \mathrm{Hz^*}$
single frequency	175.5 A	0.68 G	380 Hz^{\dagger}	$19.5 \mathrm{~Hz^{\dagger}}$

Table 2.6: Different magnetic trap parameters used in this thesis. We have the best measurements of the axial frequency in the "tight" trap. Because f_z should not depend on the bias current, we list the measured axial frequency in the "tight" trap for all of the traps. For the traps marked with an "*", we have measurements of the axial frequency with larger uncertainty. We have no measurements of the axial or radial frequency for the trap marked with a "†". The radial trap frequency for this trap is inferred from a measurement of the bias field.

2.5.2 Trap Depth

For the $m_f = 9/2$ state, the trap depth is set by the size of the science cell in the z (axial) direction. The trap depth can be calculated using the form for the magnetic field without the harmonic approximation (equation 2.5). For the $m_f =$ 9/2 state, the trap depth is then $-\frac{\mu_{9/2}}{k_b}[B(y_0, 1.27 \text{ cm}) - B(y_0, 0)] = 6 \text{ mK}$, where y_0 is the equilibrium position in the vertical direction (taking into account gravitational sag). For other states, the magnetic moment changes at high field and gravitational sag in the vertical direction is important. The Breit-Rabi formula (see chapter 4) can be used to compute the field-dependent magnetic moment and the total energy $U(y, z) = -\mu(y, z)B(y, z) + mgy$. However, it turns out for the trap parameters that we use that the trap depth is still limited by the size of the science cell for the $m_f =$ 7/2 component. In the tight trap, the trap depth for the $m_f = 7/2$ component is $-\frac{\mu_{7/2}(y_0, 1.27 \text{ cm})}{k_b}[B(y_0, 1.27 \text{ cm}) - B(y_0, 0)] = 3 \text{ mK}.$

2.5.3 Mechanical Design

The coils were wound using 1/8 inch copper refrigerator tubing and insulated with heatshrink tubing. Roughly 8 liters per minute of water flow through six parallel paths in order to remove heat created through Ohmic dissipation; there is a ~ 1 °C increase in the water temperature at the combined output of the water paths during operation. The power supply is inter-locked to the water flow via a paddle-wheel flow meter (Proteus Industries). The water and current connections to the coils are carefully strain-relieved to the optical table. A manifold constructed from stainless steel optical posts (gray in figure 2.24) fixes the coils relative to the table and to each other. In addition, Delrin plastic brackets and phenolic plastic rods help to fix the distance between the halves of the coils. More details on mechanical stabilization can be found in the sub-section (2.5.5) discussing the stability of the bias field.

2.5.4 Current Servo and Switching Electronics

The primary requirements for the magnetic trap current servo are low noise, low drift, insensitivity to the microwave field used for evaporation, and fast (compared to the harmonic trap frequencies) switching times. A low noise servo (< 6 ppm integrated current noise) is constructed from MOSFET's via feedback from closed-loop current sensors (Hall probes — F.W. Bell CLN-300). The Hall probes are temperature stabilized to minimize drift, and the control voltages used for the servo are generated from a special low drift (< 1 ppm over 6 months) circuit. Sensitivity to the ~ 1 GHz microwave field used for evaporation is reduced by strictly following a star-grounding protocol and using high-frequency filters on inter-case connections. Additional MOSFET's are used to achieve fast switching times.

The MOSFET setup is shown in figure 2.40. The MOSFET's are two Advanced Power Technology APT10M07JVR devices wired in parallel; each FET is capable of passing 225 amps and dissipating 700 W. The FET's are mounted to a water cooled plate with thermal compound in order to facilitate good thermal contact.¹⁶ **F1** and **F3** are used for switching, while **F2** is used to serve the current through the bias coils by shunting some of the total current. Typically, only a small fraction of the current is shunted so that there is some common-mode rejection of noise between the bias and center coils. **F4** is used to serve the total current through the circuit. The Hall probes are used to sense the current for the serve circuits, and there is an extra set (not shown) that can be used to measure noise and monitor the current. The power supply (HP 6682A) is a particularly robust and low noise 21 V, 240 A supply that is run in constant voltage mode. The HP6682A power supply voltage is set so that the **F4** drain-source voltage is ~ 1 V.

¹⁶ Note that excellent thermal contact between the FET's and the water cooled plate is essential to avoiding fires. The water cooled plate should be carefully cleaned with a degreasing agent and methanol before applying the thermal grease and mounting the FET's



Figure 2.40: Magnetic trap current control scheme. The large star represents the physical star ground, which is a copper bar connected to the optical table to an earth ground.

The petal and center coils and the bias coils are shown in figure 2.40 as self-resonant circuits.¹⁷ The measured average resistances of the coils are 5.2 m Ω per petal coil, 7.4 m Ω per center coil, and 9.6 m Ω per bias coil. The self-resonance frequency, 450 Hz, of the coil-FET system was measured by DC biasing the FET gates and driving the gates with an oscillating voltage across a 20 Ohm resistor. The current response was measured at different frequencies using an FFT analyzer. The low resonant frequency enforces harsh constraints on the servo and switching electronics. Large inductance, and therefore low self-resonant frequency, is the main disadvantage of the cloverleaf scheme.

The servo circuit ("main" servo) that controls $\mathbf{F4}$ is shown in figure 2.41. All connections into the circuit come through "pi" filters (figure 2.42) that reduce 1 GHz signals by ~ 40 dB (measured). The signal from the Hall probe enters at $\mathbf{J4}$ and $\mathbf{J5}$ and is differentially amplified by $\mathbf{U1}$ in order to reject pick-up. $\mathbf{C1}$ is necessary to provide local feedback to suppress oscillation at $\mathbf{U1}$. The Hall probe signal is compared to the control voltage (connected at $\mathbf{J6}$) by $\mathbf{U2}$, which is the servo loop-filter. The servo proportional-integral (PI) loop-filter corner frequency is set by the combination of $\mathbf{R1}$ and $\mathbf{C2}$, while $\mathbf{R1}$ sets the open loop gain past the PI corner. The PI corner frequency and the proportional gain are optimized to maximize servo bandwidth while avoiding oscillations. $\mathbf{D1}$ prevents the output of $\mathbf{U2}$ from going lower than -0.6 V in order to avoid integrator wind-up to -15 V.¹⁸ Connecting to the gate through $\mathbf{R2}$ is necessary to spoil the Q of a resonant circuit (with resonant frequency at typically 10's of kHz) that is formed by the FET Miller capacitance and the coils.

 $^{^{17}}$ The coils have some capacitance from the gap between turns, and therefore behave like an L-C-R resonant circuit.

 $^{^{18}}$ If the output of **U1** is allowed to go to the negative rail, at least an integrator time constant is required for the servo to respond to current turn-on.







Figure 2.42: High-frequency ("Pi") filter. It is crucial that the capacitor leads are as short as possible and connected directly to the case.

The bias servo is shown in figure 2.43. The circuit is similar to the main servo, with different values of **R1** and **C2** and some op-amp and FET protection circuitry. The protection circuitry is necessary for the bias servo since the bias coils can produce large voltages at the source (and gate via the Miller capacitance) of **F2** during switching. **D2** prevents damage to the output stage of **U2** caused by high positive voltage coupled through the Miller capacitance of **F2**. The network connected to **D3** and **D4** prevents the gate-source voltage of **F2** from exceeding 9.7 V, while the **D5-D6** network prevents the gate of **F2** from falling below 0 V. **R3** is necessary so that the protection circuitry does not overload the current capabilities of the output stage of **U2**, and **C3** bypasses **R3** to enable high-frequency response.



The control voltages used for the main and bias servos are generated by the circuit shown in figure 2.44. The primary voltage source is a LM399H voltage reference. The high frequency noise caused by the LM399H heater (which is necessary to eliminate long-term drift) are reduced by the filtering capacitors C1-C4. Sensitivity to the Zener bias voltage due to slight Ohmic response is reduced by the 7912 voltage regulator. The main servo control voltage is derived from a 4-resistor voltage divider (**R1-R4**). This scheme allows precision adjustment of the control voltage by adjusting the small resistors (R2, R4) while avoiding the drift, noise, irreproducibility and temperature sensitivity that we have observed with trim-pots. Control voltages for the bias serve for different traps are fed through a multiplexer. Low noise ramps between these control voltages are executed by an Stanford Research Systems DS345 arbitrary waveform generator (computer controlled) that is also connected to the multiplexer. The computer controlled TTL signals that connect to the multiplexer (A0-A3) do not have their signal ground connected to the circuit or case ground. This is essential to avoiding ground loops since the computer hardware is located on a different circuit and has a different ground. The control voltages supplied by this circuit have < 1ppm noise in a 10 kHz bandwidth, and drift by less than 1 ppm per six months.



Figure 2.44: Control voltage circuit. Some of the bias servo control voltage dividers are not shown. Dashed boxes are used to group voltage dividers that produce individual control voltages.

The Hall probe current sensors are mounted in a temperature controlled, grounded box. The temperature is controlled with 1 mK stability using standard JILA electronics and TEC's. The Hall probes behave like current transformers with a turns ratio of 2000; the current is then read out across as a voltage across a sense resistor. We use Vishay 50 Ohm 4-terminal resistors for the sense resistors. The sense resistors are temperature controlled in the same housing as the Hall probes.

The measured noise spectrum (using an HP FFT spectrum analyzer) of both the bias and main current when the servo's are locked is flat (within 3 dB) out to 3 kHz, and then falls off with a 10 dB/octave slope. The measured integrated noise (rms) in a 5 kHz bandwidth is less than 6 ppm for both currents, although we do not know how much of this noise is correlated between I_b and I_c .

The circuit used for switching the trap on and off is shown in figure 2.45. For the trap turn-on, the gates of F1 and F3 are switched to +30 V simultaneously (within a few nano-seconds). Both servo's are locked within 4 ms, with some overshoot so that the trap is actually fully on within 1 ms of the F1 and F3 turn-on. For turn-off, F1 is switched off after a controlled delay compared to F1 (the gates are switched to 0 V). This is basically a control on the boundary conditions for the turn-off of the bias coils. The bias coils will reverse bias F2 and turn on the drain-source protection diode if the main current is switched off while current is flowing through the bias coils. This can drastically slow down the turn-off, or even reverse the bias field as current can flow in the bias coils once the center coils have shut off. The current during the turn-off is shown in figure 2.46, along with the calculated trap frequencies. The turn off is fast compared to the harmonic trap frequencies, with a small delay (0.2 ms) between the radial and axial directions.



power is bypassed w/ 0.1 F caps on all chips

Figure 2.45: Magnetic trap switching circuit.



Figure 2.46: Magnetic trap turn off. The harmonic trap frequencies are calculated using the trap parameters and the measured main and bias current.

The grounding scheme is shown in figure 2.47. The star ground is a copper bar ($6 \times 2 \times 1$ inches) that is attached to an earth ground through the optical table. Each circuit is housed in a separate box — this simplifies diagnosing and eliminating sensitivity to the microwave field. Each case has a separate connection to the star ground. The case grounds are never attached to the circuit board grounds directly. **P1-P6** are power supply lines that are carried in shielded cable. **S1-S3** are signal lines that are transmitted on twisted shielded pair cable. **S4** and **S5** are the Hall probe temperature servo TEC and thermistor connections. The cable shields are connected only at one end, and we choose the convention that the "sending" side is connected. The microwave field used for evaporation can cause DC servo errors via rectification in op-amp input stages. Without this grounding scheme, significant sensitivity (several parts in 10^5 in the current) was observed; after implementing the star ground scheme there is less than 1 ppm sensitivity when 5 W of microwave power is applied at 1 GHz.



Figure 2.47: Magnetic trap electronics grounding scheme. Thick lines indicate ground connections made with at least 12 gauge wire. Shaded rectangular bars indicate power connections made with shielded three-conductor cable. Hollow rectangular bars indicate signal lines for the Hall probe and sense resistor temperature servo. "Squiggly" lines represent signal lines in twisted shielded pair cable. The small solid circles indicate the ground connection for cable shields. Arrows represent water flow for cooling purposes. The large star is the physical star ground.

2.5.5 Bias Field Stability

Stability of the magnetic trap bias field is important on two timescales: the evaporation time scale (seconds) and the experimental cycle timescale (minutes). Jitter in the bias field on the evaporation time scale limits the energy resolution of the evaporative knife. Excellent energy resolution has proven to be especially important for evaporative cooling of fermions [54]. Shot-to-shot reproducibility in number and temperature is affected by drift on the experimental cycle timescale.

2.5.5.1 Jitter

It is most straightforward to use the evaporation behavior itself to measure the bias field stability. The bias field noise on the evaporation timescale is determined by measuring atom loss caused by pulses of microwave radiation at different frequencies. We introduce the variable ν_0 as the resonant microwave transition frequency at the minimum of the magnetic trap (at B_0) (see Chapter 4 for more details). Atoms experience an increasing resonant microwave frequency with decreasing magnetic field, and experience an absolute minimum in magnetic field at B_0 . If the applied microwave frequency ν is greater than ν_0 , then no atoms are resonant and no atom loss from the magnetic trap is observed. For ν slightly less than ν_0 , all of the atoms can be removed. Therefore, a step-like response in atom loss around ν_0 should be observed, and the width of this "step" is broadened by noise on the bias field (see Chapter 4). The width of the "step" can be used to determine the noise on B_0 . This method was tested by adding a known amplitude of white noise to the bias coil current and measuring the effect on the width of the "step".

We assume gaussian noise on the bias field. Because there is a one to one mapping of field to microwave frequency, this noise is equivalent to a gaussian microwave frequency spectrum with width Σ (see chapter 4 for more details). The effective power for removing atoms is then the integral of this gaussian spectrum, or an error function. By mapping out the loss vs. frequency at fixed pulse length¹⁹ the width of the original gaussian is determined. The measurement is done with pulses with just enough power to remove nearly all (~ 95%) of the atoms for $\nu < \nu_0 - 3\Sigma$. Measurements of this resonance before and after optimization of the servo response is shown in figure 2.48. The gas is cooled to 1.3 μ K, and then pulses are applied on the $f = 9/2, m_f = 9/2$ to $f = 7/2, m_f = 7/2$ transition in a 1.3 gauss bias field trap. The total ($m_f = 9/2$ and $m_f = 7/2$) number of atoms is measured after the pulse. Fits to an error function give r.m.s. widths of 27(9) kHz (before) and 4(3) kHz (after), which correspond to 19 mG (before) and 2 mG (after) noise in the bias field on the evaporation timescale. The 2 mG figure for the optimized trap that we use now is really an upper limit as the measurement cannot truly resolve the linewidth within the experimental scatter in number.

¹⁹ The pulse length is typically set between 200 and 1000 ms. Experimentally, we measure that a change in the pulse time between these limits has no effect on the measured noise on B_0 .



Figure 2.48: Magnetic trap bias field stability on the evaporation timescale. The top plot is the response before optimizing the magnetic trap current servos, and the bottom plot is after optimization. See the text for details.

2.5.5.2 Shot-to-Shot Stability

The shot-to-shot stability of the bias field can be measured by measuring the temperature of the gas at a fixed evaporative cut frequency over many experimental cycles. There are many factors that drive drift of the magnetic trap. The current itself is measured to be stable to better than a part per thousand over two years, and instability in the current is not a significant source of shot-to-shot noise. The trap is primarily affected by changes in temperature and stray magnetic fields. Changes in temperature cause changes in the location and distance between the coils, while stray fields add to the bias field directly. Currently, the trap bias field shifts by less than 1 mG (the equivalent to 8 nK for two-frequency evaporation — see Chapter 4) over 20 experimental cycles after the initial warm-up.

Changes in temperature are the primary source of initial drift in the bias field. The Ohmic dissipation of the science MOT coils and the cloverleaf coils as well as the lab air conditioner drive temperature drift. The science MOT coils (in thermal contact with the cloverleaf coils) are off for most of the experimental cycle, and must be kept off between cycles. With the science MOT coils off between cycles, the trap takes about three experimental cycles (~ 9 minutes) to warm up. The bias field changes by 14 mG (~ 100 nK for the two-frequency evaporation) during this time, with the temperature of the coils changing by ~ 1 °C. If the science MOT coils are allowed to heat the magnetic trap coils between cycles, the warm-up time can increase to greater than half an hour.

The shift in B_0 due to changes in the temperature of the cloverleaf coils has been suppressed in general in two ways: by covering the coils with epoxy and by fixing the two halves of the cloverleaf coils together with phenolic plastic rods (brown in figure 2.24). Phenolic plastic is a fibrous material that has a linear thermal expansion coefficient of only $15 \times 10^{-6}/K$. Four 1 inch long and 0.5 in diameter rods were glued between the coils using thermally conductive epoxy (Tra-Con BA2151). The thermal expansion coefficient of the epoxy matches aluminum, which is not much different than copper tubing. The phenolic rods also serve to damp vibrational modes where the coils move out of phase.

The lab air conditioner also has thermal and electrical effects on the trap. The air conditioner condenser cycle is roughly 15 minutes, and each time the condenser turns on the bias field shifts by less than 3 mG.²⁰ Some of this shift is due to electrical transients that affect the Hall probe temperature servo. There is also contribution from the optical table (which oscillates with 0.25 °C amplitude on the air conditioner cycle) and air temperature changing. The effect of electrical transients has been suppressed by adding high-frequency filters to the temperature servo.

The bias field is also affected by the magnetization of low-grade stainless steel optical posts and screws that are used to mechanically stabilize the coils. The magnetic field from the stainless steel components adds directly to the bias field. Confirmed with a flux gate magnetometer located near the trap and aligned with the bias field, the steel components are magnetized by the shim and quantization coils. Shifts as large as 27 mG have been observed when using a quantization field for probing that was collinear with the bias field.²¹ These effects have been minimized by using quantization fields that point in directions orthogonal to the bias field (E/W and U/D). Further, both quantization fields are turned on for every experimental cycle, with a delay to select the field used for probing. The effect of switching between quantization fields has now been minimized to less than a 1 mG shift in the bias field.

Effects on B_0 caused by stray fields from the transfer tube magnets are avoided by replacing the last section (closest to the science cell) with electromagnets (green in figure 2.24). The last section of permanent magnets, 6 cm long, caused a 140 mG drift per hour in the bias field. This drift seems to be very long term warm up caused by

²⁰ This was 7 mG before installing the phenolic rods and fixing the Hall probe temperature servo.

 $^{^{21}}$ We observed changes this large within a few experimental cycles of switching between using a U/D vs. E/W quantization field for imaging.

heating of the magnets. Without these permanent magnets, the fill rate F for the science MOT drops by a factor of 2.5. Switching to the electromagnets (three coils arranged in a hexapole configuration, each 100 turns carrying 3 amps) actually increased F by a factor of 1.5 compared to transfer with the permanent magnets.

2.6 Absorption Imaging

We use standard destructive, absorption imaging techniques. Images are produced by illuminating the gas, after release from the magnetic trap and some expansion, with a 24 μ sec pulse of resonant light ("probe beam" — green in figure 2.24). The shadow cast by the gas into the circularly polarized probe beam is imaged onto a CCD camera by lenses ℓ 31 and ℓ 32 using the "high magnification" optics, or by the combination of ℓ 31, ℓ 32, ℓ 33 and ℓ 34 for the "low magnification" optics. The exact placement of the imaging optics is shown in figure 2.49. The optical depth of the gas is calculated from three images: one with the atoms present (shadow—S), one without the gas (light—L), and one without the probe beam (dark—D). The optical depth is then:

$$OD(y,z) = \ln\left[\frac{L(y,z) - D(y,z)}{S(y,z) - D(y,z)}\right]$$
(2.9)

which reproduces a two-dimensional image of the gas column density. This image is then fit and information on number, temperature, and energy can be extracted (see Chapter 5).

The magnification of the imaging optics is determined by watching the gas drop after release from the trap and measuring the gravitational acceleration constant (figure 2.50). We have only carefully characterized the "high magnification" optics, and the "low magnification" optics are calibrated via a direct comparison between images of identical gases. The "high magnification" imaging optics have a magnification of $6.35 \ \mu m/pixel$. The camera is focused (see figure 2.51) on a cloud with a diameter close to the diffraction limit of the optics. We measure a diffraction limit of < 10 μm FWHM (see the caption to figure 2.51), which is consistent with the calculated limit from the first lens in the optical setup of 5 μm .

HIGH MAGNIFICATION



Figure 2.49: Schematic of imaging optics. In the experiment, there is a 1.5 inch square mirror between the 12.5 and 50 cm lenses. There is also a ~ 1 inch by ~ 2 inch rectangular gold mirror between the atoms and the 12.5 cm focal length lens.



Figure 2.50: Magnification measurement. The vertical position of a $m_f = 9/2$ gas in free fall is plotted for the two quantization fields used for imaging. A fit to the E/W data indicates that the magnification is 6.39(3) μ m/pixel. The U/D data shows a small difference (6.00(3) μ m/pixel) that is consistent with the force on the atoms caused by the small gradient from this coil. These fits determine that the initial center of mass velocity, which is left as a free parameter, is zero.



Figure 2.51: Camera focus. The r.m.s. radial size σ of a gas cooled to 270 nK and expanded for 1 ms is measured for different camera positions. The fit (solid line) to the expected dependence $\sigma = \sigma_0 \sqrt{1 + (\frac{x-x_c}{R})^2}$ determines the depth of focus R, the focussed position x_c , and the actual size of the object σ_0 . From this data we find R = 2.3(2) cm, and $\sigma_0 = 1.70(8)$ pixels, where the expected size is 1.6 pixels from the temperature and harmonic oscillator frequency. This implies a ~0.7 pixel (gaussian) diffraction limit, as the observed size is $\sqrt{\sigma_0^2 + d^2}$ where d is the diffraction limited size. We treat this as an upper limit on the diffraction limit of the imaging optics.

Typically, the maximum optical depth we can measure is 2.5, which we hypothesize is limited by light that the atoms cannot absorb and by lensing. Experimental checks reveal that a gas with an OD higher than 1/2 of this maximum observable OD will have a shape distorted enough to produce significant errors in the measured energy and temperature (see Chapter 5). In order to observe larger optical depths, we switch from a quantization field collinear with the probe beam (E/W) to one at a right angle (U/D).²² This changes the polarization relative to the atoms to 1/4 left circular, 1/4right circular, and 1/2 linear. The observed optical depth is reduced by a factor 2 (figure 2.52), with no other measured effects on the image. We have observed that the factor by which the OD is reduced is a good indicator of the quality of the polarization of the light. Also, with the U/D quantization field, actual OD's as high as 7 have been observed.

 $^{^{22}}$ Normally, one would detune the probe beam from the atomic resonance in order to reduce the observed OD. However, lensing caused by detuning the probe beam introduces significant distortions in the cloud sizes and shape. See Chapter 5 for more details.



Figure 2.52: Optical transition lineshapes. The peak OD in the gas is measured for the two imaging quantization fields as the probe AOM (AO2 in figure 2.17) is varied. The atomic resonance can be precisely located this way, and the frequency shift here is from different magnitudes of the E/W and U/D fields. The OD correction factor is measured by toggling between the quantization fields with the AO on resonance. The beat frequency between the trap laser and the peak-locked laser was 575 MHz for this data. The widths from Lorentzian fits (solid lines) are 6.9(1) MHz and 6.0(3) MHz for the E/W and U/D quantization fields, respectively. This data was taken at a 20 ms expansion time with a $m_f = 9/2$ gas at 100 nK.

The thermometry techniques used for the work in this thesis rely on detecting small changes in the shape of the expanded gas. Even at our lowest temperatures, the shape only deviates slightly from a gaussian. For this reason, it is crucial to have a very flat background plane in the absorption images. To achieve this we prepare the probe beam using a non-standard method. The probe beam is spatially filtered by a polarization-preserving fiber. A small iris aperture (**IA2** in figure 2.17) is then put into the very center of the collimated beam, so that the transmitted beam has a flat intensity profile. The aperture is then imaged onto the gas by a 100 cm focal length lens (ℓ 30 in figure 2.17). It is crucial that no part of the probe beam is truncated by any optics, as this can produce stripes into the image.

2.7 Microwave Coil and Sources

Our evaporation scheme utilizes microwave transitions between hyperfine ground states to transfer atoms into magnetically untrapped states. We use microwave transitions so that we have selectivity in removing different spin components. The transitions we use primarily address the electron and not nuclear spin so that the coupling to magnetic fields is strong (see the appendix to Chapter 4). Using microwave fields for evaporation instead of radio-frequency (rf) is not easy — working at higher frequencies brings many complications. Achieving excellent performance for the evaporation requires producing a large oscillating magnetic field between 1200 and 1300 MHz at the atoms, producing a controlled frequency sweep with high frequency precision (1 kHz at the end of evaporation), and producing two frequencies simultaneously. Meeting all of these requirements is actually fairly difficult. Working at 1 GHz is awkward electrically as lumped components do not work well, lead lengths are important, the current flows in a few micron layer on the surface of conductors, and traditional microwave components are large and bulky. Impedence matching at 1 GHz is also difficult, as the inductive component of the reactance becomes very large and self-resonances tend to appear below 1 GHz. Further, working between the near and far field regions is not intuitive, and the field dependences can be unexpected. Affordable synthesizers with the required resolution and stability at these frequencies are not designed for rapid frequency sweeps, and often perform this task disastrously poorly.

In order to produce a large magnetic field a coil shaped device must be used. Antenna designs were tested and do not work well for this purpose. The coil selfresonance must be above the frequencies of interest. A modified version of the transducer from a 1950's style grid-dip meter satisfies these constraints well. Our coil design is shown in figure 2.53 (purple in figure 2.24). The self-resonance of this coil is above 2 GHz, and the large surface area permits low surface resistance so that power coupled into the coil produces large currents and therefore large magnetic field. The surface is chemically etched and tinned to improve the surface quality and to prevent oxidation; this procedure lowered the total impedence at 1250 MHz by a factor of 65. The measured (using a network analyzer) impendence of the coil at 1250 MHz is (2.5-4 i) Ω - it has only 2.5 Ω of surface resistance and the overall reactance is capacitive. We use 0.086 inch diameter semi-rigid coaxial cable to connect the coil to an amplifier for its mechanical stability and ability to transmit 1.3 GHz signals with low loss. Connecting the coil to the coaxial cable must be done carefully. A short section (5 mm) of the coaxial cable center conductor is exposed. Large diameter copper wire is soldered from the center conductor and shield to either end of the coil. Using as short as possible leads and making excellent solder joints is critical to avoiding loss of the microwave power and creating a low-inductance connection.



Figure 2.53: Microwave coil.

The coil is impedence matched using a stub-tuning technique [84]. This permits high power to be coupled into the coil, with good transmission over a relatively large bandwidth. The basic idea of stub tuning is shown in figure 2.54. A load with admittance (1/impedence) Y_L is attached to a length d of coaxial cable so that the admittance looking into the coax-load system in $1/50 \ \Omega + iB$. An open stub of length L with admittance -iB is wired in parallel to the load so that the total admittance looking into the system is $1/50 \ \Omega$, resulting in a matched condition. Our implementation is shown in figure 2.53 with d=5.7 cm and l=3.5 cm. The stub is made with a standard SMA "T" and RG-174 coaxial cable. The transmitted power into the coil is shown in figure 2.55. Greater than 60% transmission is achieved over the frequency range of interest.



Figure 2.54: Stub tuning scheme. The lumped admittances of the stub and the load-coax system are grouped inside of dashed boxes.



Figure 2.55: Transmittance of microwave coil. The power reflected from the stub-tuned load was measured using a network analyzer.

The position dependence of the field was measured using a pick-up coil, and is shown in figure 2.56. The field is predominantly (> 90%) along the axis of the coil near the center of the coil, even with large metal objects such as a mock-up of the cloverleaf coils in place. This was surprising considering the wavelength and distance scales involved. The position of the coil relative to the magnetic trap bias field means that only $\Delta m_f = \pm 1$ transitions should be driven. In practice, we find that we can drive $\Delta m_f = \pm 1, 0$ transitions equally well. Presumably, the coil is not well centered on the atomic sample and fringing fields play a large role.



Figure 2.56: Position dependence of microwave field. The power in a pick-up coil oriented to measure the field along the axis of the microwave coil was measured using a spectrum analyzer. The pick-up was translated along the coil axis.

The overall microwave setup is shown in figure 2.57. The output from two synthesizers is combined using a Mini-Circuits ZAPD-2 power combiner. A voltage con-
trolled attenuator (VCA — Mini-Circuits ZAFS-2000) allows for fast computer control of the power. A 5 W amplifier (RFGA0012-05 from Richardson Electronics) feeds the microwave coil through a short length of 0.086 inch diameter semi-rigid coaxial cable. Measurements of the spectral properties of the microwave field indicate that this system produces radiation with no (< 90 dB) harmonics and excellent phase noise (< -100 dBc at 1 kHz).



Figure 2.57: Synthesizer setup. A schematic of the driver circuit for the synthesizer analog input can be found in the appendix to this chapter.

We use E4420B Agilent frequency synthesizers, and have also had (bad) experience with the Agilent 8656, 8646, and 8647 models. There are two primary problems that are encountered with many synthesizers: long frequency switching time and uncontrolled switching behavior. Long switching times limit the evaporation duty cycle (and effective microwave power) and speed (see chapter 4). The 8656 and 8646 models suffer from this problem, with switching times as long as 100 ms. The 8656, 8646, and 8647 models all suffer from frequency "glitches" during switching. The frequency is swept by issuing successive frequency set commands via GPIB control. Excursions larger than 1 kHz from the target frequency during switching lead to uncontrolled atom loss during the last stages of our evaporation. These glitches can be caught on a spectrum analyzer,²³ or by monitoring the atom number during evaporation. In order to limit the effect of these glitches, the VCA can be used to turn off the microwave power during the glitch. Unfortunately, this again limits the duty cycle and evaporation speed and leads to poor evaporation performance. The E4420B is an excellent model, with fast (< 40 ms) switching times and no glitches for frequency switching by less than 10 MHz. The E4420B only has to be attenuated during large frequency diversions and changes in the synthesizer power level. Furthermore, the E4420B has analog voltage controlled, phase continuous frequency sweep capability over 20 MHz which we use for rethermalization measurements (see chapter 6).

2.8 Anti-Gravity Coil

In order to perform experiments with spin mixed gases and extract information separately about each component, we need a method to separately image different spin states. We choose to spatially separate and simultaneously image the different spin components by applying a magnetic field gradient and essentially perform a Stern-

 $^{^{23}}$ To catch a glitch, the spectrum analyzer must be set to sample a small range in frequency fairly rapidly (typically 1 MHz in 60-100 ms). The spectrum analyzer must also be set to operate in a sample and hold mode, and a glitch can usually be caught after watching the synthesizer sweep ten times.

To separate the spin states for imaging, an 80 G/cm magnetic field gradient is applied during the expansion [73]. The gradient is created by a coil (cyan in figure 2.24) located on the bottom port of the science MOT cell. The "anti-gravity" (AG) coil has 220 turns, a 0.52 inch inner-diameter, 0.97 inch outer diameter, is 0.3 inches thick, and is 5 cm from the atoms. Electrically, the coil has a resistance of 3.3 ohms. The coil is quickly (< 100 μ sec) switched using a FET connected to a 60 V power supply (Xantrex XHR 60-18). The circuit is fused to support less than a rated 3 A DC current in order to prevent accidental heating of the science cell. The fuse itself limits the transient current to 18 A during the expansion. The coil is switched on for 9 ms during the expansion, with the image taken after 10 ms of total expansion time. A picture of the separated spin components is shown in figure 2.58.



Figure 2.58: Absorption image of separated spin components at low temperature. Because the atoms are initially loaded into a 45 G bias field trap, only states with $m_f > 1/2$ are trapped. The separation between the centers of each component is used to measure the AG coil magnetic field gradient.

After expansion, the centers of the $m_f = 9/2$ and $m_f = 7/2$ components are separated by ~ 430 μ m. The widths and shape of the gas are also altered. This problem is exacerbated in our experiment because a large gradient is required to separate atoms with a small difference in magnetic moments. The effect of the AG coil on the gas is calibrated by taking a series of images at different temperatures with and without the coil on during the expansion. Sample data used for this calibration is shown in figure 2.59. The effect on the widths from by the Thomas-Fermi and energy fits (see Chapter 5) appears to be fractional in the cloud size, even in the quantum degenerate regime. Calibrations from different dates are detailed in tables 2.7 and 2.8; the correction factor is defined as the measured width without the coil on divided by the measured width with the coil on. The shape of the cloud is also changed by the anti-gravity coil, so that Thomas-Fermi fits do not correctly determine the fugacity. This systematic is difficult to correct for, and has not been studied extensively. Regardless, the corrected number, temperature, and value of T/T_F agree with data taken with the AG coil off (figure 2.60).



Figure 2.59: Sample anti-gravity correction data. The Thomas-Fermi fit widths are shown with the AG coil on vs. off. A fit (solid line) to a line without an intercept is used to determine the correction factor.



Figure 2.60: Data corrected for the effect of the AG coil. This is the data used for figure 2.59, explicitly plotting the corrected values of N, T, and T/T_F . Linear fits (solid line) with no intercept indicate excellent agreement within the scatter.

Date	TF z	TF y	energy z	energy y	gaussian z	gaussian y
5/8/00	1.02(1)	1.12(2)	1.000(9)	1.084(7)	0.99(2)	1.09(1)
9/22/00	0.93(2)	1.06(2)	0.93(1)	1.06(1)	0.95(1)	1.07(1)
10/18/00	0.95(1)	1.06(1)	0.959(8)	1.062(9)	0.963(3)	1.070(9)

Table 2.7: Calibration of the effect of the AG coil on different dates for the $m_f = 9/2$ component. The correction factors for the widths from different fits are shown for each direction.

Date	TF z	TF y	energy z	energy y	gaussian z	gaussian y
5/8/00	0.997(9)	1.07(1)	1.008(4)	1.064(5)	1.001(8)	1.073(3)
9/22/00	0.940(5)	1.043(5)	0.946(6)	1.049(5)	0.949(4)	1.052(3)
10/18/00	0.974(5)	1.056(6)	0.969(1)	1.047(6)	0.965(6)	1.046(3)

Table 2.8: Different measurements of AG corrections for $m_f = 7/2$ component.

The AG coil correction factors are consistent with a classical, 2-dimensional model describing individual atom trajectories during the expansion. The coil is modelled as a loop of wire, with magnetic field magnitude in the z-y plane:

$$B(y,z) = \frac{\mu_0}{4\pi} m \frac{\sqrt{z^2 + 4y^2}}{(z^2 + y^2)^2}$$
(2.10)

where m is the loop magnetic dipole moment (m = IA where A is the loop area and I the current). The total force on an atom is then:

$$F_z(y,z) = \mu \frac{\mu_0 m}{4\pi} \frac{3(z^3 + 5zy^2)}{(z^2 + y^2)^3 \sqrt{z^2 + 4y^2}}$$
(2.11)

$$F_y(y,z) = \mu \frac{\mu_0 m}{4\pi} \frac{12y^3}{\left(z^2 + y^2\right)^3 \sqrt{z^2 + 4y^2}} - Mg$$
(2.12)

where μ is the atom magnetic moment, M is the mass, and g is the value of gravitational acceleration. Newton's differential equations for the atom position are then solved numerically for specific initial positions and velocities.

We perform the simplest calculation and consider specific atom trajectories without averaging over all initial positions and velocities. We consider two cases: atoms that start at the center of the cloud with equal and opposite velocities ($z_0 = y_0 = 0$, $v_0 = \pm \sqrt{\frac{k_b T}{m}}$), and atoms that start at opposite sides of the average diameter with the average and opposite velocities ($z_0 = \pm \sqrt{\frac{k_b T}{m\omega_z^2}}, v_{z0} = \pm \sqrt{\frac{k_b T}{m}}$ and $y_0 = \pm \sqrt{\frac{k_b T}{m\omega_r^2}}, v_{y0} = \pm \sqrt{\frac{k_b T}{m}}$). The correction factor is defined as the ratio between the distance between the atoms with and without the AG coil on during the expansion. For the first case the correction factors would be (regardless of the initial velocity): 1.07 ($m_f = 9/2$ radial), 0.98 ($m_f = 9/2$ axial), 1.06 ($m_f = 7/2$ radial), 0.98 ($m_f = 7/2$ axial). For the second case we find: 1.09 ($m_f = 9/2$ radial), 0.96 ($m_f = 9/2$ axial), 1.07 ($m_f = 7/2$ radial), 0.97 ($m_f = 7/2$ axial). These results are independent of the absolute temperature. The measured values are consistent with these calculated correction factors. The calculation also indicates that there is sensitivity to the initial position of the gas relative to the coil. Since the AG coil is not very mechanically stable, this may explain the scatter in correction factors indicated in tables 2.7 and 2.8.

2.9 Loading and Release Sequences

The science MOT and magnetic trap loading sequence is shown as a timeline in figure 2.61. The magnetic trap is turned on to the "load" trap after 2.3 ms of expansion in order to mode-match to the expanded MOT as closely as possible. The bias current is then ramped linearly over 100 ms to the "evap" trap and the first stage of evaporation begins.

Evaporative cooling to the degenerate regime takes 2-3 minutes. The release sequence is shown in figure 2.62. The magnetic trap is always ramped to the "tight" trap before the release. Note that we avoid turning on any other magnetic fields while the magnetic trap is on. The exception is the "pcoil" which is turned on after the bias current is off but while some main current is still on. Even the quantization fields used



Figure 2.61: Magnetic trap loading timeline.

for imaging can drive significant center-of-mass motion in the expanding gas if they are turned on before the release. The "pcoil" preserves the spin orientation of the atoms while allowing time for the imaging quantization fields to turn on. In figure 2.62 there are two release sequences — one for E/W (with dashed segments) probing and the other for U/D probing (without dashed segments). For U/D probing, the "pcoil" provides the quantization axis. The E/W field takes ~ 2 ms to turn completely on, so the actual turn-on time has some spread.



Figure 2.62: Magnetic trap release sequence.

2.10 Computer Control

The computer control system used in the experiment is not explained thoroughly in this thesis. We copied the scheme used in Eric Cornell's BEC experiment. TTL, DAC, GPIB, and ADC computer boards are controlled from a QuickBASIC program with built-in low-level routines. Simple circuits are used to sum analog and TTL signals, as well as convert TTL signals to low-noise analog control voltages that are controlled by a potentiometer. The computer program is phase-locked to the 60 Hz line cycle, which is remarkably stable. The 60 Hz phase locked loop (PLL) is frequency multiplied to provide a 1/3 ms time base for computer events. Events requiring fast timing are controlled by external delay and pulse generator circuits, where are triggered by the computer TTL signals. This system is shown schematically in figure 2.63. Data acquisition occurs primarily from the Princeton Instruments camera used to take images. The CCD array is read-out using a combination of Princeton Instruments supplied OCX controls and home-brewed Visual Basic programs. These programs then fit and analyze the images.



Figure 2.63: Computer control schematic.

2.11 Appendix to Chapter 2

2.11.1 MOPA mounts

Figures 2.66, 2.65, and 2.64 are the most up-to-date drawings for the specialized MOPA mounts, with corrections by Neil Claussen to the original schematics. The lens mounts are fixed to the double flexure with stainless steel screws (holes "C"), and are designed to hold the DO-818 lenses. The center arm of the double flexure mount is fixed to the baseplate with two screws (holes "D"), and the PA coupling mount is attached to the double flexure mount via two screws (holes "E"). The SDL C-Mount fits snugly into the PA coupling mount. A connection to the cathode pin is usually constructed from an electrical socket and a Delrin rod — this is glued into place so that the cathode pin fits snugly into the socket.



Figure 2.64: PA double flexure mount.



Figure 2.65: PA lens mount.



Figure 2.66: PA coupling mount.

2.11.2 PA Current Driver

Figure 2.67 is the circuit diagram for the PA current driver, courtesy of James Fung-A-Fat. Switch S3 can be used to short the PA for protection when it is off. We have observed that transients can appear across a load similar to the PA connected to the current driver through a long (20 foot) cable when the supplies PS1 and PS2 are switched on or off (even with S3 shorted to ground). These transients have been suppressed to undetectable levels with the additions of switches S1, S2, and the solid-state relay K1. K1 forces the high current supply PS2 to turn on or off only on a zero crossing in the 60 Hz wall power cycle. This prevents PS2 from producing transients across the load. The turn-on procedure is: S3 is shorted to ground and the driver is set to zero current, turn-on S1, then turn-on S2. The turn-off procedure is current to zero, short S3 to ground, turn-off S2, then turn-off S1.



Figure 2.67: PA current driver.

2.11.3 Synthesizer Driver Circuit

Driving the modulation input on the E4420B carries two dangers, notably damage to the computer DAC and/or damage to the E4420B input stage. A typical computer DAC barely has the current drive capability to output 1 V across 50 Ohms. Repeatedly trying this will eventually damage the output stage of the DAC. If a high current diver is improperly connected to the synthesizer input, the 50 Ohm input stage can easily by overdriven and damaged. The circuit shown in figure 2.68 is designed to safely allow a computer DAC to drive the modulation input. This circuit has a 50 Ohm output impedence, and enough current drive capability to drive a few volts across 50 Ohms. The output is limited by 1.2 volts by the diodes. The input is divided down by a factor of ten in order to increase the DAC resolution. The circuit has a built in low-pass at ~ 3 kHz.



Figure 2.68: Synthesizer modulation input driver.

2.11.4 XTreme Diode Cooling

We learned many lessons the hard way when it came to using multiple stages of thermo-electric coolers (TEC's) to cool diode lasers. There are general design issues to keep in mind when attempting to cool diodes. Peltier cooling efficiency is actually quite low, temperature dependent, and dependent on the temperature difference across the device. The Melcor Thermoelectric Cooling handbook has a wealth of information on TEC's. A TEC can be modelled as a device the removes heat at a rate \dot{Q}_c from the cold side (at T_c) and expels heat at a rate \dot{Q}_h from the hot side (at T_h):

$$\dot{Q}_c = 2N \left[\alpha I T_c - \frac{I^2 \rho}{2G} - \kappa \left(T_h - T_c \right) G \right]$$
(2.13)

$$\dot{Q}_h = 2N\alpha I T_c \tag{2.14}$$

where I is the device current, N is the number of p-n junctions, and α, G, κ, ρ are device constants, and temperature is measured in K. For the devices used in the work described by this thesis, some approximate values are $\alpha = 2.0 \times 10^{-4} \text{ V/K}$, $\rho = 1.0 \times 10^{-3} \text{ ohm} \cdot \text{cm}$, and $\kappa = 1.5 \times 10^{-2} \text{ W/cm/K}$. The value of N is obtained from the second set of digits in the part number. For part number CP-1.4-127-045L G = 0.171, for CP-0.8-127-06L G = 0.42, and for OT-1.5-32-F0 G = 0.03. The second term in the equation for \dot{Q}_c represents Ohmic heating in the device, while the third term is reverse heat flow due to the junction thermal conductivity. Note that the device efficiency falls to zero around $\Delta T = T_h - T_c = 40^{\circ}$ C, and that the optimal operating current is always around 80% of the specified maximum current. In reality, we have found it difficult to support ΔT greater than 30 °C across a single stage.

Working with multiple stages of cooling carries unique design constraints. Multiple stages are used when trying to reach very large ΔT 's, and actually tend to reduce the available cooling power. Each stage must be able to handle the load from the previous stage as well as move some extra heat. It is not necessary to always stack the TEC's on top of each other, and, in fact, there are some good reasons not to. Experimentally, we have found that the heat load on the diode mount comes from convective transfer from the closest neighbor and the laser housing. For this reason, it is advantageous to cool the lens mount in the laser, for example, to below room temperature. Also, foam insulation between objects at different temperatures is very helpful in reducing heat loads. For example, a thin layer of insulation between the laser baseplate and the heatsink can reduce the heat load on the baseplate by 10-15%. For the type of mounts that we use, it is difficult to measure any temperature gradient across aluminum parts that are between TEC's.

In general, good thermal contact is only assured by connecting multiple stages together with stainless steel screws. The heads of the screws are insulated with a combination of nylon and fiber washers. The surfaces of aluminum parts should be lapped and thoroughly cleaned — flycutting tends to cause surfaces to bow. A very thin layer of thermal compound is helpful to avoid problems caused by pits or machine marks. The thermal compound should only be applied with a clean tool, as oils and grease can significantly diminish its effectiveness.

The total heat load when using multiple stages is multiplied by the inefficiency of each stage. It is easy to generate at least 60 W at the end of three stages of cooling that bring a diode down to -40 °C. Removing this heat can be difficult, and the heatsink thermal resistance becomes the actual limit on the lowest temperature that a system can achieve. Water cooling works best for the heatsink, but seems to introduce a lot of frequency noise into ECDL's. We have achieved a (measured) thermal resistance of ~0.2 °C/W with forced convection across heatfins. The repump laser can cool to -40 °C with forced convective cooling of the heatsink *and* the laser attached to the large aluminum MOPA plate. Without forced convection the laser can only cool to -30 °C. The fan used for forced convection should not be attached directly to the optical table in order to avoid vibrational noise.

Avoiding water condensation and frost is also important, although we have found

that most laser diodes are amazingly robust against water damage. Note that we never remove the diode from the usual hermetically sealed package. The presence of frost is indicated by a drop in the laser power with no change in the threshold current. We have found that duct taping the "seams" of the laser cavity, sealing electrical feedthroughs with epoxy, sealing countersunk holes with epoxy covered duct tape, and putting a small box of dessicant in the cavity reduces the frost accumulation to the point that the laser must only be cleaned out every three months. In addition, the latex feedthrough scheme detailed in the text is necessary of low temperature adjustment of the laser wavelength, feedback, and collimation.²⁴

Both stages of cooling in the peak-locked laser are controlled with PID servo's — the temperature is measured using a 30 kOhm thermistor. In the repump laser, the laser housing and diode mount are servo'd, while the flexure TEC is controlled with a constant-current power supply. We are careful to avoid thermal runaway, and the TEC equations are useful for checking the heat load on each stage. The repump diode is cooled to -40 °C with the diode mount TEC's at 0.422 A, the flexure TEC's at 1.35 A, the flexure at -22 °C, the heatsink at 30 °C, and the laser housing at 14 °C (the housing TEC's are at 3.85 A). Under these conditions, the TEC's under the diode mount are removing 1.1 W at their cold side, and the final stage is moving ~ 55 W into the heatsink.

 $^{^{24}}$ The laser collimation and feedback alignment shift significantly from room temperature to -40 °C.

Chapter 3

COLLISIONS — ⁴⁰K IS MAGICAL!

3.1 Overview

The collisional properties of fermionic atoms are the fundamental challenge to evaporative cooling. Experiments with ultra-cold atomic gases are usually done with spin polarized samples in order to avoid loss due to spin changing collisions that transfer atoms into untrapped states. Evaporative cooling requires high collision rates so that the gas rethermalizes to lower temperature after ejecting the highest energy atoms. However, completely identical fermions do not collide at low temperature. Therefore, a spin polarized fermionic gas cannot rethermalize and evaporation will fail.

Our plan to overcome this problem was to trap two spin states of 40 K ($m_f = 9/2$ and $m_f = 7/2$ in the f = 9/2 ground state) that are stable against spin exchange collisions. We would then simultaneously evaporatively cool the two gases, each maintaining thermal equilibrium through its contact with the other. Although sympathetic cooling of bosonic atoms to degeneracy had been demonstrated before [18], this was a more complicated scheme and one that raised many questions. For example, we were unsure when we started exactly how the mechanics of the evaporation would work. We did not know how to prepare the spin mixture that we needed, and the importance of the spin mixture purity was uncertain. At that time, the magnitude and sign of the triplet scattering length of 40 K was essentially unknown [86,87]. We did not know if the elastic s-wave cross-section would be large enough for evaporation to work. For that matter, we did not know if inelastic losses due to spin exchange, dipolar relaxation, or three-body collisions would cause trouble. We were in for a few surprises, notably that ⁴⁰K is an unusual atom (even among the alkali's) with collisional properties that are especially advantageous for evaporative cooling.

This chapter will explain some of the first measurements [42] that we made of the scattering properties of ⁴⁰K. Using cross-dimensional rethermalization, we measured the p-wave and s-wave collision cross-sections. We also measured the p-wave threshold behavior vs. temperature. We discovered a p-wave shape resonance due to a bound state (in the inter-atomic potential) close in energy to the p-wave centrifugal barrier. The simultaneous measurement of the s- and p-wave collision cross-section allowed us, with the help of JILA theorists, to make a relatively precise determination of the triplet scattering length a_t . The reader should keep in mind that these measurements were performed before we installed the AG coil, and before we developed the two-frequency evaporation technique that allowed us to cool to degeneracy. Further, we used a very simplified model of the rethermalization — a more complete treatment can be found in chapter 6. At the end of this chapter, I will mention some measurements that provided an upper limit on the spin exchange rate for certain combinations of spin states. The ultra-low inelastic rate that we found was surprising. However, comparison with theoretical calculations using the currently best known values of $C_6 = 3897$ a.u. $[88]^1$ and $a_t = 169 a_0 [89]$ where a_0 is the Bohr radius ² reveal that all inelastic rates in 40 K are very low.

Note that in this chapter we use the symbols x_{rms} and z_{rms} for the radial and axial rms size of the gas. This is done in order to be consistent with published work and avoid confusion with the collision cross-section. In other chapters, σ will be used for the size of the gas.

 $^{^{1}}$ This value comes from a recent high precision calculation, with an uncertainty of only 15 a.u.

² Recent two-photon photoassociative spectroscopy of 39 K and mass rescaling was used to determine this value, with an uncertainty of only 9 a₀.

3.2 Colliding Fermions

Quantum mechanical exchange symmetry limits the type of two-particle wavefunctions allowed in collisions between fermionic atoms. For identical fermions, the wavefunction must be anti-symmetric under exchange of the atoms. It is a common point of confusion that ⁴⁰K atoms in different Zeeman levels are, in fact, identical (or *indistinguishable*) in the context of scattering. The physical observables in a collision are the initial and final relative momenta $\vec{k'}$ and \vec{k} and the initial and final spin state of the atoms. The collision itself is a quantum mechanical "black-box" where the exact collision process for any particular collision cannot be known. Consider an elastic collision between atoms a and b with $m_f = 9/2$ and $m_f = 7/2$. In the final state we see atoms with $m_f = 9/2$ and $m_f = 7/2$, but we cannot be sure which is a and which is b because the scattering potential is spin dependent and can change the spin of an atom during the collision. For this reason, the two-particle wavefunction must be properly symmetrized for exchange.

We start by considering the properties of quantum-mechanical elastic collisions. For the sake of clarity, the properties of elastic and inelastic collisions are discussed separately. The solution to the Lippman-Schwinger equation [90] using a partial wave expansion is useful for understanding the properties of colliding fermions. Solving the Lippmann-Schwinger equation for the l^{th} partial wave of the asymptotic³ wavefunction $|\Psi\rangle$ [90] in an elastic collision gives:

$$\langle \vec{x} | \Psi \rangle_l \propto P_l(\cos \theta) \left\{ \left[1 + 2ikf_l(k) \right] \frac{e^{ikr}}{r} - \frac{e^{-i(kr-l\pi)}}{r} \right\} | \chi \rangle$$
 (3.1)

where r inter-nuclear distance, $P_l(\cos\theta)$ is the l^{th} Legendre polynomial evaluated as a function of the angle between the initial and final relative momenta (we assume k = k'), $|\chi\rangle$ is the spin wavefunction, and f_l is the scattering amplitude. The first term in the brackets is the outgoing wave, modified in amplitude and phase by the action of the

³ Meaning large inter-nuclear distance.

scattering potential. The second term represents the incoming spherical wave. The usual symmetrization of the space part via the anti-symmetrization operator has not been done so that we can consider different cases. Swapping particles involves the transformation $\cos(\theta) \rightarrow \cos(\pi - \theta)$ and changing $|\chi\rangle$ appropriately. In order to be a valid wavefunction for FD particles, we must have $\langle \vec{x} | \Psi \rangle_l \rightarrow -\langle \vec{x} | \Psi \rangle_l$ for this transformation. The overall symmetry of the wavefunction is determined by the product of the symmetry of the spatial and spin parts.

For atoms in the same spin state, the spin wavefunction is always symmetric under exchange: $|s_1, s_1\rangle \rightarrow |s_1, s_1\rangle$. Here, the position inside the Dirac bra-ket denotes which atom is in which state. The symbols s_1 and s_2 will be used to denote different spin states (as in the $m_s = \pm 1/2$ and $\pm 1/2$ states of an electron or $m_f = 9/2$ and $m_f = 7/2$ for ⁴⁰K atoms) in order to be general. Elastic collisions of atoms in the same spin state can only occur on l odd entrance channels⁴ since the spatial part of the wavefunction is anti-symmetric for l odd and symmetric for l even. This comes from the behavior of the Legendre polynomials under rotation: $P_l(cos(\theta)) = P_l(cos(\pi - \theta))$ for l even and $P_l(cos(\theta)) = -P_l(cos(\pi - \theta))$ for l odd. However, for atoms in different spin states $(s_1 \neq s_2), |\chi\rangle$ can be anti-symmetric

$$\frac{1}{\sqrt{2}} \left(|s_1, s_2\rangle - |s_2, s_1\rangle \right) \to -\frac{1}{\sqrt{2}} \left(|s_1, s_2\rangle - |s_2, s_1\rangle \right) \tag{3.2}$$

or symmetric

$$\frac{1}{\sqrt{2}} \left(|s_1, s_2\rangle + |s_2, s_1\rangle \right) \to \frac{1}{\sqrt{2}} \left(|s_1, s_2\rangle + |s_2, s_1\rangle \right) \tag{3.3}$$

under exchange. Atoms in mixed spin state gases can undergo elastic collisions on both odd *and* even partial wave channels. This is our trick for evaporation — at low temperature, where only s-wave collisions are allowed, a mixed spin-state gas can still maintain a high collision rate. Collisions occur between, but not within, the two components.

 $^{^4}$ It is rather remarkable to realize that atoms in the same spin state do not ever approach each other in a head-on collision.

Collisions on channels higher than s-wave (l=0) are suppressed at low temperature because of the centrifugal barrier. The threshold energy E_{th} for a given partial wave to contribute to scattering can be approximated by the centrifugal barrier,

$$E_{th}(l) = \frac{\hbar^2 l(l+1)}{2mb^2} - \frac{C_6}{b^6}$$
(3.4)

$$b^2 = \left[\frac{6C_6m}{\hbar^2 l(l+1)}\right]^{1/2} \tag{3.5}$$

where the currently best known value of C_6 for ⁴⁰K is 3897 a.u., m is the mass of either colliding partner (40 a.m.u. for our case), and b is the radius of the local maximum in the effective potential. The threshold energy for p-wave (l = 1) collisions and d-wave (l = 2)collisions are 100 μ K and 500 μ K, respectively. Since we start with a MOT temperature of 150 μ K and evaporate to less than 1 μ K, collisions are primarily s-wave in nature, with a significant p-wave contribution only at the beginning of the evaporation. The contribution from d-wave collisions is unknown but may play a role once the atoms are compressed and adiabatically heated after transferring from the science MOT into the magnetic trap.

Next, we consider the properties of inelastic collisions, specifically spin-exchange collisions where the incoming and outgoing spin states are different. It is important to avoid loss due to inelastic collisions in the experiment, and a measurement of the spin-exchange rate constant will appear later in the chapter. Spin-exchange collisions must not only satisfy FD statistics but also conserve the total projection of spin. For this reason a spin-exchange collision between an $m_f = 9/2$ and $m_f = 7/2$ atom is forbidden because there is no other final state at lower energy such that the total $m_f = 8$. Furthermore, the spin-exchange collision $m_f = 7/2 + m_f = 7/2 \rightarrow m_f = 9/2 + m_f = 5/2$ is forbidden by FD statistics for s-wave collisions. A mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms is then stable against spin relaxation when only s-wave collisions are allowed.⁵

⁵ Dipolar relaxation, where two atoms enter on an s-wave channel and exit on a d-wave channel, can cause spin relaxation. This rate is very low, however, because we work at temperatures where the collision energies are far below the d-wave centrifugal barrier.

In the experiment, we also have $m_f = 5/2$ and $m_f = 3/2$ atoms present during some of the evaporation. There are some spin exchange collisions that are still not allowed, $m_f = 5/2 + m_f = 9/2 \rightarrow m_f = 7/2 + m_f = 7/2$ for example. This can be understood as the time-reversed version of a forbidden s-wave collision, and, since the scattering Hamiltonian is time-reversal invariant, this collision is also forbidden. Alternatively, one can consider the l^{th} partial wave of the asymptotic wavefunction again (this time written in a way useful for considering spin-exchange collisions):

$$\langle \vec{x} | \Psi \rangle_l \propto P_l \left(\cos \theta \right) \left\{ \left[1 + 2ik f_l(k) \right] \frac{e^{ikr}}{r} |\chi_{out}\rangle - \frac{e^{-i(kr-l\pi)}}{r} |\chi_{in}\rangle \right\}$$
(3.6)

where the spin wavefunctions for the in and out states have been separated. For even partial waves, this wavefunction cannot be properly symmetrized for identical final spin states — the outgoing part cannot have odd exchange symmetry. This collision can only, therefore, occur on odd partial waves. Because we trap atoms in the lower hyperfine ground state, hyperfine changing collisions are energetically forbidden. There is not enough kinetic energy for a collision where $f = 9/2 \rightarrow f = 7/2$ below 10 mK.

3.3 Definition of Cross-Section

In the experiment, we measure collision, or rethermalization, rates in the gas, which we then connect to the collision cross-section. The parameter of interest to theorists is the scattering length, which can be connected to the collision cross-section. The point of this section is to explain the connection between the collision rate, collision cross-section, and scattering length clearly. We will pay particular attention to using definitions that work for both the experimentalists and theorists.

Jim Burke has done an excellent job of explaining the quantum mechanics of cold atom scattering in his thesis [91], and then deriving the appropriate cross-sections in terms of the scattering length. As he mentions, his method is easily adapted to fermions by changing the symmetrized wavefunctions appropriately. Modifying Jim's equation 2.21, for elastic scattering of fermions we have:

$$\Psi_i^I = \left(\frac{e^{ik_iZ} - e^{-ik_iZ}}{\sqrt{2}}\right) |s_1, s_1\rangle \tag{3.7}$$

$$\Psi_i^D = \frac{e^{ik_i Z} |s_1, s_2\rangle - e^{-ik_i Z} |s_2, s_1\rangle}{\sqrt{2}}$$
(3.8)

for the incoming plane waves, and in 2.22

$$\Psi_f^I = \left(\frac{e^{i\vec{k}_f \cdot \vec{R}} - e^{-i\vec{k}_f \cdot \vec{R}}}{\sqrt{2}}\right) |s_1, s_2\rangle \tag{3.9}$$

$$\Psi_f^D = \frac{e^{i\vec{k}_f \cdot \vec{R}} |s_1, s_2\rangle - e^{-i\vec{k}_f \cdot \vec{R}} |s_2, s_1\rangle}{\sqrt{2}}$$
(3.10)

for the outgoing state. In these equations, I refers to the case of scatterers with identical spin states and D to the case of scatterers with different spin. Jim's cases (i)-(iv) can be calculated with modifications to the sums and to the wavefunctions expanded in spherical harmonics in order to preserve the proper symmetry. The final result for the elastic collision (case (i) and (iv)) cross-section for fermions is the same as for bosons: $4\pi a^2$ for atoms with different m_f and $8\pi a^2$ for atoms with the same m_f , where a is the relevant scattering length. The s-wave elastic collisions between atoms in different spin states described in this chapter are primarily triplet in nature, and the relevant cross section is $4\pi a_t^2$ where a_t is the s-wave triplet scattering length. A scattering length can be formally defined for p-wave collisions (J. Bohn, private communication). The relevance to the collisions that we consider is that there is a factor of two difference in the p-wave collision cross-section between atoms in the same internal state and atoms in different internal states.

The factor of two difference in the cross-section between atoms in different compared to the same spin states often seems mysterious but can be understood from a simple point of view. The rate of elastic collisions between two atoms for a given partial wave is proportional to

$$\int d(\cos\theta) \, d\phi \, \left| \langle \psi_{in} | T | \psi_{out} \rangle \right|^2 \tag{3.11}$$

where the angular integration is carried out over the angles θ and ϕ between \vec{k} and $\vec{k'}$, and T is the scattering matrix. In order to satisfy exchange symmetry we must have:

$$|\psi_{in}\rangle = \frac{1}{\sqrt{2}} \left(|s_1 s_2, \vec{k}\rangle - |s_2 s_1, -\vec{k}\rangle \right) \tag{3.12}$$

$$|\psi_{out}\rangle = \frac{1}{\sqrt{2}} \left(|s_1 s_2, \vec{k}'\rangle - |s_2 s_1, -\vec{k}'\rangle \right)$$
 (3.13)

as exchanging particles involves swapping the spin of each atom $(s_1 \leftrightarrow s_2)$ and changing the sign of the relative momenta $(\vec{k} \rightarrow -\vec{k} \text{ and } \vec{k'} \rightarrow -\vec{k'})$. In general, T is only a function of k, θ , and the spin of the scatterers. Using momentum conservation, we find that equation 3.11 becomes⁶

$$\int d(\cos\theta) \, d\phi \, |\langle s_1 s_2 | T_k(\theta) | s_1 s_2 \rangle - \langle s_1 s_2 | T_k(\pi - \theta) | s_2 s_1 \rangle|^2 \tag{3.14}$$

where the first and second terms are the "direct" and "exchange" terms and we have implicitly left in a possible spin dependence for T_k . Consider p-wave collisions (which are the only collisions for our experiment that can exhibit this factor of two) so that we can have $s_1 = s_2$ or $s_1 \neq s_2$, and $T_k^l(\theta) = -T_k^l(\pi - \theta)$. If $s_1 = s_2$ then the direct and exchange terms are exactly equal and add, giving $4T_k^l(\theta)^2$ inside the integral. If $s_1 \neq s_2$, the exchange term for alkali atoms is generally much smaller than the direct term, so that we have $T_k^l(\theta)^2$ inside the integral. For $s_1 = s_2$, however, the integral over θ covers only $0 \rightarrow \pi/2$ in order to avoid double counting. In the end, then, we reproduce the factor of two difference in the rate of collisions. It is important to realize that this is only approximately a factor of two, with the deviation coming from the subtraction of the small exchange term. The importance for our system is that the rate of p-wave collisions between atoms with the same m_f is approximately twice as high as for atoms with different m_f .

In practice, these definitions are useful for connecting experimentally measured rethermalization rates to the scattering length and cross-section. It can be difficult to

⁶ To get this result you must use: $T(\vec{k}', \vec{k}) = T_k(\theta), T(-\vec{k}', -\vec{k}) = T_k(\theta), T(-\vec{k}', \vec{k}) = T_k(\pi - \theta),$ and $T(\vec{k}', -\vec{k}) = T_k(\pi - \theta).$

drag out of collision cross-section calculating theorists how their cross-section connects to the collision rate in the gas. Their theory is developed for two particles in a box, using Fermi Golden's rule to determine the number of collisions per volume per second:

$$K_2 = \hbar \frac{k}{\mu} \sigma \tag{3.15}$$

where k is the magnitude of the incident relative momentum and μ is the reduced mass. The relative speed $(|\vec{v}_1 - \vec{v}_2|)$ in the collision is just $\hbar \frac{k}{\mu}$.

For a trapped gas, we have to average K_2 over all colliding partners in order to determine the total number of collisions in the gas per second Γ . We assume that the colliding partners must be at the same place in space, since the range of the interaction is shorter than all other relevant length scales. The total number of collisions per second between a 1 and a 2 atom is then:

$$\Gamma_{12} = \frac{1}{N_1 N_2} \int d^3 \vec{r_1} \int d^3 \vec{r_2} \int d^3 \vec{p_1} \int d^3 \vec{p_2} \, \delta^3 (\vec{r_1} - \vec{r_2}) \mathcal{F}_1(\vec{r_1}, \vec{p_1}) \, \mathcal{F}_2(\vec{r_2}, \vec{p_2}) \frac{|\vec{p_1} - \vec{p_2}|}{m} \, \sigma$$

$$= \frac{1}{N_1 N_2} \int d^3 \vec{r} \int d^3 \vec{p_1} \int d^3 \vec{p_2} \, \mathcal{F}_1(\vec{r}, \vec{p_1}) \, \mathcal{F}_2(\vec{r}, \vec{p_2}) \frac{|\vec{p_1} - \vec{p_2}|}{m} \, \sigma \qquad (3.16)$$

where \mathcal{F}_1 and \mathcal{F}_2 are the statistical distribution functions for the colliding atoms, the differential cross section is assumed to be isotropic with total cross section σ , and the mass of the colliding atoms is m (which is the same for atom 1 and 2). For Maxwell-Boltzmann distributions, \mathcal{F} is separable in \vec{r} and \vec{p} , and we define $n_{12} = \int d^3 \vec{r} n_1(\vec{r}) n_2(\vec{r})$ and the mean relative speed $v = \frac{1}{N_1 N_2} \int d^3 \vec{p}_1 \int d^3 \vec{p}_2 \Pi_1(\vec{p}_1) \Pi_2(\vec{p}_2) \frac{|\vec{p}_1 - \vec{p}_2|}{m}$. Both the density and momentum distributions $n(\vec{r})$ and $\Pi(\vec{p})$ may be anisotropic if the gas is out of crossdimensional equilibrium. With substitution into equation 3.16, we end up with the total collision rate as

$$\Gamma_{12} = n_{12}\sigma v \tag{3.17}$$

which is the familiar definition of the cross section from considering a beam-target experiment. In order to determine the collision rate per particle, Γ_{12} must be divided by the number of particles N_1 or N_2 . If the mixture of particles 1 and 2 is not equal, then each gas may have a very different collision rate per particle.⁷

3.4 Elastic Collision Measurement

We determine elastic collision cross sections [42] from measurements of crossdimensional thermalization rates [92] in the magnetic trap. The sample is taken out of thermal equilibrium by changing the radial harmonic trap frequency ν_r through a ramp of the bias coil current. For the measurements reported here ν_r lies between 44 and 133 Hz. The change in ν_r occurs adiabatically (slow compared to the atomic motion in the trap) but much faster then the rate of collisions between atoms. Since the axial frequency is essentially unchanged, energy is added to (or removed from) the cloud in only the radial dimension. Elastic collisions then move energy between the radial and the axial dimensions, and the thermal relaxation is observed by monitoring the time evolution of the cloud's aspect ratio.

To avoid perturbations to the image due to the spatially dependent magnetic fields, the trap is turned off suddenly and the cloud is imaged after 2.7 ms of free expansion. The aspect ratio of the cloud is observed via absorption imaging using a 9.1 μ s pulse of light resonant with the 4S_{1/2}, f = 9/2 to 4P_{3/2}, f' = 11/2 transition. Optical depth is calculated from the image captured on a CCD array and then surface fit to a gaussian distribution to find the rms cloud size in both the radial and axial dimensions. An example of the cloud evolution following a change in trap potential is shown in figure 3.1. Since the expanded cloud sizes are proportional to the square root of the cloud energy in each dimension, the exponential time constant for the redistribution of energy, τ , can be extracted from an appropriate fit (see the appendix to this chapter) to the aspect ratio vs time. To rule out significant relaxation through trap anharmonicities,

⁷ This may seem strange at first, but consider s-wave collisions between a million $m_f = 9/2$ and a hundred $m_f = 7/2$ atoms. Each $m_f = 9/2$ atom is colliding at a much lower rate than each $m_f = 7/2$ atom, because there are far fewer available colliding partners.

we have verified that the relaxation rate $1/\tau$ scales linearly with the number of trapped atoms N.



Figure 3.1: Example of thermalization data. The inset shows the axial size, z_{rms} (\circ), and the radial size, x_{rms} (\bullet), imaged after 2.7 ms of free expansion, relaxing as the trapped atoms rethermalize via elastic collisions. At time=0 the cloud is taken out of equilibrium by changing ν_r from 133 to 44 Hz. A fit (line) to the aspect ratio vs time is used to extract the relaxation rate.

To obtain the elastic collision cross section σ from our measurements of thermal relaxation rates, we use the relation: $1/\tau = \frac{2}{\alpha}n\sigma v$, where *n* is the density-weighted density of the trapped atoms given by $\frac{1}{N}\int n(r)^2 d^3r$, *v* is the rms relative speed between two atoms in the trap, and α is the calculated average number of binary collisions per atom required for thermalization. The product nv depends on both the size and temperature, *T*, of the trapped sample. These are measured by observing the expansion of an equilibrated sample after release from the magnetic trap. The rate of expansion yields the temperature, while an extrapolation back to the release time gives the initial sizes. Using the trap potential calculated from the field coil geometry we have checked that the measured initial sizes and temperatures are consistent to within their uncertainties.

The mean number of collisions each atom undergoes, α , during one relaxation time constant was determined by Murray Holland from a numerical simulation of the experiment using classical Monte Carlo methods. ⁸ For a harmonic trapping potential, the relaxation simulation yields $\alpha_s = 2.5$ for s-wave collisions and $\alpha_p = 4.1$ for p-wave collisions, where we include explicitly the different angular dependence. The ratio α_p/α_s can also be determined analytically through an integration over the angular dependence of scattering. This gives $\alpha_p/\alpha_s = 5/3$, consistent with the Monte Carlo results.

The primary results of this measurement are shown in figure 3.2. While ordinarily one cannot measure higher order partial wave contributions to the collision cross section directly, the Fermi-Dirac statistics of 40 K allow us to probe p-wave and s-wave interactions independently. The p-wave cross section σ_p is determined from measurements using a spin-polarized sample ($|f = 9/2, m_f = 9/2\rangle$ atoms), where s-wave collisions are prohibited by the quantum statistics. The s-wave cross sections σ_s are determined from data obtained using a mixture of two spin states, $|9/2, 9/2\rangle$ and $|9/2, 7/2\rangle$. This mixture is stable against inelastic spin exchange collisions (s-wave), which must preserve the total projection of spin. The magnitude of the p-wave cross section is surprisingly

 $^{^{8}}$ For more details see M. Holland.

large and we find that ⁴⁰K has a p-wave shape resonance at a collision energy of roughly 280 μ K. At temperatures well below the resonant energy (less than 30 μ K), a fit to σ_p vs T gives $\sigma_p \propto T^{2.0\pm0.3}$. Thus, we have directly measured the expected threshold behavior $\sigma_p \propto E^2$ [93]. In contrast, σ_s exhibits little temperature dependence. With these very different temperature dependencies, the collision rate changes by over two orders of magnitude at our lowest temperatures depending on the spin mixture of the fermionic atom gas.



Figure 3.2: Elastic cross sections vs. temperature. The s-wave cross section (\circ), measured using a mixture of spin states, shows little temperature dependence. However, the p-wave cross section (\bullet), measured using spin-polarized atoms, exhibits the expected threshold behavior and is seen to vary by over two orders of magnitude. The lines are a fit to the data, as described in the text, yielding $a_t = 157 \pm 20a_0$.
To explore this effect further we measure the thermalization rate vs spin polarization at 9 μ K (see figure 3.3). We control the relative populations of $|9/2, 9/2\rangle$ and $|9/2, 7/2\rangle$ atoms in a two-component cloud with a microwave field that drives transitions to untrapped spin states in the F=7/2 ground state manifold. The trap bias magnetic field breaks the degeneracy of the hyperfine ground-state splitting (1.286 GHz at zero field [94]) so that the different spin-states can be removed selectively (see figure 3.3 inset). For the data shown in figure 3.3 the fraction of atoms in the $|9/2, 9/2\rangle$ state $f_{m_f=9/2} \equiv \frac{N_{m_f=9/2}}{N_{m_f=9/2}+N_{m_f=7/2}}$, was varied smoothly from 70 to 100% by varying the power of an applied microwave field (frequency swept) that removes a portion of the $|9/2, 7/2\rangle$ atoms.



Figure 3.3: Dependence of collision cross-section σ on spin composition at $T = 9\mu K$. A quantity proportional to σ , $1/\tau N$, is measured as a function of the fraction of atoms in one of two trapped Zeeman spin states, $f_{m_f=9/2}$. The inset shows the number of atoms remaining after application of microwaves at the indicated frequency. The three features correspond to removal of trapped atoms in particular spin states and can be used to measure or control the spin composition of the atom gas.

The thermalization of mixed spin-state samples depends on both s-wave and pwave collisions. The data in figure 3.3 can be fit to a simple model given by:

$$1/\tau = n_{1,2}(\frac{2}{\alpha_s}\sigma_s + \frac{2}{\alpha_p}\frac{\sigma_p}{2})v + (n_{1,1} + n_{2,2})\frac{2}{\alpha_p}\sigma_p v,$$

where $n_{i,j}$ is the density-weighted density between two species given by $n_{i,j} = \frac{1}{N_1+N_2} \int n_i(r)n_j(r)d^3r$ and the subscripts 1 and 2 stand for the two relevant spin states. Since the magnetic moments of the $|9/2, 9/2\rangle$ and $|9/2, 7/2\rangle$ atoms are only slightly different we make the simplifying assumption that these states have identical spatial profiles in the trap. A fit using the above model with σ_s and the ratio σ_p/σ_s as free parameters shows good agreement with the data in figure 3.3. In addition to demonstrating the type of control over collision rates that is available in a trapped gas of fermionic atoms, this measurement of σ_p/σ_s at low temperature provides a sensitive constraint on the triplet scattering length. The s-wave cross-sections shown in figure 3.2 were extracted using the above equation, however at these low temperatures σ_p is relatively small and the measured thermalization rates are due primarily to s-wave interactions.

To compute the scattering cross sections for comparison with these data, our collaborators Jim Burke and John Bohn perform fully coupled scattering calcuations, including potassium hyperfine structure, as detailed in [95]. These calculations include additional corrections that enable them to tune independently the singlet and triplet scattering lengths over their entire ranges $-\infty < a < \infty$. They set the singlet scattering length's value at $a_s = 104a_0$ [95] where a_0 is the Bohr radius, but leave the triplet scattering length a_t as a free parameter to be determined by the experiment. Scattering is strongly triplet-dominated for the collision partners used in this experiment, and they find no singlet resonance near threshold. Therefore, each partial wave's cross section has little dependence on the spin states involved. Moreover, the present experiment is relatively insensitive to the value of a_s .

After computing cross sections as a function of collision energy, they determine temperature-dependent cross sections by computing a thermal average over collision events, weighted by the collision energy. Using this type of thermal averaging to account for a temperature-dependent cross section is supported by Monte Carlo studies.⁹ To make a fit to the data, they compute χ^2 while floating both a_t and a multiplicative factor ϵ which scales simultaneously the computed $\sigma_s(T)$ and $\sigma_p(T)$. This factor is required to accommodate a $\pm 50\%$ systematic uncertainty in the experimental determination of absolute cross sections (primarily from N). Their global best fit occurs for $a_t =$ $157 \pm 20a_0$ and $\epsilon = 1.6$, with a reduced χ^2 of 3.8; the corresponding cross sections are plotted as lines in Fig. 2. The uncertainty in a_t reflects a doubling of the fit χ^2 and includes a $\sim 2a_0$ uncertainty arising from varying C_6 over its range $3600 < C_6 < 4000$ a.u. [96].¹⁰ Their nominal potential gives a p-wave shape resonance at $\sim 280\mu$ K in collision energy, with an asymmetric lineshape whose FWHM is $\sim 400\mu$ K.

The relatively small uncertainty on the value of a_t is attributable to the fact that they can simultaneously fit s-wave and p-wave collision data having little relative uncertainty. The value of a_t for ⁴⁰K determined here does not agree well with reference [41], highlighting the importance of low temperature data in determining accurate potentials. Our measurement is however in good agreement with the value $a_t = 194^{+172}_{-42}$ obtained from an analysis of photoassociation spectroscopy of ³⁹K [95]. This agreement between two fundamentally different experiments is very encouraging, and suggests that the potassium scattering lengths are now fairly well determined. ¹¹ Using their new value, they have tabulated in table 3.1 the resulting triplet scattering lengths for collisions between the different potassium isotopes.

⁹ For more details see M. Holland.

¹⁰ This was the range for C_6 in 1999, when this measurement was made.

¹¹ The result of [95] is somewhat sensitive to the value of C_6 , and is in better agreement with the present result for $C_6 = 4000$ a.u.

Isotopes	Nominal a_t	Range
40 + 40	157	$136 < a_t < 176$
39 + 39	-44	$-80 < a_t < -28$
41 + 41	57	$49 < a_t < 62$
39 + 40	3600	$a_t > 500 \text{ or } a_t < -900$
39 + 41	164	$140 < a_t < 185$
40 + 41	93	$83 < a_t < 99$

Table 3.1: Triplet scattering lengths a_t in Bohr radii for collisions between potassium isotopes.

3.5 Inelastic Collisions

We were first suspicious that 40 K might have abnormally low inelastic collision rates when we discovered some $m_f = 5/2$ atoms present after a several minute evaporation sequence. One would expect that these atoms would undergo spin exchange collisions and disappear, ultimately populating $m_f = 9/2$ and untrapped states. Further, even at peak densities exceeding 10^{13} cm⁻³ we were unable to observe any rapid, non-exponential loss when measuring the magnetic trap lifetime. We used changes in the spin composition of the gas to detect spin exchange collisions and measure the spin exchange collision rate constant. This measurement is consistent with an upper limit derived from the stability of a mixture of four spin states at lower temperature.

The number of $m_f = 9/2$ atoms produced by a mixture of $m_f = 5/2$ and $m_f = 7/2$ atoms was measured in order to determine the spin exchange rate constant K. This was a quick measurement, with few experimental checks. However, it is useful to have an order of magnitude estimate for an upper limit on K. Assuming that both $m_f = 9/2$ and $m_f = 3/2$ atoms are created and therefore present, all allowed s-wave spin exchange collisions are detailed in table 3.2. These collisions involve a change in potential energy $Q = E_f - E_i$ caused by the second order Zeeman effect. The Q value for each collision is also listed in table 3.2 for a 5 gauss field. The $m_f = 1/2$ atoms are not trapped and

collision	<i>Q</i> (nK)
$m_f = 9/2 + m_f = 3/2 \rightarrow m_f = 7/2 + m_f = 5/2$	-369
$7/2 + 5/2 \rightarrow 9/2 + 3/2$	369
$7/2 + 3/2 \rightarrow 9/2 + 1/2$	562
$5/2 + 3/2 \rightarrow 7/2 + 1/2$	366
$5/2 + 3/2 \rightarrow 9/2 + -1/2$	1100

leave the gas once they are produced.

Table 3.2: Possible spin exchange reactions for the measurement of K.

After evaporating in the "evap" trap and adiabatically ramping to the "tight" trap, all $m_f = 9/2$ atoms were removed using a swept microwave field. The gas was then initially was composed of $1.1 \times 10^6 m_f = 5/2$ and $0.9 \times 10^6 m_f = 7/2$ atoms at 20μ K. The gas was held in the "tight" trap for 200-300 seconds, and then all the atoms *except* for $m_f = 9/2$ were removed by again using a microwave sweep. The number of $m_f = 9/2$ atoms produced during the "tight" trap hold (which is the number of atoms left after the second microwave sweep) was then measured using absorption imaging after release from the magnetic trap.

A very simple model is used to determine the spin exchange rate. We assume that the spin exchange rate constant is the same for all of the collisions listed in table 3.2. There is very little heating of the gas during the measurement, and therefore we assume that the radial cloud size $x_{rms} = \left(\frac{k_b T}{m\omega_r^2}\right)^{1/2}$ in the trap is constant and the same for the spin states that we consider in this measurement. This is a fairly bad assumption, as the magnetic moments of the $m_f = 9/2$ and $m_f = 3/2$ atoms are quite different, for instance. On the other hand, this is not meant to be a very exact measurement. The total rate of spin exchange collisions is then Kn_{12} , where $n_{12} = \frac{1}{2\sqrt{2}} \frac{N_1 N_2}{(2\pi)^{3/2} x_{rms}^3}$ is the overlap integral of the density distributions for states 1 and 2. The rate Kn_{12} gives the total number of spin exchange collisions in the gas per second. The coupled differential equations describing the change in number of each state, including a term for loss due to background collisions with a 300 second magnetic trap lifetime, are numerically solved:

$$\frac{dN_{9/2}}{dt} = \frac{K}{A} \left(-N_{9/2}N_{3/2} + N_{7/2}N_{5/2} + N_{7/2}N_{3/2} + N_{5/2}N_{3/2} \right) - \frac{N_{9/2}}{300 \ s} \tag{3.18}$$

$$\frac{dN_{7/2}}{dt} = \frac{K}{A} \left(-N_{7/2}N_{5/2} - N_{7/2}N_{3/2} + N_{9/2}N_{3/2} + N_{5/2}N_{3/2} \right) - \frac{N_{7/2}}{300 \ s} \tag{3.19}$$

$$\frac{dN_{5/2}}{dt} = \frac{K}{A} \left(-N_{7/2}N_{5/2} - 2N_{5/2}N_{3/2} + N_{9/2}N_{3/2} \right) - \frac{N_{5/2}}{300 \ s} \tag{3.20}$$

$$\frac{dN_{3/2}}{dt} = \frac{K}{A} \left(-N_{9/2}N_{3/2} - 2N_{5/2}N_{3/2} - N_{7/2}N_{3/2} + N_{7/2}N_{5/2} \right) - \frac{N_{3/2}}{300 \ s} \tag{3.21}$$

where $A = 2\sqrt{2}(2\pi)^{3/2}x_{rms}^3$. We ignore threshold effects since the average collision energy is high compared to the Q values for all of these reactions. The value of K is varied to match the number of $m_f = 9/2$ atoms produced during the experiment.

We determine that $K = 1.0(3) \times 10^{-14} \text{ cm}^3/\text{s}$, where the uncertainty represents the scatter in the data points. This should really be considered an upper limit, since other, unknown processes may be contributing to the production of $m_f = 9/2$ atoms. By setting the rate constant to zero for individual reactions in the calculation explained above, we find that this experiment is most sensitive to the reaction $5/2 + 7/2 \rightarrow 9/2 +$ 3/2. This is not surprising considering that the initial conditions are a gas consisting of only $m_f = 5/2$ and $m_f = 3/2$ atoms and that relatively few $m_f = 9/2$ atoms are produced. A multi-channel calculation courtesy of John Bohn for the inverse reaction at 5 gauss and different collision energies is shown in figure 3.4. The measured value of K agrees well with the calculation, which uses the currently best known values of C_6 and a_t . The dependence of K on a_t is shown in figure 3.5. Apparently there is some interference effect in the scattering, much like ⁸⁷Rb where the fact that $a_t \approx a_s$ produces very low spin exchange collision rates [97].



Figure 3.4: Spin exchange rate constant for $m_f = 5/2 + m_f = 7/2 \rightarrow 9/2 + 3/2$ at different collision energies and 5 gauss.



Figure 3.5: Spin exchange rate constant (calculated at 5 gauss and 20 μ K to match the experimental conditions) for $5/2 + 7/2 \rightarrow 9/2 + 3/2$ for different values of a_t . The vertical lines indicate the current range on a_t . The horizontal line is the value of our measurement.

In a separate measurement, we use the spin mixture stability at lower temperature to put an upper limit on the spin exchange rate. For this measurement, after cooling to 200 nK the gas consisted of $2.1 \times 10^5 m_f = 9/2$ atoms, $1.8 \times 10^5 m_f = 7/2$ atoms, $1.2 \times 10^5 m_f = 7/2$ atoms, and $0.15 \times 10^5 m_f = 3/2$ atoms (a picture of this gas can be found in figure 2.58). The gas is held in the "tight" trap for 50 seconds, and then released. The number in the $m_f = 9/2$ and $m_f = 7/2$ components are measured using the AG coil and absorption imaging. After 50 seconds, the gas has heated to 1 μ K (see chapter 4 for a discussion of the heating), and we use this temperature to calculate the constant A. The equations equivalent to 3.7-3.10 are solved assuming that only the first collision in table 3.2 is allowed and that the others are threshold suppressed. An upper limit of 2×10^{-14} cm³/sec for K is determined by varying K to match the time dependence of the ratio $\frac{N_{9/2}}{N_{9/2}+N_{7/2}}$, which is measured to be fixed to within 4% over 50 seconds. This upper limit is consistent with the calculated value shown in figure 3.4.

3.6 Appendix to Chapter 3

3.6.1 Time dependence of the aspect ratio

We develop a simple cross-dimensional rethermalization model here that was used to analyze the data in this chapter. This model is valid under certain assumptions that will be relaxed in chapter 6. In chapter 6 we will also explain how this model can be derived from the first principles of kinetic theory.

The average energy per particle in one dimension is defined as:

$$\langle E_x \rangle = \langle \frac{1}{2} m \omega_x^2 x^2 \rangle + \langle \frac{p_x^2}{2m} \rangle \tag{3.22}$$

where a statistical average is done over the gas. A similar equation holds for z. Since the difference in magnetic moments between the $m_f = 9/2$ and $m_f = 7/2$ atoms is small, we treat both components as one gas with a single harmonic oscillator frequency. We assume that $\langle E \rangle$ relaxes exponentially in time in each direction:

$$E_z(t) = E_{zf} + (E_{zi} - E_{zf}) e^{-\frac{t}{\tau}}$$
(3.23)

$$E_x(t) = E_{xf} + (E_{xi} - E_{xf}) e^{-\frac{t}{\tau}}$$
(3.24)

where the averaging symbol is dropped for convenience. The subscripts z and x are used to denote the axial and and one of the radial directions, and f and i denote the final $(t \to \infty)$ and initial $(t \to 0)$ values. We assume that $\langle E_x \rangle = \langle E_y \rangle$ because of the cylindrical symmetry of the magnetic trap.

The physical picture behind equations 3.23 and 3.24 is that collisions are redistributing energy between the x and z directions in the gas. For that reason, both directions relax in time with the same exponential time constant.

The aspect ratio of an expanded cloud, which is the ratio of the rms sizes $\frac{z_{rms}}{x_{rms}}$, is directly connected to the ratio of the energy E_z/E_x . The expanded aspect ratio is given by:

$$\frac{z_{rms}}{x_{rms}} = \frac{z_{0rms}}{x_{0rms}} \frac{\sqrt{1 + (\omega_z t_{exp})^2}}{\sqrt{1 + (\omega_x t_{exp})^2}}$$
(3.25)

where t_{exp} is the expansion time, and the "0" indicates the size in the trap before the expansion. The appendix to chapter 6 will explain how an anisotropic density distribution can be derived that is properly normalized for N. In the classical limit, the anisotropic density distribution is

$$n(r,z) = \frac{N}{(2\pi)^{3/2} x 0_{rms}^2 z 0_{rms}} e^{-\frac{r^2}{2x 0_{rms}^2}} e^{-\frac{z^2}{2z 0_{rms}^2}}$$
(3.26)

where $r^2 = x^2 + y^2$. The key assumption here is that each direction maintains a well defined Maxwell-Boltzmann energy distribution, and that the axial and radial directions are separable. The average potential energy $(\langle u_x \rangle)$ in the x direction, for example, is then connected to the r.m.s. size in the trap through:

$$\langle u_x \rangle = \frac{\int d^3 \vec{r} \, n(r,z) \frac{1}{2} m \omega_x^2 x^2}{\int d^3 \vec{r} \, n(r,z)} = \frac{1}{2} m \omega_x^2 x 0_{rms}^2$$
(3.27)

where a similar equation holds for z. Finally, using the equipartition theorem, $\langle E_x \rangle = m\omega_x^2 x 0_{rms}^2$ and $\langle E_z \rangle = m\omega_z^2 z 0_{rms}^2$.

Combined with equations 3.23 and 3.24, energy conservation

$$E_{zf} + 2E_{xf} = E_{zi} + 2E_{xi} \tag{3.28}$$

and the condition for cross-dimensional equilibrium

$$E_f = E_{xf} = E_{zf} \tag{3.29}$$

are used to solve for the time dependence of E_z/E_y . After some algebra and substituting in the relation between the energy and the r.m.s. sizes, we find that:

$$\frac{z_{rms}}{x_{rms}} = \sqrt{\gamma \frac{1 + (s-1)e^{-\frac{t}{\tau}}}{1 - \frac{1}{2}(s-1)e^{-\frac{t}{\tau}}}}$$
(3.30)

where $s = \frac{E_{zi}}{E_f}$ and $\gamma = \frac{1 + (\omega_z t_{exp})^2}{1 + (\omega_x t_{exp})^2} \frac{\omega_x^2}{\omega_z^2}$. We leave γ and s as free parameters in the fit to the time dependence of the aspect ratio.

Chapter 4

EVAPORATIVE COOLING

4.1 Overview

After demonstrating that we could maintain rethermalizing collisions in the gas at low temperature, we needed to devise an evaporation scheme that could cool to quantum degeneracy. We adapted the evaporative cooling techniques [12,98–100] used to produce atomic Bose-Einstein condensation to the unique requirements of fermions. Cooling fermionic atoms is more difficult than bosons. Working with multiple components increases the complexity of the evaporation procedure. Further, efficient evaporative cooling relies on a density increase, even as atoms are removed, that arises from the harmonic trap compressing the cloud as the temperature is lowered. The density increase results in a increasing collision rate for decreasing temperature, producing "runaway" evaporation for bosonic and classical gases. Fermions, unfortunately, have exactly the opposite behavior in the degenerate regime. Once FD statistics start to dominate the thermodynamics, the Fermi pressure starts to "freeze" the cloud size and the density stops increasing as quickly. Even worse, the collision rate is shut off by Pauli blocking, eventually falling to zero at T = 0. So, as the Fermi gas is cooled into the degenerate regime the *collision rate must decrease*.

In practice, we do not think that the evaporation performance is currently limited by the collision rate. Ultimately, the limiting collision rate behavior described above must become a problem, but the experiments detailed in Chapter 6 will show that the rethermalization rate is high even at our lowest values of T/T_F . Because we are trying to cool to as low T/T_F as possible, technical limitations to the evaporation performance, such as heating and magnetic trap stability, are dominant. This is in contrast to experiments with BEC, where reaching 0.5 T_c is sufficient to ensure negligible thermal fraction. We believe that technical limitations are the current barrier to our ability to reach lower T/T_F .

This chapter will discuss the details of the cooling procedure, as well as explain the most important technical requirements. Our current cooling scheme represents two years of work, progressing from only being able to cool a spin polarized gas to $T/T_F = 0.5$ to producing an interacting, equal mixture of $m_f = 9/2$ and $m_f = 7/2$ components with the $m_f = 9/2$ component at $T/T_F \sim 0.2$. A factor of two in temperature may not seem large, but actually represents a huge change in the behavior of the gas. The incentive to continually improve the evaporation came from a time-dependent model constructed by Murray Holland [54] that found no fundamental limit to evaporative cooling of fermions. This chapter will develop a time-independent model that was useful for ferreting out and fixing technical limits.

4.1.1 Chapter Content

This is a long and detailed chapter, so I will give a quick list of the content here. General cooling considerations relevant to working with two components and an overview of our evaporation scheme are explained in section 4.2. Details on the first, single-frequency stage of the evaporation are found in section 4.2. Details of the second, two-frequency stage of the evaporation and results circa 1999 are found in 4.3.3.1. Improvements to the two-frequency stage and results circa 2000 are found in 4.3.3.2. The explanation of a toy-model of evaporation that was useful for learning how to improve the evaporation performance is found in 4.4. Characterization of the heating we observe for a magnetically trapped gas is detailed in 4.5. Finally, experimental details for the evaporation are listed in the appendix to the chapter.

4.2 General Cooling Considerations

There are a few general requirements that must be met for evaporative cooling to work well, independent of working with multiple components. The ratio of the rate of elastic, or "good", collisions to the rate of ("bad") collisions that lead to uncontrolled loss should be high, typically higher than 100. For our system, this is a constraint on the rethermalization rate compared to the inverse of the trap lifetime ($\sim 3 \text{ mHz}$). Also, simultaneous evaporation for multiple components requires not only energy selectivity in atom removal but also component selectivity. Finally, the cooling rate must always exceed the heating rate.

The overall rethermalization rate in the gas depends not only on the total number and temperature but the spin mixture and the component spatial overlap as well. The total collision rate (collisions/sec) in the gas is $\Gamma = n_{12}\sigma v$, where σ is the collision cross-section, n_{12} is the density overlap integral, and v is the mean relative speed. For information on the calculation of n_{12} and v, see Chapter 6. Neglecting relative sag,

$$n_{12} = \frac{N_{9/2} N_{7/2} \lambda}{(2\pi)^{3/2} \left(\frac{k_b T}{m\omega_{9r}^2}\right)^{3/2}} \frac{\xi^3}{(1+\xi^2)^{3/2}}$$
(4.1)

for classical clouds with different magnetic moments. In this equation, $\lambda = \omega_z/\omega_r$ is the same for both components, and $\xi = \omega_{7/2}/\omega_{9/2} = \sqrt{7/9}$. The mean relative speed for classical gases is

$$v = \frac{4}{\sqrt{\pi}} \sqrt{\frac{k_b T}{m}} \tag{4.2}$$

and is number independent. If we rewrite the number of atoms as $N_{9/2} = fN$ and $N_{7/2} = (1-f)N$ where $N = N_{9/2} + N_{7/2}$ and f is the fraction of atoms in the $m_f = 9/2$ component, then the collision rate is proportional to $f(1-f)N^2$. The collision rate is highest for an equal mixture, diving to zero for a spin polarized gas. In order to maintain a high collision rate, the spin mixture should be kept close to an equal mixture

of components. At a 70/30% mixture the collision rate is reduced by 20% from the maximum value, and at a 90/10% mixture the collision rate is down by a factor of 2.8.

The centers of each component are displaced from each other by gravitational sag. The displacement for one component from the minimum of the magnetic field is g/ω_r^2 , where g is gravitational acceleration. The displacement between the $m_f = 9/2$ and $m_f = 7/2$ components is then:

$$\Delta y = \frac{g}{\omega_{9r}^2} \left(\frac{1}{\xi^2} - 1\right) \tag{4.3}$$

with the same definition of ξ . For classical distributions, the value of n and therefore Γ is reduced by

$$\frac{\int d^3 \vec{r} \, n_{9/2}(x, y, z) \, n_{7/2}(x, y - \Delta y, z)}{\int d^3 \vec{r} \, n_{9/2}(x, y, z) \, n_{7/2}(x, y, z)} = e^{-\frac{m\Delta y^2 \xi^2 \omega_{9r}^2}{2k_b T(1+\xi^2)}} \tag{4.4}$$

which limits the choice of trap parameters. As we cool to low T, the spatial overlap must be preserved by choosing a trap with sufficient radial confinement. For a 100 nK *classical* gas in the "tight" trap, the collision rate is reduced by only 10%.¹

We use magnetic field dependent transitions between hyperfine ground states to remove atoms in order to have both energy and component selectivity [42]. The component selectivity is important to maintaining a 50/50% mixture as well as removing energy in a balanced way. The transitions we use are shown in figure 4.1. The transitions $|f = 9/2, m_f = 9/2\rangle \rightarrow |7/2, 7/2\rangle$ and $|9/2, 7/2\rangle \rightarrow |7/2, 7/2\rangle$ are driven together with a single microwave frequency in the first ("single-frequency") stage of evaporation, while $|9/2, 9/2\rangle \rightarrow |7/2, 7/2\rangle$ and $|9/2, 7/2\rangle \rightarrow |7/2, 7/2\rangle$ are used separately in the second stage ("two-frequency"). Atoms in the $m_f = 5/2$ and $m_f = 3/2$ states are removed when necessary using the transitions $|9/2, 5/2\rangle \rightarrow |7/2, 3/2\rangle$ and $|9/2, 3/2\rangle \rightarrow |7/2, 5/2\rangle$, which are degenerate in frequency at low field. In practice, we sometimes use other transitions when convenient.

¹ Note that this problem is reduced for a degenerate gas, since the rms size of the gas is larger than the classical expectation.



Figure 4.1: Microwave transitions used for evaporation. Note that there are many Zeeman levels that are not shown in this figure.

The frequency of these transitions can be calculated using the Breit-Rabi formula [101], which includes the second order Zeeman shift. We start by calculating the energy of each Zeeman level relative to the 4S ground state. We define the parameters

$$x(B) = \frac{(g_J + g_I r)\mu_b B}{h\nu_{hfs}} \tag{4.5}$$

$$\zeta(B, m_f) = 1 + \frac{4m_f x(B)}{2I+1} + x(B)^2$$
(4.6)

$$Z_u(B, m_f) = \sqrt{\zeta(B, m_F)} \tag{4.7}$$

$$Z_{l}(B, m_{f}) = \begin{cases} \sqrt{\zeta(B, m_{f})} & m_{f} > -7/2 \\ \sqrt{\zeta(B, m_{f})} & m_{f} < -7/2 \& x(B) < 1 \\ -\sqrt{\zeta(B, m_{f})} & m_{f} < -7/2 \& x(B) > 1 \end{cases}$$
(4.8)

where u and l refer to the upper (f = 7/2) and lower (f = 9/2) states, $g_J = 2$, I = 4, $g_I = -0.363$, $v_{hfs} = -1286$ MHz, μ_b is the Bohr magneton, and $r = \frac{1}{1836}$ is the ratio of the Bohr to nuclear magneton. The energy of any level is then:

$$E_u(m_f, B) = -\frac{h\nu_{hfs}}{2(2I+1)} - g_i\mu_n Bm_f - \frac{h\nu_{hfs}}{2}Z_u(B, m_f)$$

$$E_l(m_f, B) = -\frac{h\nu_{hfs}}{2(2I+1)} - g_i\mu_n Bm_f + \frac{h\nu_{hfs}}{2}Z_u(B, m_f)$$
(4.9)

where μ_n is the nuclear magneton. The frequency for any microwave transition is equal to $\frac{1}{h} \left[E_u(m'_f, B) - E_l(m_f, B) \right]$. Equations 4.9 can also be used to determine the field dependent magnetic moment through $\mu(m_f, B) = -\frac{\partial E(m_f, B)}{\partial B}$:

$$\mu_l(B, m_f) = g_I \mu_n m_f - \frac{(g_J + g_I r) \mu_b}{4\sqrt{\zeta(B, m_f)}} \left[\frac{4m_f}{2I + 1} + 2x(B) \right]$$
(4.10)

$$\mu_u(B, m_f) = g_I \mu_n m_f + \frac{(g_J + g_I r)\mu_b}{4\sqrt{\zeta(B, m_f)}} \left[\frac{4m_f}{2I + 1} + 2x(B)\right]$$
(4.11)

for the upper or lower hyperfine ground states.

A measurement of the frequency of these transitions in the tight trap is shown in figure 4.2. For this measurement, the gas was cooled to 1.2 μ K. The number of atoms was measured after a slow (5.9 seconds total time), linear frequency sweep from high frequency to low frequency was applied. The transition frequencies at B_0 are indicated by a "step" in atom number — as the frequency is swept over the transition at the minimum of the trap the atoms are all removed. In this data, the transitions $|9/2,7/2\rangle \rightarrow |7/2,5/2\rangle$ and $|9/2,5/2\rangle \rightarrow |7/2,7/2\rangle$ are degenerate and appear at 1276.04 MHz. The $|9/2,7/2\rangle \rightarrow |7/2,7/2\rangle$ transition is not seen because the $m_f = 7/2$ atoms have already been removed by that point in the sweep. The $|9/2,7/2\rangle \rightarrow |7/2,5/2\rangle$ transition appears at 1272.75 MHz. The location of the transition frequencies and the spacing between them is consistent with $B_0 = 5.336$ gauss.

There are three regimes for the evaporation. The kinetic energy scale k_bT sets the width of the cloud in the magnetic potential and hence the spread of microwave frequencies $\delta\nu \sim k_bT/h$ that will remove atoms from the trap. The shape of the gas in frequency space is determined by calculating the number of atoms that would remain in each component versus the applied microwave frequency (see figure 4.3). The calculation assumes that the microwave field removes all atoms above the cut energy, E_c . The cut energy as a function of the magnetic field experienced by an atom is $E_c = \frac{2}{9}m_f\mu_b(B - B_0)$, where $-\frac{2}{9}m_f\mu_b$ is the magnetic moment to first order. Ignoring second order



Figure 4.2: Measurement of location of microwave transition frequencies in the "tight" trap. Atoms are removed with a slow frequency sweep from high (starting at 1279 MHz) to low frequency.

Zeeman shifts,

$$E_{c} = \frac{2}{9} m_{f} \mu_{b} \left[\frac{\nu - 1286 \text{ MHz}}{\Lambda_{m_{f}, m'_{f}}} - B_{0} \right]$$
(4.12)

where the transition is between $|9/2, m_f\rangle$ and $|7/2, m'_f\rangle$ with applied microwave frequency ν . In equation 4.12, $\Lambda_{m_f,m'_f} = -\frac{2}{9}\frac{\mu_b}{h}(m_f + m'_f)$ is the shift in the transition frequency with field to first order. The number of atoms left after removing all the atoms above E_c is

$$N_{left} = \int_0^{E_c} dE \ g(E)\mathcal{F}(E) \tag{4.13}$$

where $\mathcal{F}(E)$ is the statistical distribution function and g(E) is the density of states (see

chapter 5). For a classical gas,

$$N_{left} = N \times \begin{cases} 1 - e^{-\eta} \left(1 + \eta + \frac{\eta^2}{2} \right) & \eta > 0 \\ 1 & \eta < 0 \end{cases}$$
(4.14)

where $\eta = \frac{E_c}{k_b T}$. The case $\eta < 0$, calculated strictly from the equation for E_c , corresponds to an applied microwave frequency that is too high to remove any atoms from the trap. This corresponds to a microwave frequency that would be resonant at a field lower than the bias field.

The three regimes and corresponding microwave spectra described by equation 4.14 are shown in figure 4.3. The trap bias field sets the Zeeman shift $\Delta \nu = \nu_{9/2} - \nu_{7/2}$ between the two transitions at the trap center. At relatively high T (case 1) where $\delta \nu >> \Delta \nu$, a single frequency removes atoms nearly equally from both spin states. At low T (case 3) where $\delta \nu \ll \Delta \nu$, the microwave lines are distinct and two microwave frequencies are needed to cool both components in parallel. In the intermediate case (case 2) where $\delta \nu \sim \Delta \nu$, the application of any relevant microwave frequency will remove unequal numbers of atoms from each species.

We call our evaporation scheme "simultaneous cooling" [53] because it is designed to always remove energy and atoms from each component, and therefore cool each component at the same time. The evaporation sequence starts at 1 mK in case 1 in the "evap" trap (1.3 gauss B_0). A swept single frequency microwave field is applied (with frequency indicated by ν_{evap} in figure 4.3) until we reach case 2. The fact that evaporation in case 2 changes the spin mixture is used to produce a 50/50% mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms by continuing to apply a single frequency. We often refer to the evaporation from case 1 through case 2 as the "evap1" stage.

After a 50/50 mixture of atoms has been prepared, the trap is then adiabatically ramped to the "tight" trap (5.3 gauss B_0) to move to case 3. Cooling is continued with the application of a two-frequency field ("evap2"), with the spacing between the frequencies adjusted to preserve an equal mixture of components. Not only is a high



Figure 4.3: Microwave lineshapes and evaporation regimes. We plot the fraction of atoms that would remain in each component versus applied microwave frequency, with the frequency appropriate for evaporation ν_{evap} indicated by a dotted line (note that lower microwave frequencies remove higher energy atoms). The calculation assumed a classical gas in which the microwave field removes all atoms above a particular energy. The high T (case 1, Fig. 2), low T (case 3), and intermediate regimes (case 2) are distinguished by the spread of frequencies $\delta\nu$ resonant with atoms in the trap compared with the Zeeman shift $\Delta\nu$ at the trap center.

collision rate sustained, but energy is removed in a balanced way. Sympathetic cooling by removing energy and atoms from only one component would fail once the number imbalance became too large. In that case, the heat capacity carried by the other component would prevent any change in temperature. At the end of the two-frequency evaporation, atoms in either spin state may be removed to produce a spin-polarized gas. The removal sweep is designed to sympathetically cool the remaining component.

4.3 Evaporation Sequence and Performance

4.3.1 Single Frequency Stage

The first stage of evaporation begins with roughly 4×10^8 atoms at 1 mK in the evap trap. The collision rate per atom is about 0.3 Hz, implying that the ratio of "good" to "bad" collisions is about 100 and that evaporation should run-away. The evaporation trajectory for the first stage is shown in figure 4.4. The data was analyzed assuming

a 50/50 mixture of spin-states, and T/T_F is plotted for the $m_f = 9/2$ component. In truth, the spin mixture starts at 70/30 ($m_f = 9/2/m_f = 7/2$) and the final cut is chosen to produce a 50/50 mixture. This data was taken without the AG coil, and the spin composition was not measured carefully throughout the entire trajectory.

The data is shown using the usual release sequence, meaning that the trap is adiabatically ramped to the "tight" trap before release. Therefore, the temperature and peak optical depth in the trap do not represent the values in the "evap" trap. For an adiabatic ramp, T/T_F is preserved. Assuming that the axial frequency is fixed, the temperature scales with $\omega_r^{2/3}$. See the appendix to this chapter for information on the ramp.

As shown in figure 4.4, the evaporation does run-away; the peak OD in the trap, which is proportional to the collision rate, increases by nearly an order of magnitude. Information on the efficiency of evaporation is contained in the plot of T/T_F versus N. The steepness of the slope is directly related to the evaporation efficiency — the most efficient evaporation accomplishes a large change in T/T_F with a small change in N. We lower T/T_F by two orders of magnitude with only a factor of 50 decrease in number.

The frequency sweep used for evaporation is broken up into four stages. Each stage exponentially sweeps the frequency in time. The speed of the sweep, the microwave power, and other parameters are optimized to produce the highest number at fixed temperature. Details on these sweeps and the optimization procedure can be found in the appendix to this chapter.

Data taken at the end of the single frequency stage (using the AG coil) are shown in figure 4.5. The cooling clearly progresses into case 2, where more $m_f = 9/2$ atoms than $m_f = 7/2$ atoms are being removed. In fact, the $m_f = 7/2$ component is being very efficiently sympathetically cooled. While the evaporative process has ceased to runaway for the $m_f = 9/2$ component, run-away is evident for the $m_f = 7/2$ component as seen in the peak OD shown at a 10 ms expansion time. A final cut of 1281.8 MHz



Figure 4.4: Single frequency evaporation trajectory. This data was taken without the AG coil, and the analysis assumes a 50/50 mixture. The total number $N = N_{9/2} + N_{7/2}$ is plotted, while T/T_F is plotted for just the $m_f = 9/2$ component. The cooling is performed in the "evap" trap, but the measurements shown in this figure are made after an adiabatic ramp to the "tight" trap.

would be chosen based on this data in order to produce a 50/50 mixture. It is also at this cut where the $m_f = 9/2$ peak OD starts to decrease.

The data in figure 4.5 also explains the advantage to using a partially spin polarized gas. If the end of the single frequency cooling is set by the constraint of producing a 50/50 mixture, a higher degree of initial spin polarization allows the evaporation to end at lower T/T_F . Because this cooling takes place in the high frequency "evap" trap, it can be more efficient than in the "tight" trap and occur much faster.

The microwave power has an effect on the behavior at the end of the single



Figure 4.5: End of single frequency evaporation trajectory. This data was taken using the AG coil, and is plotted independently for each component. Note that the evaporation occurs in the "evap" trap, but these measurements are made after an adiabatic ramp to the "tight" trap. This data was taken at a 10 ms expansion time, using the AG coil to separate the two spin components.

frequency evaporation (see figure 4.6). This data was taken by varying the final cut to always produce a 50/50% mixture. Higher power allows a deeper cut and lower T/T_F for this stage of the evaporation. This behavior is consistent with the timeindependent model described in section 4.4, where the evaporation actually becomes more efficient with higher power. The data, in conjunction with the model, suggests that the evaporation performance is power limited.



Figure 4.6: Effect of microwave power on the end of the single frequency evaporation. This data was taken using the AG coil, and the final cut was varied to always produce a 50/50 mixture. Note that this data was taken after an adiabatic ramp to the "tight" trap.

4.3.2 $m_f = 5/2$ and $m_f = 3/2$ Cleanout

After the first stage of evaporation, the trap is ramped to the "tight" trap. This is done to move to case 3 cooling by increasing the bias field, adiabatically cooling the gas, and therefore increasing $\Delta\nu$ compared to $\delta\nu$. In the "tight" trap, the different spin components can be separately addressed using multiple microwave frequencies at the temperature of the gas after the "evap1" stage.

Before the two-frequency stage, $m_f = 5/2$ and $m_f = 3/2$ atoms that are present are removed. This is done with a fast sweep on the $|9/2, 5/2\rangle \rightarrow |7/2, 3/2\rangle$ transition, which is degenerate with the $|9/2, 3/2\rangle \rightarrow |7/2, 5/2\rangle$ transition. Because there are relatively few atoms in these states, the sweep cannot cool the gas much and is as fast as possible. The temperature of the gas after the sweep is shown in figure 4.7 versus the exponential sweep rate. For more details on the sweep parameters, see the appendix to this chapter.



Figure 4.7: Fast $m_f = 5/2$ and $m_f = 3/2$ removal sweep timing. The sweep must be done with a fast exponential sweep rate α since there are not enough atoms to efficiently cool and defeat the heating.

4.3.3 Two-frequency Stage

The two-frequency stage is also broken up into multiple exponential sweeps, each of which is optimized separately (see the appendix to this chapter). There is an additional parameter, Δf , compared to the "evap1" sweeps which controls the spacing between the frequencies. Δf is adjusted for each sweep to produce a 50/50 mixture at the end of the sweep (each sweep is 0.05-0.1 MHz long).

The constraint that sets Δf in theory is that the cut energy should be identical for each component when measured in units of k_bT . If we set $\nu_{7/2} = \nu_{9/2} + \Delta f$, and solve for Δf under this constraint we find that

$$\Delta f = \frac{\Lambda_{7/2,7/2}}{\mu_{7/2}} \left[\nu_{9/2} \left(\frac{\mu_{9/2}}{\Lambda_{9/2,7/2}} - \frac{\mu_{7/2}}{\Lambda_{7/2,7/2}} \right) - \nu_{0,9/2,7/2} \frac{\mu_{9/2}}{\Lambda_{9/2,7/2}} + \nu_{0,7/2,7/2} \frac{\mu_{7/2}}{\Lambda_{7/2,7/2}} \right]$$
(4.15)

where we have introduced new parameter ν_0 , which corresponds to the frequency that is resonant at B_0 . To first order in the Zeeman shift, $\nu_{0,m_f,m'_f} = 1286 \text{ MHz} + \Lambda_{m_f,m'_f} B_0$.

According to the above equation, across a 0.15 MHz evaporation sweep in the tight trap, Δf should only change by ~ 0.020 MHz. This is consistent with a measurement of the microwave lineshapes shown in figure 4.8. For this measurement, the gas was cooled and then spin polarized by removing all of the $m_f = 9/2$ or $m_f = 7/2$ atoms in order to turn off rethermalization and heating. With a spin polarized gas, the gas cannot rethermalize during the microwave lineshape measurement because s-wave collisions are forbidden and p-wave collisions are heavily suppressed at this temperature. Heating is also heavily suppressed (see section 4.5) for a spin-polarized gas. The temperature of the remaining $m_f = 9/2$ or $m_f = 7/2$ gas was matched to 1.5 μ K within the experimental scatter on T.

A sweep from low to high frequency was applied, and the number of atoms left was measured versus the final frequency. In the plot, the fraction of atoms remaining is plotted versus the difference between the final cut frequency $\nu_{9/2}$ or $\nu_{7/2}$ and $\nu_{0,9/2,7/2}$ or $\nu_{0,7/2,7/2}$. A constant offset of 1.55 MHz, which would correspond to Δf , is removed in the plot. In 4.8 is is clear that a constant Δf within 30 to 40 kHz should work for all of the two-frequency evaporation. Within this spread in frequency, the same fraction of each gas is removed for the same difference between ν and ν_0 .



Figure 4.8: Microwave lineshapes of spin-polarized gases. The gas is cooled and then the $m_f = 9/2$ or $m_f = 7/2$ component is completely removed. The fraction of atoms left after a sweep from low to high frequency is plotted versus the final cut from the resonant frequency ν_0 for the relevant transition at B_0 .

4.3.3.1 Original Two-Frequency Evaporation Results

In 1999 we were able to first cool below $T/T_F = 1$ [53]. Data from a cooling trajectory at that time is shown in figure 4.9 as a plot of T/T_F versus N. This data was taken with a $m_f = 7/2$ removal sweep after the two-frequency stage. Working with spin polarized clouds not only simplified interpretation of the images, but the removal sweep significantly cooled the gas (see section 4.3.4). The sweep occurs on a time scale that is long compared to the collision rate so that the component that remains is in thermal equilibrium. As a check, we have taken data varying the time for the removal sweep by a factor of 10 and seen no change in the polarized gas thermodynamics. The thermometry techniques that we use to determine T/T_F will be discussed in the next chapter.



Figure 4.9: First demonstration of cooling below T_F using two-frequency evaporation. The data is taken with a $m_f = 7/2$ removal sweep, so that the data is shown for the resulting spin polarized gas.

A rather startling feature [53] appears in the evaporation trajectory once the gas is cooled to $T/T_F=0.5$. The forced evaporation becomes grossly inefficient in reducing T/T_F and many more atoms are removed to accomplish the same change in temperature. This plunge in cooling efficiency does not coincide with any observed change in the atom loss (not due to evaporation) or heating rate (see 4.5 for an exhaustive heating study). The behavior survives aggressive changes in the evaporation timing, variation in the initial number, and replacement of the two-frequency stage with single-frequency cooling in a lower B_0 trap. The evaporation always failed at $T/T_F \sim 0.5$ in traps covering $\omega_r = 2\pi \times 127$ to 373 Hz, $N = 3.5 \times 10^5$ to 1.2×10^6 , and $T_F = 0.36$ to $1.0 \ \mu$ K. Data comparing continued single-frequency evaporation in a $B_0 = 0.68$ Gauss trap (the "single-frequency" trap) to two-frequency evaporation without a $m_f = 7/2$ removal sweep is shown in figure 4.10.



Figure 4.10: Evaporation data from 1999 taken in different traps. The two-frequency stage was replaced with a single-frequency stage in a low bias field trap. This data was taken with a 50/50 mixture of components (no removal sweep), and is plotted for one of those spin states. In either case, the evaporation becomes inefficient at $T/T_F \sim 1$.

At the time, we attributed this change in the evaporation behavior to the pathological collision rate behavior of fermionic gases discussed earlier. A time-dependent calculation (see figures 4.11 and 4.12) of the evaporation process by Murray Holland [54] revealed no fundamental limit in T/T_F . However, Murray was able to show for an optimized trajectory that the collision rate did start to plummet around $T/T_F = 0.8$. The decreasing collision rate was due partly to Pauli blocking of collisions and partly to the Fermi pressure. This meant that the evaporation speed had to slow down as the temperature decreased, making the evaporation performance very sensitive to technical limits. Incidentally, his optimized trajectory in cut energy as a function of time was very close to the frequency sweeps that we were already using. Murray's optimized cut energy fits almost perfectly to an exponential in time.



Figure 4.11: Evaporation trajectory from the quantum kinetic evaporation model by Murray Holland. The open squares (connected by a line) show the optimized evaporation trajectory as determined by Murray's calculation. The only technical limitation he includes in the simulation is a finite magnetic trap lifetime. The closed circles and open triangles are experimental data taken with spin polarized $m_f = 9/2$ gases, using an evaporative $m_f = 7/2$ clean-out. The closed circles are the data originally published in [53], while the open triangles show the improvement after optimizing the magnetic trap current servos. Optimization of the current servos improved the energy resolution with which atoms are removed from the trap; the comparison in this plots highlights the importance of this resolution to the cooling.



Figure 4.12: Collision rate calculation from the quantum kinetic evaporation simulation by Murray Holland. The dashed line only includes effects due to the Fermi pressure, while the solid line includes suppression due to Pauli blocking as well. The collision rate is plotted as a function of time in the simulation, and the open circles and labels indicate the value of T/T_F at different times.

We began to try to understand and overcome the technical limits to the cooling process. One of the biggest breakthroughs was improving the bias field stability and consequently the evaporation energy resolution. This improvement is detailed in 2.5.5, and involved reducing the bias field noise on the evaporation timescale by a factor of almost 10. A trajectory after this improvement is shown in figure 4.13. Now we were able to cool a spin polarized gas (still using an evaporative $m_f = 7/2$ removal sweep) to $T/T_F = 0.3$. This was a huge gain in our ability to observe the FD statistics.

We revisited the limit on our cooling at this point. Data taken with different initial N is shown in figure 4.14. The initial number was reduced by a factor of two by limiting the number of atoms loaded into the science MOT. Again, the evaporation failure seemed to be related to the degree of degeneracy.



Figure 4.13: Evaporation trajectory after optimization of the magnetic trap current servo's. This trajectory should be compared directly to the data in figure 4.9. Optimizing the current servo's reduced the noise in the bias field on the evaporation timescale by a factor of 10. This data was taken with a spin polarized gas, evaporatively removing the $m_f = 7/2$ component after the "evap2" stage.


Figure 4.14: Evaporation trajectory with different initial number, taken under the same conditions as the data in figure 4.13. The evaporation is not sensitive to changes in the initial number on the order of a factor of two.

4.3.3.2 Evaporation Results circa Late 2000

Our next big breakthrough was discovering the problems with various frequency synthesizers and also increasing the microwave power. We switched from using an HP E4420B and HP 8656B synthesizer to two HP E4420B's.² We also replaced a coil that was impedence matched using a resonance created with surface mount components with the stub tuned coil described in section 2.7. We were now able to cool a 50/50 mixture to $T/T_F \sim 0.35$ (see figure 4.15). The cooling power of the cleanout sweep after two-frequency evaporation was no longer needed in order to cool into the degenerate regime.

Several other improvements allowed us to get even colder. Motivated by the model explained in 4.4, we investigated the sensitivity of the evaporation to the position of the "bottom" frequency of the sweep compared to the resonant frequency at the bottom of the trap, ν_0 . The sweep "bottom" is the frequency that the exponential sweep approaches at long times (see the appendix to this chapter for more details). The data in figure 4.16 for a 50/50 mixture (although only $m_f = 9/2$ data is plotted) shows that this sweep "bottom" frequency should be around 0.15 MHz higher than ν_0 .³ Note that ν_0 is measured as described in chapter 1 by measuring T versus the final evaporation frequency.

The two-frequency evaporation is also sensitive to the bias field noise on a scale that is difficult to measure directly. The current servo's are optimized for a specific **F4** drain-source voltage drop.⁴ Data for the $m_f = 9/2$ component for a deep two-frequency cut is shown in figure 4.17 versus this voltage (altered by changing the power supply voltage). The servo's do their job well DC — this data was taken at a fixed cut and the

 $^{^2}$ This switch essentially increased the microwave duty cycle, and effective power, by a factor of 3-4. With the 8656B, the microwave power had to be attenuated for long periods of time in order to avoid frequency glitches.

³ This means that the sweep "bottom" is past the actual minimum of the magnetic trap.

 $^{^4}$ This voltage drop sets various capacitances associated with the FET and changes the servo performance.



Figure 4.15: Two-frequency evaporation trajectory taken mid 2000 with a spin-mixed gas. This data was taken with *no* removal sweep and using the AG coil. Compared to the data in figure 4.13, the improvement here was from using better frequency synthesizers and higher microwave power.

temperature (which is related to B_0) did not systematically shift with the FET voltage within the normal experimental scatter. The evaporation efficiency, measured by the number left after the sweep, does depend on the voltage with a peak close to 0.9 V.



Figure 4.16: Evaporation performance vs. the sweep "bottom" frequency. This data was taken with 50/50 mixed clouds and the AG coil, although data for only the $m_f = 9/2$ component is plotted. The absolute temperature T was kept constant by changing the final evaporation cut. The changes in T/T_F are then due to changes in the number of atoms left after evaporation. The vertical line indicates the location of ν_0 . For high values of the sweep bottom frequency, the evaporation still works well but needs to be slightly reoptimized (triangle). For sweep bottom frequency values to close to the trap bottom, no amount of reoptimization can help the evaporation performance.



Figure 4.17: Two-frequency evaporation performance vs. the main FET F4 drain-source (DS) voltage drop. The DS voltage was varied by controlling the power supply voltage and affects the current servo performance on a scale that is difficult to measure directly. This data was taken with 50/50 mixed clouds and the AG coil, although data for only the $m_f = 9/2$ component is plotted. The efficiency of the evaporation, indicated by the highest N left after the evaporation, occurs at a voltage drop ~ 0.9 volts. The temperature does not shift systematically with changes in the DS voltage, and the scatter is consistent with the normal scatter in T at a fixed evaporative cut.

A typical example of an evaporation trajectory circa late 2000 is shown in figure 4.18. Without too much work, we are able to cool a spin mixed gas to $T/T_F = 0.25$. With careful manipulation of Δf throughout the evaporation sequence and particularly in the last stage, $T/T_F = 0.2$ can be reached (figure 4.19). The data in figure 4.19 also demonstrates the sensitivity to Δf . The data was taken at fixed $\nu_{9/2}$, so that Δf changes the final cut into the $m_f = 7/2$ component. This is clear in the plot of N versus Δf . Note that the Δf that we find experimentally does not agree with the result of equation 4.15, which predicts 1.67 MHz for the value of $\nu_{9/2}$ used in this data. This is probably due to unequal power in driving the relevant microwave transitions and relative gravitational sag.⁵

The data in both figures 4.18 and 4.19 were taken with a 500 ms microwave pulse to remove $m_f = 5/2$ and $m_f = 3/2$ atoms before the last stage of two-frequency evaporation. We observe that atoms in these states are actually produced during the two-frequency stages. The heat capacity carried by atoms in these states can foil the last stage of cooling.

⁵ Since the microwave transitions we use are driven by different components of the microwave magnetic field, it is unlikely that they are driven equally well. This is compensated for by changing Δf and cutting deeper into one component. In addition, relative sag causes the components to experience slightly different bias fields.



Figure 4.18: Typical two-frequency evaporation trajectory from late 2000. This data was taken with the AG coil, and T/T_F is measured and plotted independently for each component. Without too much work, an equal mixture of components can be cooled to $T/T_F \sim 0.25$. The failure of the cooling in T at deep cuts is particularly bad in this data, but often much colder absolute temperatures can be reached (see the next plot, for example). Also, the spin mixture is this data set was not controlled well because Δf was not set properly — with slightly more effort we can keep the mixture fixed to $50 \pm 2\%$.



Figure 4.19: Sensitivity to Δf , the difference between $\nu_{9/2}$ and $\nu_{7/2}$ during "evap2". These are our coldest temperatures with a spin mixed cloud $(T/T_F = 0.2 \text{ for the } m_f = 9/2 \text{ component})$. The final $\nu_{9/2}$ cut was fixed to 1272.41 MHz for this data. Note that this data was taken at a slightly different bias field than the data in figure 4.18.

4.3.4 $m_f = 7/2$ Removal Sweep

If a spin-polarized, non-interacting gas is desired, one component or the other can be removed evaporatively after the two-frequency evaporation is finished. This is done with a frequency sweep, which is optimized to produce the coldest temperature in the remaining component (see the appendix to this chapter for details). An optimized trajectory for a $m_f = 7/2$ removal sweep is shown in figure 4.20. The data is plotted vs. the final $\nu_{7/2}$ cut of the frequency sweep that is used to remove the $m_f = 7/2$ atoms. The $m_f = 9/2$ component is sympathetically cooled, with the temperature falling by a factor of 2 in this data.⁶ Since there is no loss in atom number, T/T_F is lowered by the same amount.

 $^{^{6}}$ Note that this is the theoretically best (fractional) reduction in temperature that can be accomplished for a classical gas starting with a 50/50 mixture of spin components.



Figure 4.20: Optimized $m_f = 7/2$ removal sweep. The $m_f = 9/2$ gas is sympathetically cooled in this process from its initial ~ 350 nK temperature. See the appendix to this chapter for details on the single frequency sweep that is used to remove the $m_f = 7/2$ atoms.

4.4 Toy Model of Evaporation

4.4.1 Overview

After measuring cross-dimensional relaxation times as short as 60 ms at our lowest temperatures (see chapter 6) with 85% of the $m_f = 7/2$ component removed (from an initially equal mixture of spin states), we guessed that our evaporation was not failing because of low collision rate. This hypothesis was advanced by the observation that optimal two-frequency evaporation performance occurred for a frequency sweep with a 0.07 Hz rate that took over 10 seconds. It took this long to cool the gas from 600 to 100 nK, while the rethermalization time was probably shorter than 30 ms! We constructed a toy model [102] of the evaporation in order to explore technical limits that were not included in Murray's model.

We assume that we have experimentally found the best parameters (such as the evaporation speed) for the evaporative cooling procedure given our technical limits. The point behind this model is to then look at how changes in the technical limits impact the performance. The model is not designed to look at extensive re-optimization, but only small deviations from the existing conditions. Our hope was to find easy changes to the experiment that could lead to gains in the lowest achievable T/T_F .

We develop two models, one for simultaneous evaporation of two species and one for single species evaporation. The single species model is equivalent to the two species model for perfectly balanced evaporation, and includes a finite energy resolution (caused by magnetic field noise) in removing atoms. The two component model is computationally intensive, and does not include the finite energy resolution in the interest of expediency. We use these models only to explore the behavior of the twofrequency evaporation trajectory. The model uses a frequency sweep identical to the sweeps described in section 4.7.2.2 and the appendix to this chapter, and calculates the cut energy as described by equation 4.12. The sweeps are not divided into stages as in the experiment, but rather Δf is set according to equation 4.15 for balanced evaporation. Δf can also be adjusted in the model by hand to make the cut energy different for the two components.

4.4.2 Calculation Mechanics

The model divides the evaporation sweep into 100 ms steps. The choice of step size dt=100 ms is arbitrary and serves only as a convenient way to divide up the frequency sweep. As a check, the step size has been varied by a factor of two (both higher and lower) and no changes in the simulation results were observed. For each step *i* (which occurs at time *i*dt), the frequencies $\nu_{9/2,i}$ and $\nu_{7/2,i}$ and corresponding cut energies $E_{c,i}$ are calculated. See the appendix to this chapter for an explanation of the time dependence of the microwave frequency.

At step i in the iterative process, the energy and temperature removed by the microwave knife is calculated for either species. For the single species case, we have (see chapter 5)

$$dN_i = \frac{1}{2\lambda(\hbar\omega_r)^3} \int_{E_{c,i}}^{\infty} d\epsilon \frac{\epsilon^2}{\frac{1}{3_i} e^{\frac{\epsilon}{k_b T_i}} + 1} \mathcal{E}\frac{\operatorname{erf}(\epsilon/\Sigma) + 1}{2}$$
(4.16)

$$dU_i = \frac{1}{2\lambda(\hbar\omega_r)^3} \int_{E_{c,i}}^{\infty} d\epsilon \frac{\epsilon^3}{\frac{1}{3_i} e^{\frac{\epsilon}{k_b T_i}} + 1} \mathcal{E}\frac{\operatorname{erf}(\epsilon/\Sigma) + 1}{2}$$
(4.17)

where \mathfrak{Z}_i is the fugacity at step *i* and \mathcal{E} is the total probability to make a transition to the untrapped state. The integral in each case is done over the energy of the atoms ϵ in the magnetic trap.

The error function in equations 4.16 and 4.17 takes into account the finite energy resolution in removing atoms caused by noise in the bias field. We assume gaussian noise on B_0 with r.m.s. size δB_0 , which translates into an r.m.s. spread in cut energy $\Sigma = \frac{2}{9}m_f\mu_b\Delta B_0$ (see equation 4.12). The idea is then that as an atom moves in the harmonic potential past the cut energy, the gaussian noise is integrated into an error function.⁷ So, many Σ above E_c the atom has probability \mathcal{E} to be removed and many Σ below E_c the atom is not addressed by the microwave field and has no probability for removal. Note that for the two component case, equations identical to 4.16 and 4.17 are used except the $[\operatorname{erf}(\epsilon/\Sigma) + 1]/2$ is left out. This is done to speed up the calculation.

The fugacity at each step is solved via the number normalization condition:

$$\operatorname{Li}_{3}(-\mathfrak{Z}_{i}) = -\frac{1}{6\left(\frac{T_{i}}{T_{F,i}}\right)^{3}}$$

$$(4.18)$$

where $T_{F,i}$ is the Fermi temperature at step *i* and Li_n is the polylogarithmic function of order *n*.

For each step, the new N and U is calculated for each species:

$$N_i = N_{i-1} - dN_i - N_{i-1} \frac{dt}{350 \text{ sec}}$$
(4.19)

$$U_i = U_{i-1} - dU_i + 3\left(8 + 8\frac{N_{i-1}}{10^6}\right) \times 10^{-9} k_b N_{i-1} dt$$
(4.20)

where dt is the time step, 350 sec is the trap lifetime, and the third term in the equation for U_i is the experimentally measured heating rate in the "tight" trap (see the section on heating in this chapter). For the two component case, we have these expressions for each spin component.

For the single species case, the new temperature is calculated by solving the equations:

$$N_i = -\frac{1}{\lambda} \left(\frac{k_b T_i}{\hbar \omega_r}\right)^3 \operatorname{Li}_3(-\mathfrak{Z}_i) \tag{4.21}$$

$$U_i = -\frac{3}{\lambda} \frac{(k_b T_i)^4}{(\hbar \omega_r)^3} \operatorname{Li}_4(-\mathfrak{Z}_i)$$
(4.22)

simultaneously for the new value of T_i and z_i . For the two species case, we solve

$$N_{i,9/2} = -\frac{1}{\lambda} \left(\frac{k_b T_i}{\hbar \omega_{r,9/2}} \right)^3 \text{Li}_3(-\mathfrak{Z}_{i,9/2}) \quad (4.23)$$

$$N_{i,7/2} = -\frac{1}{\lambda} \left(\frac{k_b T_i}{\hbar \omega_{r,7/2}} \right)^3 \text{Li}_3(-\mathfrak{Z}_{i,7/2}) \quad (4.24)$$

$$U_{i,9/2} + U_{i,7/2} = -\frac{3}{\lambda} (k_b T_i)^4 \left[\frac{1}{(\hbar \omega_{r,9/2})^4} \text{Li}_4(-\mathfrak{Z}_{i,9/2}) + \frac{1}{(\hbar \omega_{r,7/2})^4} \text{Li}_4(-\mathfrak{Z}_{i,7/2}) \right]$$
(4.25)

 $^{^{7}}$ We use the same definition of the error function as Mathematica.

simultaneously, since $U_{i,9/2} + U_{i,7/2}$ is the conserved quantity.

4.4.3 Simulation Results

We start the simulation with the same conditions (N and T) as the data in figure 4.18. We find that we have to set $\mathcal{E} = 0.2$ to match the experimentally observed trajectory. The trajectory for the single species case as a function of \mathcal{E} is shown in figure 4.21. For this data, the energy resolution was set to the experimentally measured upper limit ($\delta B_0 = 2 \text{ mG}$, which gives $\Sigma/k_b = 135 \text{ nK}$). With only a modest increase in power, the removal efficiency would no longer be a limit to the evaporation. A transition probability of 0.2 is actually not unreasonable considering our experimental parameters (see the appendix to this chapter).

The effect of the width of the microwave knife is shown in figure 4.22. This calculation was done with $\mathcal{E} = 0.2$. Before the current servo optimization, the field noise was equivalent to 1350 nK energy resolution. At that time, the width of the knife was severely limiting the evaporation. The width (now equivalent to 135 nK) is currently not much of a limitation, however.



Figure 4.21: Evaporation simulation dependence on microwave knife efficiency. This simulation was done for a single species, with an energy resolution of $\Sigma/k_b = 135$ nK. Although it is not indicated in this plot or the next, there is a change in the number of steps in the simulation for different experimental parameters.



Figure 4.22: Evaporation simulation dependence on microwave knife energy resolution. This simulation was done for a single species with $\mathcal{E} = 0.2$.

Moving to the two species calculation, we show the trajectory without any heating (with $\mathcal{E} = 0.2$) in figure 4.23. The model suggests that the evaporation is actually limited by the combination of heating and an inefficient microwave knife. If the microwave removal efficiency is improved *or* the heating is eliminated, then the evaporation can reach lower T/T_F .



Figure 4.23: Evaporation simulation without heating for two components. This simulation was done with $\mathcal{E} = 0.2$.

Data showing the sensitivity to Δf is shown in figure 4.24. The case plotted with thin lines, with Δf set perfectly, compares to the experimental trajectories qualitatively well (compare to figure 4.18 — the difference in the frequencies is caused by a difference in B_0 between that data and the simulation). The case plotted with thick lines is with Δf systematically high by only 3 kHz throughout the entire trajectory. This eventually produces an imbalance in the spin mixture. However, the $m_f = 9/2$ component is still cooled efficiently. At some point, an error in Δf becomes catastrophic, as all of the atoms in one component will be removed before reaching low T.



Figure 4.24: Evaporation simulation dependence on Δf . This simulation was done with $\mathcal{E} = 0.2$. The bottom two plots are for the perfect Δf case, and can be compared qualitatively to data.

It is possible that we are not as efficient at removing the $m_f = 7/2$ atoms since the magnetic field from the microwave coil should predominantly produce circular polarization. To explore this effect, we show a trajectory in figure 4.25 with $\mathcal{E} = 0.3$ for the $m_f = 9/2$ component and $\mathcal{E} = 0.1$ for the $m_f = 7/2$ component. Δf is forced systematically high by 10 kHz during the trajectory to fix the spin composition to roughly 50%. This means that the microwave knife is cutting deeper in the $m_f = 7/2$ component compared to the $m_f = 9/2$ component. Having an unbalanced microwave knife is very detrimental to the evaporation performance — we find that the trajectory reaches only $T/T_F=0.4$ compared to 0.25 for the perfect case.



Figure 4.25: Evaporation simulation with unequal removal probability for different spin components. This simulation was done with $\mathcal{E} = 0.1$ for the $m_f = 7/2$ component and $\mathcal{E} = 0.3$ for the $m_f = 9/2$ component. Δf was forced systematically high by 10 kHz in order to fix the spin composition to roughly 50%.

Finally, we explore the dependence on the location of the "bottom" of the frequency sweep compared to ν_0 in figure 4.26. This data was taken for the spin mixed case, with $\mathcal{E} = 0.2$, extracting the best value of $T/T_{F,9/2}$ reached in the trajectory. This data is remarkably similar to the experimentally obtained results shown in figure 4.16. The best performance is with an aggressive cut at the end of evaporation, faster than an exponential with ν_0 as the endpoint. If heating is removed from the simulation, the dependence shown in this plot disappears. Presumably, the aggressive cut is needed to boost the cooling rate compared to the heating rate.



Figure 4.26: Evaporation simulation dependence on the location of the sweep "bottom" compared to ν_0 . The frequency resonant at B_0 , ν_0 , is indicated by a vertical line in this plot. The simulation was done for two species with $\mathcal{E} = 0.2$, although the minimum value of T/T_F is plotted only for the $m_f = 9/2$ component.

4.4.4 Evaporation Simulation Conclusions

There are some important conclusions to draw from the results of this toy model. First of all, the simulation can qualitatively describe the evaporation behavior without any information on the collision rate. This confirms our suspicion that the evaporation is not limited by low collision rate. Magnetic trap bias field noise and imperfect spin composition caused by errors in Δf are also not currently limits. Locating the frequency sweep "bottom" is important to the evaporation, but this does not really limit T/T_F .

According to the model, there are three limits to the evaporation performance: limited microwave power, heating, and the microwave power balance between transitions. A modest increase in microwave power (by a factor of two) should remove this as a limit, although other factors may then come into play. Experimentally, it might be fruitful to investigate the dependence on the relative power between $\nu_{9/2}$ and $\nu_{7/2}$, which we have not done. Reduced heating could also lead to the ability to reach lower T/T_F .

4.5 Heating Study

4.5.1 (Lack of) Models

We spent a great deal of time investigating the heating of ⁴⁰K atoms held in the magnetic trap. This detailed study was done in hopes of defeating the heating as a limit to the evaporative cooling (see section 4.4). Other experiments have observed heating due to inelastic collisions and hot "Oort" clouds of alkali atoms left behind by evaporation or populated by collisions with residual gas atoms. Two BEC experiments at JILA observe much lower heating rates than we do under similar conditions, while a third observes a heating rate comparable to ours. Carl Wieman's ⁸⁵Rb experiment has a 3-5 nK/sec heating rate in a "baseball" coil Ioffe-Pritchard trap.⁸ Eric Cornell's

⁸ This heating rate is measured using a Feshbach resonance to set the scattering length to $\sim 100 a_0$.

double-MOT ⁸⁷Rb experiment has only 1 nK/sec heating rate with atoms held in an 8 Hz TOP trap. The Cornell "People's BEC" experiment observes a ~ 10 nK/sec heating rate with ⁸⁷Rb atoms held in a $\omega_r = 2\pi \times 100$ Hz Ioffe-Pritchard trap. Our heating rate, with a similar number of atoms, tends to be between 10 and 20 nK/sec.

We find that the heating rate in our experiment scales linearly with the number of trapped atoms N, has some weak dependence on the trap parameters, and depends on the spin mixture in the same way as the collision rate. The heating does not depend on temperature, is not affected by a microwave "shield", and is not due to magnetic trap field noise. Unfortunately, these dependencies make it difficult to significantly reduce the heating rate relative to the cooling rate.

Our heating data are not consistent with any known heating models [103,104]. Because the states $(m_f = 9/2 \text{ and } m_f = 7/2)$ used in these studies do not undergo inelastic collisions, we rule out heating due to inelastic processes.⁹ Heating can also be caused by glancing collisions with residual gas atoms that result in ⁴⁰K atoms with energy less than the trap depth [103]. Glancing collision heating gives rise to an energy transfer rate (\dot{U}) to the gas proportional to N, or a heating rate $(\dot{T} \propto \dot{U}/N)$ independent of any of the ⁴⁰K gas properties. For this process, the ⁴⁰K gas is modelled as a stationary, spherical target. The probability for an incident residual gas atom to collide with a ⁴⁰K atom is proportional to $n_K \sigma l$, where l is the cloud size, σ the collision cross section, and n_K the ⁴⁰K gas density. The number of residual gas atoms that are incident on the ⁴⁰K "target" per second is proportional to $n_b l^2 v_b$, where n_b is the number density of residual gas atoms and v_b is the mean residual gas atom velocity. If $n_K \propto N/l^3$, then the overall rate of collisions with residual gas atoms is proportional to N and has no dependence on l. The rate of residual gas atom collisions is probably much lower than the collision rate in the gas, so that the energy deposition rate \dot{U} is then proportional to N.

⁹ Again, these states are affected by dipolar relaxation, but the rate should be so low that it cannot account for our heating rates.

Another process than can drive heating is secondary collisions between high energy ⁴⁰K atoms (created by glancing collisions or other processes) and cold ⁴⁰K atoms. As explained in [104], this kind of heating and "Oort" cloud heating give rise to a $\dot{U} \propto \Gamma n_K l$, which is temperature dependent through the dependence on the collision rate in the cold alkali gas Γ and l. Since we observe no temperature dependence over a wide range, we conclude that this heating mechanism is not present.

4.5.2 Characterization

The heating rate is measured by measuring the temperature of the gas for different hold times in the magnetic trap. The data in this section cover the classical to quantum regime. The heating rate is remarkably constant in time as seen in figure 4.27 for data taken in the tight trap. The temperature fits very well to a linear function of time, which implies that the heating rate is temperature independent. For the data shown in figure 4.27, the density in the gas decreased by over an order of magnitude and the collision rate by a factor of 30 because of the increase in T. This data also, therefore, suggests that the heating rate is independent of density and collision rate. It *is* possible that the heating rate is dependent on a combination of parameters such that the temperature dependence cancels. For example, if the heating rate was proportional to Γ/n_{pk}^2 , where Γ is the collision rate and n_{pk} the peak density in the gas, then the heating rate would appear to be temperature independent. Further, the linear heating implies that parametric heating caused by magnetic trap noise is not the culprit, since, in that case, the temperature should increase exponentially in time.

The data in figure 4.28 demonstrates the linear dependence of the heating rate on the number of atoms N. This data was taken while trying to determine if heating due to a hot Oort cloud was at work. By skipping various stages of evaporation, we tried to create a hot Oort cloud (or modify an existing one) of atoms and then measure the impact on the heating rate. This data in shown in figure 4.28. All of the data has



Figure 4.27: Heating of the trapped gas. A linear fit (solid line) is used to measure the heating rate, which is remarkably constant in time. This data was taken in the "tight" trap, with 2×10^6 total atoms and an approximately equal mixture of $m_f = 9/2$ and $m_f = 7/2$ atoms. The measured heating rate is 31.6(8) nK/sec.

the same scaling with N, which matches the normal dependence without any changes in the evaporation sequence. For the data in figure 4.28, the beginning of evaporation was varied from 1190-1250 MHz, and sections were skipped starting from 1269-1280 MHz. The final cut was fixed at 1282 MHz, which fixed the initial temperature to 1.43 μ K. This data suggests that Oort heating is not at work, as the character should change with the properties of the Oort cloud. This is also consistent with the observation that a microwave "shield" (an applied microwave field that is resonant far from the edge of the gas) has no effect on the heating.



Figure 4.28: Heating rate dependence on properties of hypothetical Oort cloud. The initial evaporative cut was varied and sections of the evaporation sequence were skipped in order to try and change the properties of an Oort cloud that might be causing heating. This data suggests that Oort heating is not a dominant source of heating in our experiment. This data was taken in the "tight" trap, with an approximate 50/50 mixture of spin components.

Noise on the magnetic trap current can also drive heating through parametrically driving energy into individual atom trajectories [105]. We measured the heating rate dependence on N for different traps before the servo optimization; the results are shown in table 4.1. We do not understand the dependence on the trap parameters (bias field and harmonic trapping frequencies), and the base heating with N = 0 seems to vary significantly over months and years. We found no significant difference in the heating rates immediately after the servo optimization. This suggests that the heating is not being caused by magnetic field noise, since the servo optimization reduced the bias field noise by a factor of 10. Further, we measured the effect on the heating rate caused by mechanically driving the coils. A PZT that was mounted rigidly to the optical table was butted up against one half of the bias coils and driven with a white noise source spanning DC-2 kHz. There was no difference in the heating rate with and without the drive on. Finally, we tested the heating rate dependence on the main FET drain-source voltage drop. We know that this voltage drop changes the servo behavior and the noise, but we observed no effect on the heating rate as shown in figure 4.29.

trap	N dependence
"load"	$\dot{T} = -14(6) + 19(4) \cdot N/10^6 \text{ nK/sec}$
"tight"	$\dot{T} = 8(6) + 8(1) \cdot N/10^6 \text{ nK/sec}$
"evap"	$\dot{T} = 32(7) + 8(3) \cdot N/10^6 \text{ nK/sec}$

Table 4.1: Heating rate dependence on trap parameters before servo optimization. Note that the heating rate in the "evap" and "load" trap were determined using temperatures measured after a ramp to the "tight" trap.



Figure 4.29: Heating rate dependence on magnetic trap power supply voltage. Changing the voltage changes the main FET drain-source voltage and the servo and current noise characteristics. From this data, which has the usual dependence on N, we conclude that the servo noise on this low level is not playing any role in heating. This data was taken in the tight trap, with an approximately 50/50 mixture of spin components.

We do have the ability to investigate the heating characteristics in a way that other experiments do not. We can change the spin composition at fixed total number and temperature. This varies the overall collision rate in the gas without changing the density or temperature. The results for this measurement at 190 nK in the tight trap is shown in figure 4.30. The heating rate is plotted versus both the spin composition and a quantity proportional to the overall collision rate. The heating rate depends on the spin composition in the same way as the overall collision rate in the gas.



Figure 4.30: Heating rate dependence on spin composition at 190 nK in the tight trap. The heating rate is plotted against the fraction of atoms in the $m_f = 9/2$ component, and a quantity proportional to the overall collision rate in the gas. The total number of atoms $N = N_{9/2} + N_{7/2}$ was kept relatively fixed to 2×10^5 in this measurement. Note that this data, which was taken ~ 1 year later than the other data in this section, shows a higher heating rate than we might expect from the previously measured dependence on N.

4.6 Chapter Conclusions

We have implemented a unique, two-component cooling scheme that allows us to cool the fermionic atom gas into the quantum degenerate regime. Since the spin degrees of freedom are frozen in our system, we tend to think of the 40 K gas as two separate gases, each of which maintains thermal equilibrium only through its contact (through collisions) with the other. Our "simultaneous cooling" scheme is designed to remove energy and atoms from each component in a balanced way by taking advantage of the adjustable magnetic trap parameters and multiple microwave frequencies. Simultaneous cooling allows us to cool the an equal mixture of spin components to temperatures as cold as $T/T_F \sim 0.2$.

By comparing the evaporation performance to a time dependent and a toy model, we find that the evaporation performance is probably technically limited. A model constructed by Murray Holland suggests that there is no fundamental limit to our cooling scheme. Indeed, the rethermalization rate in the gas should be high enough for evaporation to work well even at our lowest temperatures. And, technical improvements to the magnetic trap stability and microwave frequency sources over several years have allowed us to reach lower and lower T/T_F . The toy model suggests that the evaporation performance is currently (as of late 2000) limited by heating and insufficient microwave power.

The prospects are still bright for cooling the gas to even lower T/T_F . The microwave power could be increased by installing a higher power amplifier (used 20 W travelling wave tube amplifiers are easy to find), and the resulting effect on the evaporation performance could be readily tested. The toy model suggests that a study of the relative balance of microwave power between the transitions used for evaporation may be key to reaching lower T/T_F . Also, reducing the heating rate relative to the cooling rate may still be possible, although more heating studies are needed. Identifying the heating mechanism and understanding the heating rate dependence on the magnetic trap parameters may prove fruitful.

4.7 Appendix

4.7.1 Ramps

As mentioned in chapter 2, the ramps are executed by an Stanford Research Systems DS345 arbitrary waveform generator that is multiplexed with the stable control voltages. The ramp is programmed via GPIB into the waveform generator as a series of voltage-time coordinates. The generator is then TTL triggered to activate the ramp, and the multiplexer is switched at the beginning and end of the ramp. The requirements for adiabaticity are to ramp slower than the collision rate and harmonic trap frequencies. Also, the SRS voltage must be carefully matched to the control voltages at the beginning and end of the sweep in order to avoid exciting slosh. The slosh is excited by sudden changes in the trap center due to sag. In practice, the ramp time is limited on the long side by heating. The ramp time is chosen to produce a ramp that is as adiabatic as possible. The ramps are linear in the bias coil current vs. time. The ramp from the "load" to "evap" trap takes 100 ms, and the ramp from the "evap" to "tight" trap takes 525 ms.

4.7.2 Frequency Sweeps

4.7.2.1 Single Frequency

The frequency sweeps for single-frequency evaporation follow the form

$$f(t) = f_b + (f_0 - f_b)e^{-\alpha t - \beta t^2}$$
(4.26)

where f_b is the sweep "bottom frequency", f_0 is the start frequency, α is the evaporation rate, and β is an acceleration parameter. An exponential sweep removes an equal fraction of the energy in the gas per unit time for a classical gas, and is also generally very flexible. Exponential sweeps can easily be turned into linear sweeps, and multiple sweeps can be used to approximate any function. The frequency sweep is executed by

```
a loop of QuickBASIC code
```

```
starttime = TIMER
time = 0
IF beta = 0 THEN
    duration! = -1! / alpha! * LOG(1! - (gotofreq - startfreq) /
        (zerofreq - startfreq))
ELSE
    C = LOG(1! - (gotofreq - startfreq) / (zerofreq - startfreq))
duration! = (-alpha! + SQR(alpha! ^ 2 - 4 * C * beta!)) / 2 / beta!
END IF
PRINT "Evaporation will take "; duration!
CALL attenon(power)
WHILE time < duration!
   frq# = startfreq + (zerofreq - startfreq) *
       (1 - EXP(-time *(alpha! + beta! * time)))
   CALL IBWRT(microhp%, "FREQ " + LEFT$(STR$(frq#), 11) + " MHZ")
   gwait 40
   time = TIMER - starttime
WEND
CALL IBWRT(microhp%, "FREQ " + STR$(gotofreq) + " MHZ")
qwait 100
```

which sends a frequency set command to the synthesizer based on the time. In this way the most accurate sweep is produced. The qwait command (which adds a 40 ms pause) after the frequency set command is necessary in order to allow the synthesizer GPIB processing time.

The evaporation in the "evap" trap is broken up into four sweeps, from 1190 MHz to ~ 1282 MHz. The sweep parameters for each stage are optimized to maximize Nat the end of "evap1".¹⁰ The synthesizer power is constant for all the sweeps. The bottom frequency is always set to coincide with the $|9/2, 9/2\rangle \rightarrow |7/2, 7/2\rangle$ transition frequency at B_0 , typically 1282.7 MHz. The sweep parameters are given in table 4.2.

stage	α	β	duration
1190-1250 MHz	$0.03~\mathrm{Hz}$	$0 \ Hz^2$	$35 \sec$
1250-1270	0.07	0	13.7
1270-1278	0.07	0.001	14.8
1278-x	0.05	0	33-44

Table 4.2: Single frequency evaporation sweeps. The end frequency x for "evap1" ranges from 1281.8-1282.2 MHz.

¹⁰ Note that the evaporation parameters should be optimized for the highest N at fixed temperature.

The $m_f = 5/2$ and $m_f = 3/2$ removal sweep in the "tight" trap also uses this routine, with typical parameters $\alpha=10$ Hz, sweep from 1279.15 to 1279.25 MHz, sweep bottom at 1279.35 MHz. This sweep is optimized by shifting all three frequencies around together to minimize the temperature of the gas and the number of $m_f = 5/2$ and $m_f = 3/2$ atoms that are left afterward. The sweep rate α is set as described in the text. The width of the sweep may need to be changed if the initial cloud temperature changes significantly.

The $m_f = 7/2$ cleanout sweep at the end of "evap2" uses a single frequency evaporative sweep as well. Typical parameters are $\alpha = 0.07$ Hz, sweep from 1273.9-1274.2 MHz, sweep bottom at 1274.3 MHz. This sweep is optimized by minimizing the temperature of the remaining $m_f = 9/2$ component. A standard procedure that works well is: set the start of the sweep to $\nu_{9/2} + \Delta f$, move the sweep end and bottom together to minimize $T_{9/2}$ after the sweep, vary α to minimize $T_{9/2}$, then finally vary the sweep end relative to the sweep bottom.

4.7.2.2 Two Frequency

The sweeps for two-frequency evaporation follow the form

$$f1(t) = f_b + (f_0 - f_b)e^{-\alpha t}$$
(4.27)

$$f2(t) = f1(t) + \Delta f \tag{4.28}$$

where an acceleration parameter is not used. The frequencies f1 and f2 correspond to $\nu_{9/2}$ and $\nu_{7/2}$, respectively. Similar QuickBASIC code is used to execute the sweep: starttime = TIMER time = 0 CALL attenon(power) tau! = 1! / alpha! duration! = -tau! * LOG(1! - (gotofreq - startfreq) / (zerofreq - startfreq)) PRINT "Evaporation will take "; duration! WHILE time < duration! frq# = startfreq + (zerofreq - startfreq) * (1! - EXP(-time / tau!)) frq2# = frq# + deltaf! CALL IBWRT(microhp%, "FREQ " + LEFT\$(STR\$(frq#), 11) + " MHZ")

```
qwait 40
CALL IBWRT(micro%, "FREQ " + LEFT$(STR$(frq2#), 11) + " MHZ")
qwait 40
time = TIMER - starttime
WEND
CALL IBWRT(microhp%, "FREQ " + STR$(gotofreq) + " MHZ")
CALL IBWRT(micro%, "FREQ:CW " + STR$((gotofreq + deltaf! - fslope! *
(zerofreq - gotofreq))) + " MHZ")
qwait 100
```

where a pause after *each* GPIB command is required. The synthesizers tend to glitch if there is any activity on the GPIB bus during frequency switching.

The two-frequency evaporation is also broken up into stages. These stages are also individually optimized, and the sweep bottom frequency is set as described in this chapter. Although the single-frequency stage optimized parameters are constant over many months, the two-frequency stage parameters typically must be tweaked on a daily basis. A typical breakdown is described in table 4.3. A microwave pulse before the last stage is used to remove $m_f = 5/2$ and $m_f = 3/2$ atoms that are created in previous sweeps. We usually set α to be constant across all the sweeps as variation seems to have little effect. The sweeps in table 4.3 use 1272.7 MHz as a bottom frequency.

stage	α	Δf	duration
1272.15-1272.3 MHz	$0.07~\mathrm{Hz}$	$1.670 \mathrm{~MHz}$	$4.5 \sec$
1272.3-1272.4	0.07	1.677	$4.1 \sec$
1272.4-1272.45	0.07	1.676	$2.6 \sec$
500 ms pulse at 1279.25 MHz			
1272.45-	0.07	1.674	

Table 4.3: Two-frequency evaporation sweeps.

4.7.3 Landau-Zener Transitions

A dressed state description must be used to accurately calculate the probability of removing atoms from the trap via microwave transitions [106]. Consider the case of trying to remove an atom in state $|1\rangle$ by transferring it to $|2\rangle$. The microwave field couples the states at a location in space where the magnetic field from the trap causes the frequency to be resonant with the $|1\rangle \rightarrow |2\rangle$ transition. As the atom approaches the resonant region, the states adiabatically transform into each other in order to create an avoided crossing. A transition to the untrapped $|2\rangle$ state occurs if the resonant region is traversed adiabatically.

The probability to undergo a diabatic [107], or Landau-Zener transition, is given by

$$P = \exp\left[-2\pi \frac{|V_{12}|^2}{\hbar (dE/dt)}\right]$$
(4.29)

where V_{12} is the matrix element connecting states 1 and 2 and dE/dt is the rate of change of the energy difference between the states as the atom passes through the resonant region. The probability for the atom to be removed from the trap is then 1 - P.

In order to calculate P we must calculate the matrix element V_{12} . The term in the Hamiltonian that couples the states is $-\vec{\mu} \cdot \vec{B}$ where $\vec{\mu}$ is the total magnetic moment. We evaluate the matrix element by breaking up the total angular momentum into it's components: $\vec{F} = \vec{S} + \vec{L} + \vec{I}$. For the ground states, $\vec{L} = 0$ and we ignore the magnetic field coupling to the nuclear spin compared to the electron spin. The relevant matrix element is then $V_{12} = \langle 2|2\mu_b\vec{\sigma} \cdot \vec{B}|1 \rangle$ where $\vec{\sigma}$ is the Pauli spin matrix, μ_b is the Bohr magneton, and the factor of 2 is for the electron g-factor.

Let us consider the particular case of the transition $|9/2, 9/2\rangle \rightarrow |7/2, 7/2\rangle$. These states in the I,S basis are given by

$$|9/2, 9/2\rangle = |m_I = 4, m_s = +1/2\rangle$$
 (4.30)

$$|7/2,7/2\rangle = -\frac{1}{3}|m_I = 3, m_s = +1/2\rangle + \frac{2\sqrt{2}}{3}|m_I = 4, m_s = +1/2\rangle$$
(4.31)

using the appropriate Clebsch-Gordan coefficients. We also need an estimate for the magnitude of the oscillating magnetic field \vec{B} . We estimate the current in the coil by

assuming that all 2.5 W¹¹ are coupled into the coil as a 50 Ω load. If we treat the coil as a circular loop of wire with a 1 cm radius located 5 cm away from the atoms, we find that the magnitude of the field experienced by the atoms is ~ 1 mG. We assume that the field is equally distributed over all three directions, giving $\vec{B} = \frac{1}{\sqrt{3}} \frac{mG}{\sqrt{3}} (\hat{x} + \hat{y} + \hat{z})$. Writing the Pauli spin matrices in a spherical basis as $\vec{\sigma}_x = \frac{1}{2} (\sigma_+ + \sigma_-)$ and $\vec{\sigma}_y = \frac{1}{2} (\sigma_+ - i\sigma_-)$, we find that the matrix element is $V_{12} = -\frac{2\sqrt{2}}{3} 2\mu_b \frac{1}{2} (B_x + B_y)$.

In order to calculate P we also need dE/dt. We approximate dE/dt using the uncoupled states. The relevant microwave transition frequency changes at $\Lambda_{9/2,7/2}$ =-2.5 MHz/gauss at low field, so that $dE/dt = \Lambda_{9/2,7/2}h(dB/dt)$ where dB/dt is the rate of change of magnetic field (from the trap) that the atom experiences. In the radial direction, we can write the magnitude of the trap magnetic field as $B = B_0 + B''r^2$, so that dB/dt = 2B''r(dr/dt) for the moving atom. Our final evaporative cuts are about 10 μ m from the center of the trap, and we assume that we are removing atoms with ~ 300 nK of kinetic energy. Using the trap parameters, we find that 1 - P = 0.2, which is a fairly low probability of atom removal. With only a factor of two increase in microwave power, 1 - P increases to 0.37 and the model outlined in section 4.4 predicts that the microwave power should no longer be a limit to the evaporative cooling.

 $^{^{11}}$ For the two-frequency evaporation, only 2.5 W from the 5 W amplifier are available to drive either transition.

Chapter 5

THERMOMETRY AND THERMODYNAMIC MEASUREMENTS

5.1 Overview

In the summer of 1999 we were first able to detect the emergence of quantum behavior in an evaporatively cooled gas of 40 K atoms [53]. We detected FD statistics through measurements of the momentum distribution and the energy of the gas. In our lab, we now had the world's first and only trapped Fermi gas of atoms. In early 2001, two other groups (see [68] and C. Salomon, unpublished) succeeded in cooling a gas of the fermionic atom ⁶Li into the degenerate regime.

This chapter will begin by deriving a laundry list of thermodynamic relations for a trapped Fermi gas, as well as other useful quantities such as the density distribution and expanded column density. We will assume in this chapter than mean-field type interactions are negligible; this is equivalent to dealing only with an ideal gas. Final equations will be boxed for easy reference. The equations derived in this chapter will be put into a form that is convenient for people working with trapped Fermi gases.

I will then explain, in detail, the thermometry techniques that we have developed. Our first results probed the momentum distribution of the gas using a fitting technique that did not assume FD statistics. Later, we moved on to probing the Thomas-Fermi (TF) shape of the gas directly. These fitting techniques not only allow us to determine N and T, but can provide an independent " T/T_F -meter". I will also explain the details of a method that we have devised for measuring the mean energy per particle in the
gas.

At the end of the chapter I will highlight experimental results. I will show the data where we first detected deviation from classical statistics at low T in the momentum distribution of the gas. Results from the TF thermometry will also appear, demonstrating the ability to directly measure the fugacity of the gas, which is only a function of T/T_F . The TF thermometry also provides a very clean method for determining the absolute temperature T of the gas. Direct comparisons between the observed momentum distributions and calculated classical distributions for gases with the same N and T will dramatically demonstrate the effects of FD statistics in the gas. Finally, measurements of the mean energy per particle in the gas show a factor of two "excess" energy in the gas compared to the classical expectation at our lowest temperature.

5.2 Everything You Wanted to Know About a Trapped Fermi Gas but were Afraid to Ask

5.2.1 Trapped Fermi Gases — Statistical Mechanics

All the thermodynamics of a non-interacting (or ideal) trapped Fermi gas can be derived from straightforward statistical mechanics. The single particle Hamiltonian for an atom in a harmonic potential is

$$H = \frac{1}{2m} \left(p_x^2 + p_y^2 + p_z^2 \right) + \frac{m\omega_r^2}{2} \left(x^2 + y^2 + \lambda^2 z^2 \right)$$
(5.1)

where $\lambda = \omega_z / \omega_r$ is the trap asymmetry parameter. The density of states is given by [27]

$$g(\epsilon) = \frac{\epsilon^2}{2\lambda(\hbar\omega_r)^3} \tag{5.2}$$

as a function of the energy ϵ . This is an approximation to the density of states that is valid when the temperature of the gas is much greater than $\hbar\omega_r$. And, of course, the Fermi-Dirac distribution function [108,109] is

$$\mathcal{F}(\epsilon) = \frac{1}{\frac{1}{3}e^{\frac{\epsilon}{k_bT}} + 1}$$
(5.3)

with \mathfrak{Z} as the fugacity. The fugacity is a convenient parameter to deal with as it is a function of only T/T_F .

The Fermi temperature T_F is defined by the Fermi energy E_F , which is the energy of the highest occupied energy level in the harmonic potential at T = 0. The Fermi energy can be determined through integrating the equation for N at T = 0:

$$N = \int_0^{E_F} g(\epsilon) \ d\epsilon \tag{5.4}$$

where every level up to E_F is fully occupied. The integral is trivial, and yields:

$$T_F = \frac{E_F}{k_b} = \frac{\hbar\omega_r}{k_b} \left(6\lambda N\right)^{1/3}$$
(5.5)

which is only a function of the trap parameters and N. Note that the harmonic trapping frequencies depend on the mass and magnetic moment of the atom.

Fortunately, all the common thermodynamic integrals have an analytic representation in terms of a hyper-geometric function that is known to Mathematica. For n > 1(where n is an integer or half-integer),

$$\int_0^\infty d\epsilon \frac{\epsilon^n}{\frac{1}{3}e^{\frac{\epsilon}{k_bT}} + 1} = -\left(k_bT\right)^{1+n} \Gamma(1+n)Li_{1+n}(-\mathfrak{Z})$$
(5.6)

where Γ is the Euler gamma function (use Gamma[1+n] in Mathematica) and Li_n is the Poly-Logarithmic function of order n (use PolyLog[n,-z]).¹ Maxwell-Boltzmann (classical) statistics holds for $\mathfrak{Z} << 1$, with $Li_n[-\mathfrak{Z}] \to -\mathfrak{Z}$ to first order in \mathfrak{Z} . In order to perform some of the integrals in this chapter, a series expansion for Li_n

$$Li_n(-\mathfrak{Z}) = \sum_{k=1}^{\infty} (-\mathfrak{Z})^k / k^n$$
(5.7)

¹ All of the thermodynamic equations in this chapter can be modified for Bose gases by taking $-Li_n[-3] \rightarrow Li_n[3]$ and including a term for the ground state which is not counted properly because of the approximation used for $g(\epsilon)$.

must be used and the integrals evaluated term by term.

We can now determine the integrals for $N = \int_0^\infty g(\epsilon) \mathcal{F}(\epsilon) \ d\epsilon$ and the total energy in the gas $U = \int_0^\infty \epsilon g(\epsilon) \mathcal{F}(\epsilon) \ d\epsilon$:

$$N = -\frac{1}{\lambda} \left(\frac{k_b T}{\hbar \omega_r}\right)^3 Li_3(-3)$$
(5.8)

$$U = -\frac{3}{\lambda} \frac{(k_b T)^4}{(\hbar \omega_r)^3} Li_4(-3)$$
(5.9)

Rearranging the equation for N and substituting in the definition of T_F leads to:

$$Li_3(-\mathfrak{Z}) = -\frac{1}{6(T/T_F)^3}$$
(5.10)

which must be solved numerically in order to determine \mathfrak{Z} .

Another quantity of interest that is related to U is the mean energy per particle E = U/N, given by

$$E = 3k_b T \frac{Li_4(-3)}{Li_3(-3)}$$
(5.11)

which reduces to the classical expectation for $\Im << 1$. We will discuss measurements of $E/3k_bT$, which is only a function of T/T_F .

Universal thermodynamic plots for harmonically trapped ideal FD gases are shown in figure 5.1. As $T \to 0$, E approaches $\frac{3}{4}E_F$. This means that the extra fractional energy $E/3k_bT$ compared to a classical gas diverges at T = 0. This "excess" energy is present because atoms cannot all fall to the lowest levels of the harmonic potential. The fugacity is shown on a log plot (it is a very steep function of T/T_F at low T) and approaches 0 as T/T_F becomes large. We also plot the chemical potential $\mu = k_bT \cdot \ln(\mathfrak{Z})$, which approaches E_F at T = 0. The chemical potential crosses zero at $T/T_F \sim 0.55$ which coincides with a phase space density equal to unity.



Figure 5.1: Universal thermodynamic plots for a harmonically trapped Fermi gas. The classical expectation is plotted as a dashed line.

5.2.2 Density and Momentum Distribution

The Thomas-Fermi (TF) approximation [27] is valid for the range of temperatures that we access in this experiment. In the TF picture, a semi-classical phase-space distribution [27,108] $w(\vec{r}, \vec{k})$ is used to describe the many-body wavefunction of the gas. Each point in phase-space is assigned a definite position and momentum, with the distribution given by:

$$\mathbf{w}(\vec{r},\vec{k}) = \frac{1}{(2\pi)^3} \frac{1}{\frac{1}{\frac{1}{3}e^{\frac{H(\vec{r},\vec{k})}{k_bT}} + 1}}$$
(5.12)

where $\vec{p} = \hbar \vec{k}$. Since we tend to think in terms of \vec{p} , the phase space distribution can be written as $w(\vec{r}, \vec{p}) = \frac{1}{\hbar^3} w(\vec{r}, \vec{k})$ with the \hbar^3 present to fix the normalization for number.

The density and momentum distributions are then determined by appropriate integrals over w. The density distribution is given by

$$n(\vec{r}) = \frac{1}{(2\pi\hbar)^3} \int d^3 \vec{p} \, \mathbf{w}(\vec{r}, \vec{p})$$
(5.13)

$$= \frac{1}{(2\pi\hbar)^3} 4\pi \int_0^\infty dp \, p^2 \, \mathbf{w}(\vec{r}, p^2) \tag{5.14}$$

$$n(\rho) = -\frac{(k_b m T)^{3/2}}{(2\pi)^{3/2} \hbar^3} Li_{3/2} \left(-\Im e^{-\frac{m\omega_r^2}{2k_b T}\rho^2} \right)$$
(5.15)

where $\vec{\rho}$ is a scaled coordinate with magnitude $\rho^2 = x^2 + y^2 + \lambda^2 z^2 = r^2 + \lambda^2 z^2$. Likewise, the momentum distribution² is given by

$$\Pi(p) = \frac{1}{2\pi\hbar^3} \int d^3 \vec{r} \, \mathbf{w}(\vec{r}, \vec{p})$$
(5.16)

$$= \frac{1}{(2\pi\hbar)^3\lambda} 4\pi \int_0^\infty \rho^2 d\rho \,\mathrm{w}(\vec{\rho}, p) \tag{5.17}$$

$$\Pi(p) = -\frac{1}{(2\pi)^{3/2}\hbar^3\lambda} \left(\frac{k_b T}{m\omega_r^2}\right)^{3/2} Li_{3/2} \left(-\Im e^{-\frac{p^2}{2mk_b T}}\right)$$
(5.18)

where $p^2 = p_x^2 + p_y^2 + p_z^2$.

² The trick to doing these integrals over space is to define a new variable $Z = \lambda z$. The integral can then be transformed to an integral over dZ by substituting $dz = dZ/\lambda$. Then integrate over the scaled (elliptical) coordinate $\vec{\rho}$ with magnitude $\rho^2 = r^2 + Z^2$ and $d^3\rho = \rho^2 d\rho d(\cos\theta) d\phi/\lambda$.

It is also worth determining the density and momentum distributions in the classical limit, which are:

$$n_c(\rho) = \frac{\lambda N}{(2\pi)^{3/2} \sigma_r^3} e^{-\frac{\rho^2}{2\sigma_r^2}}$$
(5.19)

$$\Pi_c(p) = \frac{N}{(2\pi)^{3/2} \sigma_p^3} e^{-\frac{p^2}{2\sigma_p^2}}$$
(5.20)

where $\sigma_r^2 = \frac{k_b T}{m \omega_r^2}$ and $\sigma_p^2 = m k_b T$.

5.2.3 Expansion from the Magnetic Trap

The images we use to extract thermodynamic data are always taken after some expansion from the magnetic trap. Theoretically, Bruun and Clark [61] have been able to show that a free expansion from a harmonic trap of an *ideal* Fermi gas amounts to rescaling the spatial coordinates in $n(\rho)$. By considering the evolution of the manybody wavefunction with the boundary condition $\omega_i \to 0$ at t = 0, they find that the expansion transformation is $x_i \to x_i/\sqrt{1 + \omega_i^2 t^2}$ (where x_i refers to x, y, or z) and $n(\rho) \to \frac{n(\rho)}{(1+\omega_r^2 t^2)\sqrt{1+\omega_z^2 t^2}}$. This is a remarkable result — the shape of the density distribution does not change in time after the release from the magnetic trap (other than a rescaling). The shape invariance under expansion is particular to harmonic traps, and does not hold for a general potential.

5.2.4 Expanded Column Density

We actually measure the column density, or the density profile integrated through in one direction. The form for the column density in the trap can be determined by integrating $n(\vec{\rho})$ through in x (expanding the hypergeometric function according to equation 5.7 and then integrating term by term). The observed optical depth (OD) is the column density multiplied by the photon absorption cross-section $\sigma_{\lambda} = \frac{3\lambda^2}{2\pi}$, where λ is the transition wavelength.³ After expanding for time t, we end up with:

$$OD(y,z) = -\frac{\sigma_{\lambda}}{2\sqrt{1+(\omega_r t)^2}\sqrt{1+(\omega_z t)^2}} \frac{m(k_b T)^2}{\pi\hbar^3\omega_r} Li_2\left(-\Im e^{-\frac{y^2}{2\sigma_r^2}} e^{-\frac{z^2}{2\sigma_z^2}}\right)$$
(5.21)

$$\sigma_r^2 = \frac{k_b T}{m\omega_r^2} \left[1 + (\omega_r t)^2 \right] \qquad (5.22)$$

$$\sigma_z^2 = \frac{k_b T}{m\omega_z^2} \left[1 + (\omega_z t)^2 \right] \qquad (5.23)$$

which reduces to a gaussian in the classical limit. The coordinate rescaling has been lumped into the widths σ in this expression. At long expansion times $((\omega t)^2 \gg 1)$ the aspect ratio, $\frac{\sigma_z}{\sigma_y} = \frac{\omega_r}{\omega_z} \sqrt{\frac{1+(\omega_z t)^2}{1+(\omega_r t)^2}}$, goes to unity so that the gas appears spherical. In this limit, the initial size of the gas is unimportant compared to the final size. The final size of the gas then reflects the kinetic energy and the expanded distribution is equivalent to the momentum distribution, which is isotropic.

5.3 Minimal Assumption Thermometry — The "Mixture" Fits

In order to use a thermometric method that does not assume FD statistics, we fit the expanded OD images to the following functional form [53] which varies smoothly between the T = 0 and classical limits:

$$OD(\mathcal{R}) = \begin{cases} A \left(1 - \frac{\mathcal{R}^2}{R^2}\right)^2 & , 1 - \frac{\mathcal{R}^2}{R^2} \ge L \\ Be^{-\frac{\mathcal{R}^2}{2}} & , \text{otherwise} \end{cases}$$
(5.24)

Here, \mathcal{R} is a scaled distance $\mathcal{R} = \sqrt{\frac{y^2}{\sigma_{m,y}^2} + \frac{z^2}{\sigma_{m,z}^2}}$ from the peak of the distribution, and A (the peak OD in the image), $\sigma_{m,y}$, $\sigma_{m,z}$, and L are fit parameters.⁴ The requirement of continuity of the function and its first derivative at the boundary of the inner quartic and the outer gaussian form fixes the parameters $B = AL^2e^{2\frac{1-L}{L}}$ and $R = 2/\sqrt{L}$. The

³ This neglects effects that may change the effective absorption cross-section as the light passes through the gas. The value for σ_{λ} assumes that the light is right circularly polarized, the light intensity is very low compared to I_{sat} , the atoms are in the $f = 9/2, m_f = 9/2$ state, and the light is resonant with the $f = 9/2, m_f = 9/2$ to $f' = 11/2, m_f = 11/2$ transition.

⁴ In practice, for all fits to the images we allow for a background plane (a term $+b + m_y \cdot y + m_z \cdot z$ in the OD) in order to account for imaging imperfections. Also, in these equations y and z have an offset in order to locate the center of the cloud. See the TF fit for more details.

parameter L characterizes the deviation from the classical gaussian profile with L = 1at $T/T_F >> 1$ and L = 0 at $T/T_F = 0$.

We call these fits "mixture fits" since a gaussian is connected, or mixed, with the T = 0 TF shape at some radius in the gas. In order to figure out how the fit parameters are connected to T, we generated simulated OD images using equation 5.21 and fit them to the mixture fit form. From these fits, we determine correction factors to the temperature calculated directly from σ_m . This is cheating a bit on our desire to avoid assuming FD statistics, but we find that the corrections are relatively small.⁵

To determine T and T/T_F , first we find the number of atoms from the fit:

$$N = \frac{2\pi}{\sigma_{\lambda}} \cdot A \cdot \sigma_{m,y} \cdot \sigma_{m,z} M^2 \frac{L^3 + 2}{3L}$$
$$= 22.4 \mu \mathrm{m}^{-2} \cdot A \cdot \sigma_{m,y} \cdot \sigma_{m,z} M^2 \frac{L^3 + 2}{3L}$$
(5.25)

where σ_m is measured in pixels and M is the imaging magnification measured in We then determine an uncorrected $T = \frac{m\omega^2 \sigma_m^2}{k_b [1+(\omega^2 t^2)]}$ from either direc- $\mu m/pixel.^6$ tion and corresponding uncorrected T/T_F with T_F determined by the trap parameters and N. From these values, we calculate the actual $T_{corr} = T \left[1 - 0.8e^{-\frac{T/T_F}{0.234}} \right]$ and corresponding $(T/T_F)_{corr} = T_{corr}/T_F$. These corrections come from testing the "mixture" fits on simulated distributions.

$\mathbf{5.4}$ **TF** Thermometry

The thermometry technique we use most commonly now is a fit of the image to the TF profile [55]. The exact function we fit to, based on equation 5.21 is

$$OD(y,z) = A \cdot \frac{Li_2 \left[-\Im e^{-\frac{(y-yc)^2}{2\sigma_{TF,y}^2}} e^{-\frac{(z-zc)^2}{2\sigma_{TF,z}^2}} \right]}{Li_2(-\Im)} + b + m_y y + m_z z$$
(5.26)

where the fitting parameters are detailed in table 5.1. The temperature is determined

⁵ For the data shown in figure 5.7, the correction was at most 7%. ⁶ The factor of 22.4 uses $\sigma_{\lambda} = \frac{3\lambda^2}{2\pi}$.

parameter	
$\sigma_{TF,y}$	$\sqrt{rac{k_bT}{m\omega_r^2}\left[1+(\omega_rt)^2 ight]}$
$\sigma_{TF,z}$	$\sqrt{rac{k_bT}{m\omega_z^2}}\left[1+(\omega_z t)^2 ight]$
A	peak OD
b	background plane offset
m_y, m_z	background plane slopes
yc, zc	cloud centers
3	fugacity

Table 5.1: TF fit parameters.

from the widths σ_{TF} , while the number of atoms is determined from a simple gaussian fit to the same image.⁷ The gaussian fit makes at most a 2% error in N across the temperatures that we sample $(T/T_F > 0.15)$.

⁷ The number from a gaussian is $N = \frac{2\pi}{\sigma_{\lambda}} \cdot A \cdot \sigma_{g,x} \cdot \sigma_{g,z} M^2$ where σ_g from the simple gaussian fit is measured in pixels and M is the imaging magnification measured in μ m/pixel.

5.5 Imaging and Thermometry Systematics

Because all of our thermometric techniques rely on detecting relatively small deviations from a gaussian (see figure 5.7 inset), we are particularly sensitive to imaging imperfections. Most imaging imperfections can mimic the effects of quantum degeneracy on the shape of the expanded distribution and/or introduce systematic errors in T. Images of clouds with a width too close to the imaging resolution are significantly distorted. This is avoided by only using images for expansion times greater than or equal to 10 ms, where the r.m.s. cloud size is at least 18 times the imaging resolution. In this limit, the observed cloud size should never be distorted by more than 0.2%. In fact, we do not observe a *significant* effect on the thermometry, in the measured T, z or L, for expansion times from 10-25 ms.

Any light present in the probe beam that the atoms cannot absorb that is not properly subtracted will distort the shape of the image. Light present in the probe beam that is of the wrong polarization or far off of resonance will not be absorbed by the atoms. Because the light is present only in the probe beam and not the "dark" frame, it will not be properly subtracted in the OD math (see Chapter 2). If $\alpha = e^{OD}$ where OD is the actual OD, then the measured OD is

$$OD_{meas} = ln\left(\frac{L+B}{L/\alpha+B}\right) \tag{5.27}$$

where L is the amount of absorbable light and B is the amount of unabsorbable light that is not subtracted in the dark frame. If we define c = L/B as the ratio of absorbable to uncounted unabsorbable light, then the observed OD is:

$$OD_{meas} = ln\left(\frac{c+1}{ce^{-OD}+1}\right) \tag{5.28}$$

and the maximum observable OD is ln(1+c). With our maximum observable OD of 2.5, we estimate c = 11. The effect of fitting a gaussian column density with the OD modified by equation 5.28 is shown in figure 5.2. From data taken with classical (gaussian) clouds, it appears that we are actually operating at much higher c. This discrepancy may be explained if lensing or scattering of light outside of the gas is playing a role in limiting the OD. Experimentally we find that shape distortions are avoided by probing clouds with a peak OD at least a factor of two smaller than the maximum observable OD.



Figure 5.2: Effect of unabsorbable, uncounted light on the mixture fit thermometry. The plot on the left is experimental data and the plot on the right is from a theoretical calculation; in both cases the value of the non-gaussian mixture fit parameter L should be 1. The theoretical plot (right) is the result fitting a simulated, gaussian image that has been modified by the presence of "bad" light for two different ratios of "good" to "bad" light (c). At high OD, the "bad" light mimics the effects of quantum degeneracy, driving the measured L to below 1. The data plot (left) was taken with clouds for $T/T_F > 1.7$, so that the image should have a gaussian (classical) profile. We observe a very weak systematic with the measured peak OD that does not match the theoretical prediction (we estimate that we have c = 11). It is difficult to understand this result, and this systematic deserves more investigation.

In order to avoid a similar systematic caused directly by saturation of the atomic resonance, we probe with $I/I_{sat} < 0.05$. A measurement of I_{sat} is shown in figure 5.3. The number of counts measured by the CCD camera per pixel in the probe beam is plotted vs. the measured peak OD of the gas. The vertical axis in these plots is essentially an experimental measure of the probe intensity. By fitting the dependence to the expected form, $I = I_{sat} \left(\frac{OD - OD_{meas}}{1 - e^{-OD_{meas}}} \right)$ [110] (here OD is the actual OD and OD_{meas} is the observed OD) we find that I_{sat} corresponds to 1400(200) counts/unbinned pixel for a 24 μ s probe pulse.⁸ We also find that for this particular gas, the actual peak OD was 2.29(5).



Figure 5.3: The effect of saturation of the atomic resonance on imaging. The probe intensity, measured in CCD counts per unbinned pixel for a 24 μ sec probe pulse, is plotted vs. the observed peak OD. A fit (solid line) to the expected dependence allows to extract I_{sat} in experimental units. The data on the right is data taken across the region in which we normally operate. There is no observed effect of saturation on the peak OD.

⁸ The CCD camera is capable of fast hardware binning, with the result of multiple pixels summed into a single pixel.

Lensing caused by the density gradient of the cloud [111] will also cause image distortions. In order to avoid this systematic, we probe on resonance and carefully focus the camera. The effect of probe detuning on the cloud parameters measured via the TF fit is shown in figure 5.4. Both the fugacity and the widths used to determine Tare systematically affected. While not shown here, widths measured using gaussian fits (which are used to determine N) are also affected. We find experimentally that the all of the observed thermodynamics are only consistent when probing on resonance. We have not taken enough data to know that the effects of detuning are conclusively due to lensing. However, the fact we have observed the cloud width dependence on detuning change sign after a camera refocusing is rather suggestive.



Figure 5.4: Effect of detuning on the TF thermometry. We observe that the probe beam detuning affects the measured shape (fugacity) of the image and the measured widths. These effects may be caused by lensing. As the probe beam is detuned from resonance, the gas acts like a lens and causes shape distortions in the image. This systematic affects both degenerate and classical gases. Unfortunately, these effects are not very reproducible day to day and cannot be easily corrected, presumably because the lensing action of the gas depends on the OD and size.

The worst systematic in our thermometry is caused by features in the background field of the image. These can be difficult to spot by eye, but seem to be characterized by curvature, stripes, or patches in the image background. These features appear to be caused by spatial variations in the intensity profile of the probe beam or fringes introduced by clipping of the beam. Data taken with an almost undetectable curvature in the image background field is shown in figure 5.5. The thermometry cannot be trusted under these circumstances.

Problems caused by the probe beam have been solved by the probe beam preparation scheme described in Chapter 2. Our probe beam is formed from the image of a small aperture placed in the center of a gaussian beam — this produces a relatively flat intensity profile. In addition, reducing background light (in the "dark" frame) by shielding the camera is important.

We find that the "mixture" and TF profile fits have trouble with images with a peak OD less than 0.35. Typically, the fits lose the ability to determine that the profile is non-gaussian (in the degenerate regime) at this point. We suspect that this problem is caused by narrow (3-4 pixels wide) stripes in the background image field that seem nearly impossible to eliminate. These stripes have a peak-to-peak variation of 0.01 in OD.



Figure 5.5: Bad effect of features in the probe beam on the mixture fit thermometry. This data was taken with an almost imperceptible curvature in the background field of the image. The solid line is the theoretical expectation. These kinds of problems can be detected by looking for systematics $(L \neq 1 \text{ or } \mathfrak{Z} \neq 0)$ for high T/T_F (classical) clouds.

5.6 Check on the Expansion

An important check on the thermometry is that the expansion follows the form described in section 5.2.3. We must be sure the temperature can be accurately extracted from one point during the expansion. The width of the cloud measured with the TF fit as a function of the expansion time for a $m_f = 9/2$ gas at 150 nK and $T/T_F = 0.4$ is shown in figure 5.6. A fit to the time dependence indicates that the expansion is consistent with form from 5.2.3: $\sigma(t)_{TF} = \sigma_{0,TF}\sqrt{1+(\omega t)^2}$ where $\sigma_{0,TF}$ is the size in This is also a statement that T extracted from each point is consistent, the trap.⁹ where we find that T is 152(8)nK and 144(8) nK from the axial and radial directions, respectively. There is no systematic shift in the measured T with the expansion time. However, there is a systematic difference between the radial and axial directions. We are not sure where this error comes from, but extensive studies have revealed that the vertical direction (radial) result produces self-consistent thermodynamics.¹⁰ The scatter in the measured T across different expansion times is consistent with the noise at fixed expansion time for repeated runs of the experiment.

⁹ The data in figure 5.6 cannot actually resolve the initial vertical size.

¹⁰ Actually, this sort of aspect ratio error seems to change slightly in the experiment from month to month.



Figure 5.6: Measured TF fit widths and corresponding temperatures as a function of the expansion time. The vertical expansion time is actually 0.2 ms longer than shown in the plot (for an explanation of the magnetic trap turn-off, see chapter 2), which is added into the fits (solid lines) to the expected form for the expansion.

5.7 Energy Measurements

One of the most dramatic features of the trapped Fermi gas is the fact that the average energy per particle E does not go to zero at T = 0. The trapped atoms have relatively large kinetic energy, $\frac{3}{8}E_F$ per atom, at T = 0. This physics tends to make the biggest impression on laypeople, as most everyone is taught in grade school that "all motion ceases at absolute zero". We are able to directly measure E of the gas from absorption images.

E can be connected to the second moment of the image at any expansion time. We define the second moment (in scaled coordinates) of the image as:

$$J_{2} = \frac{\int_{0}^{\infty} dz \int_{0}^{\infty} dy |OD(y, z)| m\omega_{r}^{2} \left[\frac{2y^{2}}{1 + (\omega_{r} t)^{2}} + \frac{\lambda^{2} z^{2}}{1 + (\omega_{z} t)^{2}} \right]}{\int_{0}^{\infty} dz \int_{0}^{\infty} dy OD(y, z)}$$
(5.29)

where the factor of 2 in $2y^2 + z^2$ takes into account that we are looking at the column density which is a projection of the 3 dimensional density onto a plane (we also assume that the radial symmetry is not broken). In this equation, y and z refer to the distance from the center of the cloud. The integrals can be done directly in cartesian coordinates using the expansion for the polylogarithmic function and the equation for the expanded column density, giving

$$J_2 = m\omega_r^2 \left[\frac{2\sigma_y^2}{1 + (\omega_r t)^2} + \frac{\lambda^2 \sigma_z^2}{1 + (\omega_z t)^2} \right] \frac{Li_4(-3)}{Li_3(-3)}$$
(5.30)

which, after substituting for σ (which is $\sqrt{\frac{k_b T}{m\omega^2} [1 + (\omega t)^2]}$ — see the section on the expanded column density distribution) gives the correct value of E. To measure E, we measure the second moment of the image as defined by J_2 .

This correspondence between the second moment of the image and E in general is no accident. At long expansion times, the second moment of the image, as defined here, is proportional to the twice kinetic energy in the gas. Simply put, each atom can be viewed as expanding from a point source by r = vt in this limit, and then the second moment (the average of r^2) is proportional to v^2 . In the opposite limit — in the trap — the second moment is proportional to twice the potential energy. What is surprising is that this measurement works in between long and short expansion times. This is a consequence of the behavior of a free expansion from a harmonic trap. The definition of J_2 that I give above is designed to extrapolate back to the conditions in the trap and calculate twice the potential energy.

We have used two slightly different methods of measuring E. One, the "direct" method, calculates J_2 directly from the image. The other, the "fitting" method, uses a specially weighted fit to extract the second moments in each direction separately. It is important to note that any method of measuring J_2 must deal with the low signal to noise section of the image at large ρ . Noise in this part of the image can cause large errors in E since it does not average away (note the absolute value in J_2) and is heavily weighted. Both of our methods use a gaussian fit to the image to help with this problem. The widths from the gaussian fits, σ_g , are used to define an inner and outer part of the image. As with the TF and "mixture" fits, we find that at least a peak OD of 0.35 is necessary to get accurate results. Presumably, problems caused by small stripes in the image are also the culprit here.

5.7.1 Direct Method

In the direct method, the raw image data is used to compute the second moment of the image. Data inside a cut region defined by $\frac{y^2}{\sigma_{gy}^2} + \frac{z^2}{\sigma_{gz}^2} = j^2$ (where j is a number) is used to calculate J_2 on a pixel by pixel basis. The background plane of the image is subtracted off using the values from the TF fit, since the simple gaussian fit (and the "mixture" fit, for that matter) makes significant systematic errors in the fit parameter b (the offset for the image background plane) at low T/T_F .

A correction must be made for the contribution to the second moment outside of the cut region. To this end, only the data outside of the cut region is fit to a gaussian, returning widths σ_{ex} . For large enough j ("large enough" depends on T/T_F), the outer part of the image actually is gaussian, since it corresponds to atoms in energy levels of the harmonic trap that are sparsely populated where FD statistics are not important. The correction to J_2 from the data that we "left out" is calculated from $\sigma_{ex,y}$ and $\sigma_{ex,z}$ according to the equation for J_2 , where the integrals are only done outside of the cut region. From comparison to generated theory images, we find that j=2.5 to 3.5 produces only small errors (< 1%) in the measured value of E for our range of T/T_F . The best value of j depends on the signal to noise in the images.

5.7.2 Fitting Method

We can also extract E using a fit to the image with a modified criteria for the fit χ^2 . Using a fitting routine to determine E is slightly more robust and less complicated than the direct method. We use a simple gaussian fit with standard parameters included for the background plane. The cutting procedure is used as outlined in the above subsection, except the fit only uses data *inside* the cut region. The center of the cloud (yc and zc) and the background plane parameters $(b, m_y, \text{ and } m_z)$ are fixed to the values from the TF fit. This must be done because this "energy" fit does not have the information from the wings of the image that allow it to determine the cloud center and image background.

For the fit χ^2 , which the fitting routine (we use a two-dimensional Levenberg-Marquardt routine) attempts to minimize, we use:

$$\chi^2 = \sum_{y,z} \left[OD(y,z) - f(y,z) \right]^2 (y^2 + \alpha^2 z^2)^2$$
(5.31)

where the sum runs over all the pixels in the image, f is the value of the fit, and $\alpha = \sigma_z/\sigma_y$ is the expanded aspect ratio.¹¹ This criteria forces the fit to minimize the deviation of the second moment of the fit gaussian from the second moment of the image. Using the definition of J_2 , E is given by $m\omega_r^2 \left[2\frac{\sigma_{e,y}^2}{1+(\omega_r t)^2} + \frac{\lambda^2 \sigma_{e,z}^2}{1+(\omega_z t)^2}\right]$ where σ_e are

¹¹ Note that the expanded aspect ratio is calculated "on the fly" using the widths in the fitting routine.

the gaussian widths from the "energy" fit.

We find from applying this fit to theoretical images that j = 2.5 - 3.5 makes a negligible error in E even down to $T/T_F = 0.1$. Experimentally, j = 3 gives the best signal to noise for clouds with peak OD's between 0.35 and 1.3. We have also tested this fit on theoretical images of the thermal component of a trapped Bose gas above and below T_c . We find that this method works equally well for Bose, Fermi and, of course, classical gases.

5.8 Sources of Uncertainty

5.8.1 Shot-to-Shot Reproducibility

There is some scatter in N, T, the spin composition, and T/T_F for repeated, identical runs of the experiment. We believe that the major contributors to this irreproducibility are changes in the initial number of atoms loaded into the magnetic trap and changes in the magnetic trap bias field. Changes in the initial number directly affect the final number for a fixed evaporative cut, and may also affect the evaporation performance and therefore the final temperature. Changes in the magnetic trap bias field result in an effective change in the final evaporative cut (affecting both T and N) for a fixed microwave frequency.

We have studied the reproducibility of the experiment at very low temperature. The data that follows was taken (on August 12, 2000) after the "evap2" stage of cooling, with a fixed final cut of 1272.69 MHz. The anti-gravity coil was used to separate the components and extract information on each (using the "mixture" fits) at a 10 ms expansion time for 10 experimental cycles. For this data, the average spin mixture was $44\% m_f = 9/2$ atoms, with a standard deviation of 2% (in the value — not fractional). In the raw data, the average values and standard deviations were 1.8(1) for the peak OD, 11.5(2) pixels for σ_y , and 8.5(1) pixels for σ_z for the $m_f = 9/2$ component, and 2.11(7) for the peak OD, 12.4(2) pixels for σ_y , and 8.6(1) pixels for σ_z for the $m_f = 7/2$ component. This translates into 2.2(2) ×10⁵ for N, 109(3) nK for T, and 0.292(9) for T/T_F for the $m_f = 9/2$ component and 2.8(1)×10⁵ for N for the $m_f = 7/2$ component. This level of reproducibility (essentially 8% in N and 3 nK in T) is the best that we have seen. Typically, we have 10% shot-to-shot reproducibility in N, 5% fractional reproducibility in the spin composition, and 8 nK reproducibility in T (these values represent the standard deviation for repeated cycles of the experiment). Note that the uncertainty from the fitting routine in the fit parameters is typically less than 0.5%.

5.8.2 Systematic Uncertainties

Our ability to determine N, T, and T/T_F is systematically affected by uncertainty in the experimental parameters. Note that we assign a 50% systematic uncertainty to N determined from absorption imaging because absorption imaging and flourescence imaging only agree to this level. The primary source of this uncertainty is the lack of a measurement of the intensity of the light experienced by the atoms for flourescence imaging (using the science MOT beams).

The dominant contributors to systematic uncertainty are the uncertainty in the imaging magnification, U/D imaging OD correction, anti-gravity corrections, and the 50% systematic uncertainty in N. We have 0.5% uncertainty in the measurement of the imaging magnification (this affects N and T). To estimate the uncertainty in the U/D imaging OD correction and AG corrections, we use the standard deviation for all measurements. There is a 3.4% standard deviation across all measurements of the U/D imaging OD correction (which affects N), and a 3.3% standard deviation in measurements of the anti-gravity corrections to the widths from the fits (which affects T and N). Using standard uncertainty analysis, the systematic uncertainty in N is then 50%, in T is 7%, and in T/T_F is 17%. Obviously, the uncertainty in N and T/T_F is dominated by the 50% uncertainty in N.

5.9 Experimental Results

5.9.1 Deviation from Classical Momentum Profile — "Mixture" Fits

In figure 5.7 we show the first data that we acquired that indicated degeneracy [53]. This data was taken at a 20 ms expansion with a spin polarized $m_f = 9/2$ gas. The expansion time was sufficiently long so that we were probing the shape of the momentum distribution.¹² As T decreases below T_F , we see a clear deviation from a classical (L=1) momentum distribution. This nongaussian character of the images can also be seen in an analysis of fit residuals that uses azimuthally averaged data. When low T/T_F images are fit to the classical gaussian distribution, a pattern appears in the fit residuals as a function of the scaled radius ρ . This pattern is consistent with our expectation from fitting theoretical distributions (from equation 5.21) to gaussians. At these low temperatures, fits to the mixture function typically give a factor of 3 improvement in the reduced χ^2 compared with a simple gaussian fit.

¹² This is an excellent approximation for the radial direction, as $(\omega_r \cdot 20 \text{ ms})^2 = 288 \gg 1$. In the axial direction, this approximation is not as good: $(\omega_z \cdot 20 \text{ ms})^2 = 6$.



Figure 5.7: Original data demonstrating the emergence of quantum degeneracy as seen in the shape of the momentum distribution. Surface fits to the images reveal the nongaussian character of the momentum distribution at low T/T_F . For a particularly low noise image at $T/T_F = 0.5$ the inset shows fit residuals normalized by the peak OD vs. the scaled cloud radius ρ . A classical gaussian fit (solid points) is contrasted with the mixture fitting function from (hollow points). The main figure shows the fit parameter L of the non-gaussian form vs. T/T_F . For a classical gas L = 1; for a Fermi gas at $T/T_F = 0, L = 0$. The data compare well to theory (line) generated by fitting simulated TF distributions to the mixture function.

5.9.2 Measurements of the Fugacity — Experimental Results with the TF Fits

Compiled data showing the measured fugacity vs. T/T_F are shown in figure 5.8. The data in 5.8 are not typical — usually we see a small shift in the T/T_F axis (as in figure 5.7) that we attribute to a systematic error in determining T or N. The fugacity is determined *independently* from N and T, and can be used as a " T/T_F -meter".

Figure 5.9 shows a comparison between the TF and mixture thermometry by analyzing the data in figure 5.7 using both fits. Either technique works equally well. In practice, we use a fit to the TF profile to determine T, a gaussian fit to determine N, and then calculate T/T_F . The TF fit provides a very clean method for determining Tthat requires no corrections, unlike the "mixture" fit.



Figure 5.8: Measured fugacity using the TF thermometry. The solid line shows the expectation from thermodynamic theory. The fugacity \mathfrak{Z} is independent in the fits from N and T, which are used to calculate T/T_F in this figure.



Figure 5.9: Comparison between mixture and TF fits using the data in 5.7. The solid line is a linear fit with no intercept, and finds a slope 0.985(3).

5.9.3 Measured Momentum Profiles

The effects of FD statistics can be seen directly in measured momentum profiles of the gas. The OD in the two-dimensional image is averaged at constant scaled radius ρ for different values of T/T_F and displayed in figure 5.10. The data is compared with the profile for a classical gas with the same N and T and with the TF profile generated with the fit parameters.¹³

At $T/T_F \sim 1$, the data is consistent with a classical distribution. At lower T/T_F , the profile clearly becomes non-gaussian (the curvature of the distribution is too "flat" in the center of the cloud) and has suppressed density and higher r.m.s. momentum compared to the classical case. At $T/T_F = 0.2$, the peak density of the gas is suppressed by a little over a factor of two. This suppression is caused by the Fermi pressure, which resists the compression of the magnetic trap. In a semi-classical picture, the low momentum states are all occupied at low T/T_F . The Pauli exclusion principle forbids the occupancy of these states from being as high as it could classically. It is this Fermi pressure which is responsible for stabilizing neutron stars against gravitational collapse.

¹³ Note that the OD in these averaged profiles must be weighted by ρ in order to calculate N.



Figure 5.10: Momentum profiles of a $m_f = 9/2$ gas at different values of T/T_F . From top to bottom, $N = 2 \times 10^6$, $N = 1.7 \times 10^6$ and $N = 4 \times 10^5$ and T = 700 nK, T = 350 nK, and T = 90 nK.

5.9.4 Measurements of Energy

The data showing deviation in E from $3k_bT$ that we originally published in *Science* is shown in figure 5.11. We used the direct method to analyze this data. This method is fairly sensitive, and can detect deviations from the classical expectation as small as 3-5%. Data taken and analyzed with the fitting method are shown in figure 5.12 for partially spin polarized gases. Data for the $m_f = 9/2$ and $m_f = 7/2$ component are shown together in figure 5.12. For our lowest temperatures $T/T_F \sim 0.17$, we observe almost twice the classical energy per particle in the gas. The data in figure 5.12 was taken with the AG coil and used the image subtraction method described in the appendix to this chapter.

At low T/T_F , the "excess" energy compared to the classical expectation appears because of the Pauli exclusion principle. The lowest energy levels of the harmonic potential become highly occupied at low T/T_F . At $T/T_F = 0.17$, for example, the ground state of the trap is 99.5% occupied. Classically, the atoms would all fall into the low energy levels of the potential at low T. However, the Pauli exclusion principle prevents this, resulting in the appearance of "excess" energy for fermionic atoms. Like the electrons in an atom, the atoms in the trapped gas "stack up" one per quantum state in the energy levels of the harmonic potential.



Figure 5.11: Original data showing deviation in the energy of the gas from the classical expectation. We plot $\delta E/E_{cl} = (E - 3k_bT)/(3k_bT)$ vs T/T_F . *E* is determined using the "direct method" outlined in the text. This data was taken with spin polarized $m_f = 9/2$ gases.



Figure 5.12: Energy in the gas as a function of T/T_F . The extra energy in the gas compared to the classical case is plotted vs. T/T_F . This data was analyzed using the "fitting" method. Data taken using partially spin polarized mixtures for both the $m_f =$ 9/2 and $m_f = 7/2$ components are shown together. The thermodynamic expectation is shown by the solid line, while the classical expectation is shown by the dashed line. The data in this figure is the same as in figure 6.1.

5.10 Appendix

5.10.1 Image Subtraction

Using the methods to determine E in this chapter for images of spin polarized clouds is straightforward. However, images of spin mixed gases taken with the AG coil introduce a new wrinkle. In order to fit to the TF profile and to extract E using the energy fitting method, we must use the simultaneous image of both components to create separate images of each. It turns out that the mixture fit is the perfect tool for this task. The mixture fit can fit almost any individual image well (even images with a peak OD that is equal to the maximum observable OD) if all the parameters are allowed to vary freely. Also, fitting to two "mixture" functions added together is relatively fast.

We fit the double image to a sum of two mixture functions (a "double mixture" function), with shared parameters for the image background plane. A "double" FD fit, although it should work as well as a "double mixture" fit, is not used because it would be very computationally intensive (it would require evaluating the hyper-geometric function twice for each pixel). The "double mixture" fit is then used to subtract off either component and produce two images, one of each component separately. This procedure works rather well — no effect in the residuals in the processed single image can be detected. In addition, there is no difference between the energy or TF fit parameters between carefully matched images at the same T/T_F taken with one or both components present. This method will fail, however, when the separation between the centers of the components becomes small compared to their sizes.

This image subtraction method can even be used when one component has a peak OD that is equal to the maximum observable OD. In this case, the double mixture fit is restricted to fit pixels with an OD lower than 1.5. This fit is then used to subtract off the high OD component where only pixels with an OD less than 1.5 are processed. Fits to the resulting image are constrained to ignore any high OD pixels that are left behind. The TF and energy fits are then performed on the processed image of the low OD component.

Chapter 6

A TWO-COMPONENT DEGENERATE FERMI GAS

6.1 Introduction and Overview

This chapter will present the first experiments [32] with a two-component degenerate Fermi gas of atoms, starting with thermodynamic measurements. By manipulating the spin mixture, we change the value of E_F of one component relative to the other. We observe that we can create an imbalance in the mean energy per particle between the two gases in this way. In the degenerate regime, this is thermodynamically expected since E is not only a function of T but also T/T_F . Although the two gases are in thermal equilibrium, T/T_F is different for the two components if there are unequal numbers of atoms in each spin state. This same physics is observed in other systems — nuclei, for example. However, in nuclei the two components (protons and neutrons) can interconvert. An imbalance in the relative Fermi energies in the nucleus can therefore drive beta or inverse beta decay [112]. In particular, ⁴⁰K decays to ⁴⁰Ca with a 1.3 billion year half-life via this mechanism (see the online ENSDF database at http://ie.lbl.gov, for example).

In this chapter I will also present the direct observation of Pauli blocking of collisions through measurements of cross-dimensional rethermalization rates. The effect of Pauli blocking can be viewed as reducing the effective collision cross-section between colliding atoms. Using techniques similar to those explained in Chapter 3, we measure a reduction in the effective collision cross-section of a factor of two (compared to the classical case) at our lowest temperatures. Note that this effect is fundamentally different from the effect of the Pauli exclusion principle on collisions between indistinguishable fermions presented in Chapter 3. In that case, individual collisions are affected by the FD statistics of the particular colliding atoms. For the measurement presented in this chapter, individual collisions are affected by the FD statistics of the entire gas.

6.2 Two-Component Thermodynamics

The emergence of quantum degeneracy at low T is observed through measurements of the equilibrium thermodynamic properties (see Chapter 5) of the two-component gas. In the quantum degenerate regime $(T/T_F < 1)$, the average energy per particle E rises above the classical expectation $3k_BT$. For the roughly equal (46% $m_f=9/2$ for this data) mixture of spin states used for evaporation, figure 6.1(a) shows E vs T/T_F for each component. The excess energy characteristic of quantum degenerate Fermi systems can clearly be observed in both components. For this data, T and Eare determined independently for each component from fits to absorption images of the expanded gas. A fit to the Thomas-Fermi shape expected for an ideal Fermi gas is used to measure T, while E is determined from a gaussian fit that is weighted to minimize the fit deviation from the second moment of the image (see Chapter 5 for a discussion of both techniques). The widths of both fits are adjusted by roughly 6% to account for distortions introduced by curvature in the Stern-Gerlach field (see Chapter 2). Recall that these distortions cause the TF fit to return an incorrect fugacity (see chapter 2 for more details). The measured temperatures of the two components match to within experimental uncertainty, as expected for thermal equilibrium.


Figure 6.1: Thermodynamics of the interacting gas. The average energy per particle E, extracted from absorption images such as the examples shown in the insets, is displayed for two spin mixtures, 46% $m_f=9/2$ (a) and 86% $m_f=9/2$ (b). In the quantum degenerate regime, the data deviate from the classical expectation (dashed line) as the atoms form a Fermi sea arrangement in the energy levels of the harmonic trapping potential. The data in (a) represent the spin mixture used for evaporation, where we reach $T/T_F \sim 0.25$ at 90 nK and $N = 2.8 \times 10^5$ atoms. The data agree with the ideal Fermi gas prediction for a harmonic trap, shown by the solid line. Misalignment of corresponding $m_f=9/2$ and $m_f=7/2$ points on the T/T_F axis reflects a difference in the Fermi energies for the two components.

Thermodynamic data for a different spin mixture, 86% $m_f=9/2$ and 14% $m_f=7/2$, are displayed in 6.1(b). Here the spin composition is controlled by removing some $m_f=7/2$ atoms after the bulk of the evaporation. Changing the spin mixture manipulates the Fermi energies E_F since E_F depends on the number of atoms N through $E_F = k_B T_F = \hbar \overline{\omega} (6N)^{1/3}$ [27], where $\overline{\omega} = (\omega_r^2 \omega_z)^{1/3}$ is the geometric mean of the harmonic trap frequencies.¹ In the thermodynamics for this less balanced spin mixture, again both components reach quantum degeneracy.

For both mixtures, the thermodynamic data agree well with the ideal gas theory prediction; this indicates that the mean-field energy due to inter-particle interactions in the gas must be small compared to the kinetic and potential energy. For Bose-Einstein condensates (BEC) with similar number, temperature, and scattering lengths, the mean field is quite significant. However, the influence of the mean field for a Fermi gas is drastically reduced because the Fermi gas always has higher energy and lower density than a Bose-Einstein condensate. In fact, for our range of experimental conditions $E_{int}/k_BT_F < 0.4\%$, where E_{int} is the interaction (mean field) energy per particle [59].

Also apparent in 6.1 is a misalignment of the corresponding $m_f=9/2$ and $m_f=7/2$ points on the T/T_F axis, indicating that the two components are not equally degenerate. This is particularly true for the 86% $m_f=9/2$ case where E_F is roughly twice as high for the $m_f=9/2$ component, and therefore the $m_f=9/2$ component is always more degenerate. Figure 6.2 displays the effect of unequal E_F by plotting the energy ratio $E_{9/2}/E_{7/2}$ vs $T/T_{F,9/2}$ (T/T_F for the $m_f=9/2$ component). For the gas with 86% $m_f=9/2$ atoms, $E_{9/2}/E_{7/2}$ is measured as high as 1.4 in the quantum degenerate regime, strongly violating the classical expectation $E_{9/2}/E_{7/2} = 1$. However, when the gas has roughly equal numbers of $m_f=9/2$ and $m_f=7/2$ atoms E_F is matched to within 13%,

¹ The two spin-states have different magnetic moments and therefore experience slightly different harmonic oscillator frequencies. The magnetic trap frequencies for all of the experiments described in this chapter are 134.0 Hz radial and 19.9 Hz axial for the $m_f=9/2$ atoms. The corresponding frequencies for the $m_f=7/2$ atoms are reduced by the square root of the ratio of the magnetic moments ($\sqrt{7/9}$).



and E is roughly equal for both components irrespective of T/T_F .

Figure 6.2: Effect of Pauli blocking on the equilibrium thermodynamics of the gas. Using the same data shown in figure 6.1, the ratio of energy, $E_{9/2}/E_{7/2}$, for pairs of clouds from each double image is plotted vs. $T/T_{F, 9/2}$. Each point in this plot represents the average of two runs of the experiment. For comparison, the prediction for an ideal Fermi gas is shown by the solid lines. The energy imbalance revealed at low T/T_F is maintained by Pauli blocking of collisions.

The observed imbalance in E must arise from a change in the collisional interactions in the gas since s-wave collisions would normally redistribute energy equally between the two components. Collisions are predicted to be suppressed by Pauli blocking, a phenomenon common to all Fermi systems such as semiconductors, liquid ³He, and nuclear matter. At low T/T_F , the lowest energy states of the trap are highly occupied and any collision resulting in a final atom state at low energy is suppressed by the Pauli exclusion principle. The energy imbalance is then maintained because collisions that remove energy from the more degenerate component are the most strongly suppressed.

6.3 Two-Component Dynamics

We have directly observed Pauli blocking of elastic collisions in measurements of the thermal equilibration time. The gas is taken out of thermal equilibrium by a rapid removal of high energy atoms from the $m_f=7/2$ component (see the appendix to this chapter). Because of gravitational sag in the trap, energy is preferentially removed from the radial direction. As a result, the $m_f=7/2$ component is initially both out of crossdimensional equilibrium as well as out of equilibrium with the $m_f=9/2$ component. A sample data set showing the rethermalization of the energy of the $m_f=7/2$ component is shown in figure 6.3. The cross-dimensional relaxation rate can be used to obtain the elastic collision cross-section (see Chapter 3 and [92]). One can define an *effective* collision cross-section σ_{eff} that encapsulates the total effect of Pauli blocking on collisions independent of changes in the density and temperature of the gas. The value of σ_{eff} is predicted to vary from 0 at T=0 to the s-wave cross-section $\sigma = 4\pi a^2$ in the classical regime [37], where a is the s-wave triplet scattering length.



Figure 6.3: Rethermalization data. A plot of the aspect ratio (axial size/radial size) of the expanded $m_f=7/2$ component vs. time shows cross-dimensional energy rethermalization. The inset, a plot of the average energy per particle of the $m_f=7/2$ component, shows the simultaneous transfer of energy from the $m_f=9/2$ to $m_f=7/2$ component. The data in this figure were taken for a gas with $T/T_{F, 9/2}=0.5$, $N_{9/2}=4.6 \times 10^5$, and $N_{7/2}=7.7 \times 10^4$. A fit (solid lines) of the time dependent aspect ratio is used to measure the effective collision cross-section.

Figure 6.4 shows the measured σ_{eff} vs $T/T_{F, 9/2}$ for an 86% $m_f=9/2$ gas. The rethermalization time constant τ is extracted from a fit to the time dependence of the $m_f=7/2$ component aspect ratio, assuming that the energy difference $\delta = E_x - E_z$ relaxes exponentially (E_z and E_x refer to the $m_f=7/2$ energy in the axial and one of the radial directions, respectively). The cross-section is then determined through $\frac{1}{\tau} = \frac{n\sigma_{eff}v}{\alpha}$, where $n = \frac{1}{N_{7/2}} \int d^3\vec{r} \ n_{9/2}(\vec{r})n_{7/2}(\vec{r})$ is the density-weighted-density, v is the mean relative speed for a collision between $m_f=9/2$ and $m_f=7/2$ atoms, and α is the average number of collisions per $m_f=7/2$ atom required for cross-dimensional equilibration. The product nv is determined from gaussian fits to the expanded images of each component (see the appendix to this chapter). See the appendix to this chapter for a complete explanation of the rethermalization model.



Figure 6.4: Collisional Pauli blocking. A factor of two decrease in the effective elastic collision cross-section, σ_{eff} , is observed at low T/T_F . The error bars in σ_{eff} are predominately from uncertainty in the fits to the time dependence of the $m_f = 7/2$ aspect ratio, while scatter in number and temperature set the error bars in $T/T_{F, 9/2}$. In addition, there is at most a 20% systematic uncertainty in $T/T_{F, 9/2}$ and a 50% systematic uncertainty in σ_{eff} from uncertainty in the number determined from absorption imaging. The solid line shows the result from a quantum kinetic calculation of the collision rate. At high T/T_F the data agree with the known value of the s-wave collision cross-section, and at low T/T_F the observed decrease in σ_{eff} agrees with the quantum kinetic prediction.

For our data, the time dependence of the aspect ratio is complicated by energy transfer from the $m_f=9/2$ component to the $m_f=7/2$ component. This is included in the fit to the aspect ratio by assuming that the difference $\Delta = E_{9/2} - \eta E_{7/2}$ also relaxes exponentially with time constant $\left(\frac{n\sigma_{\text{eff}}v}{A}\frac{N_{7/2}+\eta N_{9/2}}{N_{9/2}(1+\eta)}\right)^{-1}$, where η is the equilibrium ratio $E_{9/2}/E_{7/2}$, and A is a constant similar to α . We have used a classical kinetic theory calculation to determine that $\alpha = A = 0.75$ (see the appendix to this chapter).

In the classical regime, the measured value of σ_{eff} agrees, within the 50% systematic uncertainty in atom number, with the best known value of the scattering length for 40 K $a = 169a_0$ [89], where a_0 is the Bohr radius. The effective cross-section drops by a factor of two at $T/T_{F, 9/2} = 0.2$. Within our uncertainty, the observed size of the Pauli blocking effect agrees with the theoretical value of σ_{eff} from a quantum kinetic calculation [54] shown in figure 6.4. In the calculation, the reduction in σ_{eff} represents the effect of Pauli blocking averaged over all possible initial and final colliding atom states.

6.4 Chapter and Thesis Conclusion

This thesis has presented the production of and measurements with the world's first degenerate Fermi gas of atoms. We developed a novel cooling scheme in order to cool a gas of ⁴⁰K atoms to well below to Fermi temperature. New techniques of thermometry were devised in order to observe the emergence of quantum degeneracy through measurements of energy and the momentum distribution. After spending a year exploring and improving the performance of evaporative cooling in the quantum regime, we were able to produce the world's first and currently only two-component Fermi gas of atoms. The work in this chapter details the investigation of the effect of interactions on the thermodynamics and dynamics of the two-component degenerate gas.

As described in the introduction (chapter 1), there are many experiments left to

do with an atomic DFG. The work detailed in this thesis makes the realization of the equivalent of BCS-like pairing [67,113,40] in an atomic gas possible in the near future. The possibility of the phase transition to a paired state requires an attractive interaction between the atoms. Although the interaction for the states we use, $m_f = 9/2$ and $m_f = 7/2$, is repulsive, theoretical calculations predict [69] an experimentally accessible Feshbach resonance for the $m_f = -9/2$ and $m_f = -7/2$ states. A degenerate gas of these states could be created by loading the gas of the positive m_f states into an optical trap and quickly reversing the direction of the magnetic field. Then, the inter-atomic interactions could be made *strongly* attractive by tuning the magnetic field magnitude. In fact, unpublished work by Murray Holland predicts that a phase transition to a paired state at $T/T_F = 0.5$ is possible using the Feshbach resonance in this way.

6.5 Appendix

6.5.1 Rethermalization Model

6.5.1.1 Overview

This appendix is not a rigorous mathematical derivation of a rethermalization model. Rather, I will informally explain the basics of the model that we use to analyze the data presented in this chapter. This appendix follows the notation used in *Statistical* and *Thermal Physics* by Reif [114], which is an excellent reference for kinetic theory. The material in this appendix is meant to be used in conjuction with [114], and may be difficult to follow without some cross-reference.

Physical intuition suggests that the rethermalization rate for quantities like δ and Δ (the difference in energy between the directions of the $m_f=7/2$ component and between the two components, respectively) should be connected to the collision rate per particle in the gas. One might guess that the rethermalization should be exponential in time, as well. In fact, a Monte Carlo simulation by Murray Holland for a single component, classical gas found that cross-dimensional relaxation is nearly exponential in time, with a rethermalization rate that is proportional to the collision rate per particle [42]. However, this is difficult to justify *a priori* for a two-component Fermi gas. All we know for sure is that the time rate of change of Δ and δ must be zero in thermal equilibrium.

We use the kinetic theory derived in [114] to explore the time dependence of Δ and δ . In [114], Enskog's equations are derived from the first principles of kinetic theory. These equations connect the time rate of change of any average quantity in the gas to a collision integral.² For our case, the kinetic theory describes how collisions rethermalize energy between the components and between the directions in the $m_f = 7/2$ component.

Unfortunately, including the Pauli blocking explicitly in the kinetic theory equations is difficult. We solve the equations for the classical case, and then "patch" the results for the quantum regime. The patch is designed to take into account the effects of FD statistics that we do not want to probe (thermodynamics) and lump what we do wish to probe (Pauli blocking) into σ_{eff} . Numerical calculations are used to justify approximations that are made in order to produce tractable (and analytic) rethermalization equations. In this appendix the label "1" is used to refer to the $m_f = 9/2$ component and "2" to the $m_f = 7/2$ component.

In this appendix, we will assume that the gas maintains a quasi-equilibrium at all times. We assume that the statistical distribution function can be approximated by a FD distribution with a well defined fugacity $\mathbf{3}$ and mean energy per particle E. Ergodicity is broken by allowing different values of E in the radial and axial directions in the gas (see the appendix to chapter 3). However, we do assume that the x and ydirections in the gas stay in thermal equilibrium because of a separation of the x-yequilibration timescale and the timescale for the microwave frequency sweep that is used

² Because the force from the harmonic trap is conservative and the gas has no net momentum, only the collision term in Enskog's equation drives changes in δ and Δ .

to bring the $m_f = 7/2$ component of out equilibrium. The equilibration timescale for the x and y directions should on the order of the inverse of the radial trap frequency, ~ 8 ms. The microwave frequency sweep that takes the $m_f = 7/2$ component out of equilibrium between the radial and axial direction takes ~ 150 ms. In the experiment we do not observe any rapid change in the properties of the gas that would reflect an x-y re-equilibration after the microwave cut.

Under these assumptions, we characterize the gas by $E = E_x + E_y + E_z = 2E_y + E_z$. In addition, it is assumed that the energy is evenly distributed (on average) between kinetic and potential energy in each direction: $E_y = \frac{1}{2}m\omega_r^2 \langle y^2 \rangle + \frac{1}{2}m \langle v_y^2 \rangle$, for example (the averaging symbols refer to a statistical average over the gas).

6.5.1.2 Time Dependence of Δ

We first consider the time dependence of $\Delta = E_1 - \eta E_2$, which is the difference in energy between the components with an adjustment so that $d\Delta/dt = 0$ in thermal equilibrium in the quantum regime. The adjustment factor, η , is the equilibrium ratio E_1/E_2 .

From [114], the kinetic equation for Δ is

$$\frac{1}{2}\frac{d\Delta}{dt} = \int d^3\vec{p}_1 \int d^3\vec{p}_2 \int d\Omega_f \frac{|\vec{p}_1 - \vec{p}_2|}{m} n_{12} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2} \frac{d\sigma_{eff}}{d\Omega_f} \times \frac{1}{2m} \left\{ \frac{(p_1\,')^2}{N_1} - \eta \frac{(p_2\,')^2}{N_2} - \left[\frac{(p_1)^2}{N_1} - \eta \frac{(p_2)^2}{N_2} \right] \right\}$$
(6.1)

Here, the change in Δ is driven by collisions between a 1 and a 2 atom, with initial momenta $\vec{p_1}$ and $\vec{p_2}$ and final momenta $\vec{p_1}'$ and $\vec{p_2}'$. The final momenta are implicitly a function of the initial momenta and the solid angle Ω_f between the initial and final relative momenta. The statistical momentum distributions $\Pi(\vec{p})$ (which will be addressed later in this appendix) are normalized so that $\int d^3\vec{p_i} \ \Pi_i(\vec{p_i}) = N_i$ for i = 1 or 2. The density overlap integral is given by

$$n_{12} = \int d^3 \vec{r} \, n_1(\vec{r}) \, n_2(\vec{r}) \tag{6.2}$$

with a similar normalization $(\int d^3 \vec{r} n_i(\vec{r}) = N_i)$ for the statistical density distribution $n(\vec{r})$ (which will be addressed later in this appendix). Note that equation 6.1 is missing a factor of $\frac{1}{2}$ on the right side of the equation compared to [114] because the colliding atoms are not identical.

The physics behind equation 6.1 can be understood by breaking the equation into pieces. The time rate of change of Δ (as an average quantity) is the change in Δ for any particular collision (the second line of the equation) averaged over all possible collisions in the gas. Note that only the kinetic energy of a colliding atom is changed by a collision, so that the potential energy does not appear in equation 6.1. This is also the reason for the factor of $\frac{1}{2}$ on the left side of the equation — we are only considering the time dependence of the kinetic half of Δ . The integral over $d\Omega_f$ covers the average over all final colliding states, while the integral over $d^3\vec{p}_1d^3\vec{p}_2$ covers the average over all initial colliding states. The factor $\frac{|\vec{p}_1-\vec{p}_2|}{m}n_{12}\Pi(\vec{p}_1)\Pi(\vec{p}_2)\frac{d\sigma_{eff}}{d\Omega_f}$, combined with the differentials, gives the rate at which a collision with specific parameters $\vec{p}_1', \vec{p}_2', \vec{p}_1, \vec{p}_2$ occurs in the gas.

For the time being, we consider the classical case and take $\eta \to 1$ and $\frac{d\sigma_{eff}}{d\Omega_f} = \frac{\sigma_{eff}}{4\pi}$. The integral over $d\Omega_f$ can be done analytically by using center of mass and relative coordinates. We find that the kinetic equation 6.1 (still with $\eta = 1$) with this integration is:

$$\frac{1}{2}\frac{d\Delta}{dt} = -n_{12}\frac{N_1 + N_2}{2N_1N_2}\sigma_{eff}\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{|\vec{p}_1 - \vec{p}_2|}{m}\frac{\Pi_1(\vec{p}_1)}{N_1}\frac{\Pi_2(\vec{p}_2)}{N_2}\frac{1}{2m}\left[(p_1)^2 - (p_2)^2\right]$$
(6.3)

We now define A by

$$A = \frac{\left\{\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{|\vec{p}_1 - \vec{p}_2|}{m} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2}\right\} \times \left\{\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{1}{2m} \left[(p_1)^2 - (p_2)^2\right] \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2}\right\}}{\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{|\vec{p}_1 - \vec{p}_2|}{m} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2} \frac{1}{2m} \left[(p_1)^2 - (p_2)^2\right]}{(6.4)}$$

and rewrite equation 6.3 as:

$$\frac{1}{2}\frac{d\Delta}{dt} = -n_{12}\frac{N_1 + N_2}{2N_1N_2}\,\sigma_{eff}\frac{1}{A}\left\{\int\!d^3\vec{p_1}\int\!d^3\vec{p_2}\frac{|\vec{p_1} - \vec{p_2}|}{m}\frac{\Pi_1(\vec{p_1})}{N_1}\frac{\Pi_2(\vec{p_2})}{N_2}\right\}\times$$

$$\left\{ \int d^3 \vec{p}_1 \int d^3 \vec{p}_2 \frac{1}{2m} \left[(p_1)^2 - (p_2)^2 \right] \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2} \right\}$$
(6.5)

$$= -n_{12} \frac{N_1 + N_2}{2N_1 N_2} \sigma_{eff} \frac{1}{A} \langle |\vec{v}_1 - \vec{v}_2| \rangle \langle \frac{1}{2m} \left[(p_1)^2 - (p_2)^2 \right] \rangle$$
(6.6)

where the angle brackets refer to a statistical average: $\langle \rangle \rightarrow \int d^3 \vec{p}_1 \int d^3 \vec{p}_2 \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2}$. The average of $|\vec{v}_1 - \vec{v}_2|$ is the mean relative speed between a "1" and a "2" atom, which we denote by v. The average of $\frac{1}{2m} [(p_1)^2 - (p_2)^2]$ is just $\frac{1}{2}\Delta$ (still with $\eta = 1$), so equation 6.6 turns into:

$$\frac{d\Delta}{dt} = -\frac{1}{A} \frac{N_1 + N_2}{2N_1 N_2} n_{12} \sigma_{eff} v \Delta$$
(6.7)

which gives exponential relaxation of Δ in time, assuming that A is a constant. We find that the exponential time constant is proportional to one over the average of the collision rate per particle in each component, $\left[\frac{1}{2}\left(\frac{n_{12}\sigma_{eff}v}{N_1} + \frac{n_{12}\sigma_{eff}v}{N_2}\right)\right]^{-1}$.

We now revisit the quantum case, for which $\eta \neq 1$. The effective differential cross-section $\frac{d\sigma_{eff}}{d\Omega_f}$ is properly a function of the final momenta since Pauli blocking (of the form $1 - n(\vec{p}')$) terms should be included explicitly in the kinetic equations. Pauli blocking is what allows the gas to reach an equilibrium where $E_1 \neq E_2$ (or $\eta \neq 1$). To "patch" the classical result (equation 6.7), we consider the average over the final states (over $d\Omega_f$) from a different perspective. The total energy in a collision must be divided up, on average, into the final states consistently with thermodynamic equilibrium. For thermodynamic equilibrium, $d\Delta/dt = 0$ and $\Delta = 0$. This constraint is satisfied if the total energy $e_t = e_1 + e_2$ in a collision divides up, on average, fractionally to the 1 and 2 atom after the collision according to $e'_1 = \frac{\eta}{1+\eta} e_t$ and $e'_2 = \frac{1}{1+\eta} e_t$. Here, e_1 and e_2 are the initial kinetic energies of atom 1 and 2 in the collision, and e'_1 and e'_2 are the final kinetic energies of atom 1 and 2 in the collision.

We denote the average over $d\Omega_f$ as $\langle \rangle_{d\Omega_f}$. On average,

$$\left\langle \frac{e_1'}{N_1} - \eta \frac{e_2'}{N_2} - \left(\frac{e_1}{N_1} - \eta \frac{e_2}{N_2}\right) \right\rangle_{\Omega_f} = -\frac{N_2 + \eta N_1}{N_1 N_2 (1+\eta)} \left(e_1 - \eta e_2\right)$$
(6.8)

must be true.³ Comparing to the classical result (equation 6.7), we guess that the time dependence of Δ is approximated by

$$\frac{d\Delta}{dt} = -\frac{1}{A} \frac{N_2 + \eta N_1}{N_1 N_2 (1+\eta)} n_{12} \sigma_{eff} v \Delta$$
(6.9)

The final answer for $\Delta(t)$ is:

$$\Delta(t) = (E_{10} - \eta E_{20}) e^{-t/\tau_{\Delta}}$$
(6.10)

where $\tau_{\Delta} = \left(\frac{1}{A}n_{12}\sigma_{eff}v\frac{N_2+\eta N_1}{N_1N_2(1+\eta)}\right)^{-1}$. The initial conditions are given by $E_{10} = E_1(t=0)$ and $E_{20} = E_2(t=0)$.

The value of A can be evaluated analytically if $n(\vec{v})$ is isotropic in velocity space and the distribution is gaussian (classical case). In this case, we find $A = \frac{3}{4}$, and Δ varies exponentially in time. In order to describe the experimental case, however, the value of A must be evaluated using FD distributions that allow different mean energies in the different directions. A distribution that is correctly normalized for N (so that $\int d^3\vec{p} \Pi(\vec{p}) = N$) is

$$\Pi(\vec{p}) = -\frac{1}{(2\pi)^{3/2}\hbar^3 \lambda \sqrt{\zeta}} \left(\frac{k_b T_r}{m\omega_r^2}\right)^{3/2} Li_{3/2} \left(-\Im e^{-\frac{p_r^2}{2mk_b T_r}} e^{-\frac{p_z^2}{2mk_b T_z}}\right)$$
(6.11)

where $\zeta = T_r/T_z$ and z is now a function of ζ . This distribution gives $\langle \frac{p_z^2}{2m} \rangle = k_b T_z \frac{Li_4(-3)}{Li_3(-3)}$ and $\langle \frac{p_r^2}{2m} \rangle = 2k_b T_r \frac{Li_4(-3)}{Li_3(-3)}$ (here the angle brackets refer to an average done as $\langle \frac{p_y^2}{2m} \rangle = \frac{1}{N} \int d^3 \vec{p} \Pi(\vec{p}) \frac{p_y^2}{2m}$, for example). The variables T_r and T_z should not be viewed as temperatures (since the gas is not in equilibrium), but rather convenient parameters to allow each direction to have a different mean kinetic energy.

The value of A can now be calculated using the explicit definition given previously (equation 6.4). We use the anisotropic FD distribution (equation 6.11) to numerically investigate the case for $T/T_F = 0.25$. The range $\zeta = 0.5$ (and $E_1/E_2 = 2$) to $\zeta = 1$ (and $E_1/E_2 = 1.3$) approximately covers the rethermalization data. For $\zeta = 0.5$ and

³ To get this equation, all we've done is substitute the expressions for e'_1 and e'_2 into the left side of the equation.

 $E_1/E_2 = 2$, we find that A = 1.3, and for $\zeta = 1$ and $E_1/E_2 = 1.3$ we find that A = 1.0. Since A is not constant during the equilibration at low T/T_F , the time dependence of Δ is only approximately exponential. For high T/T_F ($T/T_F = 2$), we find that $A \sim 0.75$ for the range $\zeta = 0.5$ to 1.0 (and $E_1/E_2 = 2$ to $E_1/E_2 = 1$). Numerically we determine that A ranging from 1.3 to 0.75 makes little difference ($\sim 10\%$ in the value we determine for σ_{eff}) in the fit that determines the rethermalization time constant as described in this chapter.

6.5.1.3 Time Dependence of δ

Next, we investigate the time dependence of $\delta = E_{2y} - E_{2z}$, which is a measure of the degree of cross-dimensional equilibrium of the "2" component. The analysis of the time dependence of δ closely follows the investigation of the time dependence of Δ from the previous section.

The kinetic equation (from [114]) for the time evolution of δ is:

$$\frac{1}{2}\frac{d\delta}{dt} = \frac{n_{12}}{N_2} \int d^3\vec{p}_1 \int d^3\vec{p}_2 \int d\Omega_f \frac{|\vec{p}_1 - \vec{p}_2|}{m} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2} \frac{d\sigma_{eff}}{d\Omega_f} \times \frac{1}{2m} \left\{ (p_{2y}')^2 - (p_{2z}')^2 - \left[(p_{2y})^2 - (p_{2z})^2 \right] \right\}$$
(6.12)

As for the kinetic equation for Δ , this equation can be viewed as an average over all possible collisions. In equation 6.12, $\vec{p_1}$ and $\vec{p_2}$ are the initial momenta of the 1 and 2 atom for a particular collision. For the "2" atom, p_{2y} and p_{2z} are the projections of the initial momentum onto the axes of the harmonic trap. The quantities p_{2y} and p_{2z} are similar quantities for the final momentum of the 2 atom. The final momenta are implicitly functions of the initial momenta and the solid angle Ω_f , as in the kinetic equation for Δ .

I will jump straight to the simplified final state averaging procedure for δ . The procedure works the same way as for Δ , except that we break up the kinetic collision energy e_1 and e_2 into directions. Here e_1 and e_2 are the initial total kinetic energy for the "1" and "2" atom in a collision. The total energy in the collision again breaks up, on average, fractionally into the 1 and 2 atoms after the collision as before. Now we assume that the energy is equally divided up after the collision, on average, into each direction: $e'_{2y} = \frac{1}{3}e'_2$ and $e'_{2z} = \frac{1}{3}e'_2$ and $e'_{1y} = \frac{1}{3}e'_1$ and $e'_{2z} = \frac{1}{3}e'_1$. Here the "primed" quantities refer to the kinetic energies of the 1 and 2 atoms after the collision in the y and z directions (recall that these directions refer to the axes of the harmonic trap). Equal division of the energy into the different directions guarantees that in thermal

equilibrium $E_y = E_z$. We find that $\langle e'_{2y} - e'_{2z} - (e_{2y} - e_{2z}) \rangle_{d\Omega_f} = -(e_{2y} - e_{2z}).^4$

After completing the angular average, equation 6.12 simplifies to:

$$\frac{1}{2}\frac{d\delta}{dt} = -\frac{n_{12}}{N_2}\sigma_{eff}\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{|\vec{p}_1 - \vec{p}_2|}{m} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2} \frac{1}{2m} \left[(p_{2y})^2 - (p_{2z})^2 \right]$$
(6.13)

Introducing α (similar to A) defined by:

$$\alpha = \frac{\left\{\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{|\vec{p}_1 - \vec{p}_2|}{m} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2}\right\} \times \left\{\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{1}{2m} \left[(p_{2y})^2 - (p_{2z})^2\right] \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2}\right\}}{\int d^3\vec{p}_1 \int d^3\vec{p}_2 \frac{|\vec{p}_1 - \vec{p}_2|}{m} \frac{\Pi_1(\vec{p}_1)}{N_1} \frac{\Pi_2(\vec{p}_2)}{N_2} \frac{1}{2m} \left[(p_{2y})^2 - (p_{2z})^2\right]}{(p_{2z})^2} \frac{(6.14)}{(p_{2z})^2}$$

we rewrite equation 6.13 as:

$$\begin{aligned} \frac{1}{2}\frac{d\delta}{dt} &= -\frac{n_{12}}{N_2}\,\sigma_{eff}\frac{1}{\alpha}\left\{\int\!\!d^3\vec{p}_1\int\!\!d^3\vec{p}_2\frac{\Pi_1(\vec{p}_1)}{N_1}\frac{\Pi_2(\vec{p}_2)}{N_2}\frac{1}{2m}\left[(p_{2y})^2 - (p_{2z})^2\right]\right\} \times \\ &\left\{\int\!\!d^3\vec{p}_1\int\!\!d^3\vec{p}_2\frac{|\vec{p}_1 - \vec{p}_2|}{m}\frac{\Pi_1(\vec{p}_1)}{N_1}\frac{\Pi_2(\vec{p}_2)}{N_2}\right\} \quad (6.15) \\ &= -\frac{n_{12}}{N_2}\sigma_{eff}\frac{1}{\alpha}\langle\frac{1}{2m}\left[(p_{2y})^2 - (p_{2z})^2\right]\rangle\langle|\vec{v}_1 - \vec{v}_2|\rangle \quad (6.16) \end{aligned}$$

where the angle brackets refer to a statistical average over $\vec{p_1}$ and $\vec{p_2}$ as in the analysis of Δ .

With the same definition for v as before and realizing that $\frac{1}{2}\delta = \langle \frac{1}{2m} \left[(p_{2y})^2 - (p_{2z})^2 \right] \rangle$, we have:

$$\frac{d\delta}{dt} = -\frac{1}{\alpha} \frac{n_{12}}{N_2} \sigma_{eff} v_{12} \delta \tag{6.17}$$

Again, we find that the time dependence of δ is exponential if α is constant, and that the rethermalization rate is proportional to the collision rate per particle in the "2" component, $\frac{n_{12}}{N_2}\sigma_{eff}v_{12}$. With initial conditions $E_{2y}(t=0) = E_{2y0}$ and $E_{2z}(t=0) = E_{2z0}$, we have:

$$\delta(t) = (E_{2y0} - E_{2z0}) e^{-t\frac{1}{\alpha}\frac{n_{12}}{N_2}\sigma_{eff}v_{12}}$$
(6.18)

Using the FD anisotropic momentum distributions from equation 6.11, numerical calculations indicate that α is fixed to 0.75 to within 8% across our range of experimental conditions. The model therefore predicts a nearly exponential time dependence for δ .

⁴ Again, to obtain this equation we've just substituted for e'_{2y} and e'_{2z} into the left side of the equation.

For the classical and quantum cases, α varies from ~ 0.69 to ~ 0.74 as ζ varies from 0.3 to 0.95 (numerically investigating $\zeta = 1$ presents some technical problems, but the calculations indicate that α approaches 0.75 as ζ approaches 1).⁵ This result agrees with the result of a Monte-Carlo simulation for a single component, classical gas [42]. Note that α and A can be viewed as the average number of collisions per particle to rethermalize δ and Δ , respectively.

6.5.1.4 Functional Form for the Aspect Ratio vs. Time

We wish to find E_{2z}/E_{2y} as a function of time, which is proportional to the expanded cloud's aspect ratio squared (see the appendix to chapter 3). If total energy in the gas is $U = N_1 E_1 + N_2 E_2$ then $\Delta = U/N_1 - (\beta + \eta) E_2$ giving $E_2 = \frac{1}{\beta + \eta} (U/N_1 - \Delta)$ where $\beta = N_2/N_1$. Using $E_2 = 2E_y + E_z$, $\delta = E_y - E_z$, and some algebra, we find that:

$$\frac{E_{2z}}{E_{2y}} = \frac{E_2 - 2\delta}{E_2 + \delta} \tag{6.19}$$

which gives for the aspect ratio

$$\sqrt{\gamma \frac{[U/N_1 - \Delta(t)] - 2\delta(t)(\beta + \eta)}{[U/N_1 - \Delta(t)] + \delta(t)(\beta + \eta)}}$$
(6.20)

with $\gamma = \frac{1+(\omega_{7z}t_{exp})^2}{1+(\omega_{7r}t_{exp})^2}\frac{\omega_{7r}^2}{\omega_{7z}^2}$, where ω_{7r} and ω_{7z} are the harmonic oscillator frequencies for the $m_f = 7/2$ component and t_{exp} is the expansion time. Substituting in and doing some more algebra, we get the function that we actually fit the aspect ratio vs. time to:

$$\sqrt{\gamma \frac{h+\beta-(h-\eta)e^{-\frac{t}{\tau}\frac{1}{A}\frac{N_2+\eta N_1}{N_1N_2(1+\eta)}}{h+\beta-(h-\eta)e^{-\frac{t}{\tau}\frac{1}{A}\frac{N_2+\eta N_1}{N_1N_2(1+\eta)}} - 2s(\beta+\eta)e^{-\frac{t}{\tau}\frac{1}{\alpha}\frac{1}{N_2}}}$$
(6.21)

where γ , $\tau = \frac{1}{n_{12}\sigma_{eff}v_{12}}$, and $s = \left(1 - \frac{E_{2z0}}{E_{2y0}}\right) / \left(2 + \frac{E_{2z0}}{E_{2y0}}\right)$ are left as free parameters in the fit. The measured values of N_1 , N_2 , β , η , and $h = E_{10}/E_{20}$ are fixed in the fit.

 $^{^5}$ To calculate α for the classical case, the classical limit of equation 6.11 is used.

6.5.1.5 Determination of n and v

After we use the fit to the aspect ratio vs. time to extract τ , we need the product $n_{12}v$ in order to determine σ_{eff} . Again, these parameters are defined as:

$$n_{12} = \int d^3 \vec{r} \, n_1(\vec{r}) \, n_2(\vec{r}) \tag{6.22}$$

$$v = \frac{1}{N_1 N_2} \int d^3 \vec{p}_1 \int d^3 \vec{p}_2 \ \Pi_1(\vec{p}_1) \ \Pi_2(\vec{p}_2) \frac{|\vec{p}_1 - \vec{p}_2|}{m}$$
(6.23)

where an FD anisotropic density distribution can be defined as:

$$n(\vec{r}) = -\frac{(k_b m T_r)^{3/2}}{\sqrt{\zeta} \hbar^3 (2\pi)^{3/2}} Li_{3/2} \left[-\Im e^{-\frac{m\omega_r^2}{2k_b T_r} \left(r^2 + \lambda^2 \zeta z^2\right)} \right]$$
(6.24)

which is normalized properly for $N\left(\int d^3\vec{r} n_i(\vec{r}) = N_i \text{ for } i = 1, 2\right)$ and gives $\langle \frac{1}{2}m\omega_r^2 r^2 \rangle = k_b T_r \frac{Li_4(-3)}{Li_3(-3)}$ and $\langle \frac{1}{2}m\omega_z^2 z^2 \rangle = \frac{1}{2}k_b T_z \frac{Li_4(-3)}{Li_3(-3)}$. Here, the statistical average is done as $\langle \frac{1}{2}m\omega_r^2 r^2 \rangle = \frac{1}{N} \int d^3\vec{r} n(\vec{r}) \left(\frac{1}{2}m\omega_r^2 r^2\right)$, for example. Again, because the gas is out of equilibrium, the parameters T_r and T_z should not be regarded as temperatures, but rather parameters that allow the different directions in the gas to have different mean potential energies.

In order to determine $n_{12}v$ for the data, we fit the images to gaussians. The product $n_{12}v$ is then calculated using anisotropic gaussian momentum and density distributions. Gaussian distributions that are the high T/T_F limit of equations 6.11 and 6.24 are:

$$n(\vec{r}) = \frac{\lambda N \sqrt{\zeta}}{(2\pi)^{3/2}} \left(\frac{m\omega_r^2}{k_b T_r}\right)^{3/2} e^{-\frac{m\omega_r^2}{2k_b T_r} \left(r^2 + \lambda^2 \zeta z^2\right)}$$
(6.25)

$$\Pi(\vec{p}) = \frac{N\sqrt{\zeta}}{(2\pi)^{3/2} (mk_b T_r)^{3/2}} e^{-\frac{1}{2mk_b T_r} \left(p_r^2 + \zeta p_z^2\right)}$$
(6.26)

These distributions give $\frac{1}{N} \int d^3 \vec{r} n(\vec{r}) \left(\frac{1}{2}m\omega_r^2 r^2\right) = k_b T_r$ and $\frac{1}{N} \int d^3 \vec{p} n(\vec{r}) \left(\frac{1}{2}m\omega_z^2 z^2\right) = \frac{1}{2}k_b T_z$ as average values of potential energy. Likewise, the average values of kinetic energy are: $\frac{1}{N} \int d^3 \vec{p} \Pi(\vec{p}) \left(\frac{1}{2m} p_r^2\right) = k_b T_r$ and $\frac{1}{N} \int d^3 \vec{p} \Pi(\vec{p}) \left(\frac{1}{2m} p_r^2\right) = \frac{1}{2}k_b T_z$.

The accuracy of using gaussian fits, rather than TF fits, is checked numerically. The gaussian fitting procedure is tested on images generated using the anisotropic FD density distributions (equation 6.24). Of course, in the classical regime the gaussian fits determine $n_{12}v$ exactly. At worst, at $T/T_F = 0.2$ and $\zeta = 0.5$, the gaussian fits make a 2% error in $n_{12}v$.

6.5.2 Fast $m_f = 7/2$ Component Removal

A single frequency sweep (as described in chapter 4) is used to quickly remove atoms from the $m_f = 7/2$ component and bring the gas out of thermal equilibrium. Generating the frequency sweep via GPIB commands to the HP E4420B synthesizer is too slow and coarse for this purpose. For this reason, we use the analog voltage controlled, phase-continuous sweep capability of this synthesizer. The frequency is swept by stepping an analog voltage (that is connected to the synthesizer using the driver circuit described in Chapter 1) in time, and uses the same exponential sweep parameters as described in the appendix to Chapter 4. Typical parameters for this sweep are: sweep rate of 1 Hz, sweep start of 1274 MHz, sweep end of 1274.057 MHz, sweep "bottom" at 1274.4 MHz. We optimize the sweep to provide some cooling but still bring the $m_f = 7/2$ component out of thermal equilibrium.

Bibliography

- M. H. Anderson, J. R. Ensher, M. R. Matthews, C. E. Wieman, and E. A. Cornell, Science 269, 198 (1995).
- [2] K. B. Davis, M.-O. Mewes, M. R. Andrews, N. J. van Druten, D. S. Durfee, D. M. Kurn, and W. Ketterle, Phys. Rev. Lett. 75, 3969 (1995).
- [3] C. C. Bradley, C. A. Sackett, J. J. Tollett, and R. G. Hulet, Phys. Rev. Lett. 75, 1687 (1995), *ibid.* 79, 1170 (1997).
- [4] E. L. Raab, M. Prentiss, A. Cable, S. Chu, and D. E. Pritchard, Phys. Rev. Lett. 59, 2631 (1987).
- [5] C. Monroe, W. Swann, H. Robinson, and C. Wieman, Phys. Rev. Lett. 65, 1571 (1990).
- [6] D. E. Pritchard, Phys. Rev. Lett. **51**, 1336 (1983).
- [7] A. L. Migdall, J. V. Prodan, W. D. Phillips, T. H. Bergeman, and H. J. Metcalf, Phys. Rev. Lett. 54, 2596 (1985).
- [8] V. S. Bagnato, G. P. Lafyatis, A. G. Martin, E. L. Raab, R. N. Ahmad-Bitar, and D. E. Pritchard, Phys. Rev. Lett. 58, 2194 (1987).
- [9] E. A. Cornell, C. Monroe, and C. E. Wieman, Phys. Rev. Lett. 67, 2439 (1991).
- [10] N. R. Newbury, C. J. Myatt, and C. E. Wieman, Phys. Rev. A 51, R2680 (1995).
- [11] C. J. Myatt, N. R. Newbury, R. W. Ghrist, S. Loutzenhiser, and C. E. Wieman, Opt. Lett. 21, 290 (1996).
- [12] H. F. Hess, Phys. Rev. B **34**, 3476 (1986).
- [13] J. M. Doyle, J. C. Sandberg, I. A. Yu, C. L. Cesar, D. Kleppner, and T. J. Greytak, Phys. Rev. Lett. 67, 603 (1991).
- [14] W. Petrich, M. H. Anderson, J. R. Ensher, and E. A. Cornell, in *Fourteenth International Conference on Atomic Physics (ICAP XIV)* (AIP, College Park, MD, 1994), Chap. 1M, p. 7.
- [15] K. B. Davis, M.-O. Mewes, M. A. Joffe, M. R. Andrews, and W. Ketterle, Phys. Rev. Lett. 74, 5202 (1995).

- [16] C. S. Adams, H. Lee, N. Davidson, M. Kasevich, and S. Chu, Phys. Rev. Lett. 74, 3577 (1995).
- [17] W. Petrich, M. H. Anderson, J. R. Ensher, and E. A. Cornell, Phys. Rev. Lett. 74, 3352 (1995).
- [18] C. J. Myatt, E. A. Burt, R. W. Ghrist, E. A. Cornell, and C. E. Wieman, Phys. Rev. Lett. 78, 586 (1997).
- [19] M.-O. Mewes, M. R. Andrews, N. J. van Druten, D. M. Kurn, D. S. Durfee, and W. Ketterle, Phys. Rev. Lett. 77, 416 (1996).
- [20] E. Bashkin and A. Meyerovich, J. Phys. Colloq. (France) 41, C7 (1980).
- [21] S. Inouye, M. R. Andrews, J. Stenger, H.-J. Miesner, D. M. Stamper-Kurn, and W. Ketterle, Nature **392**, 151 (1998).
- [22] P. Courteille, R. S. Freeland, D. J. Heinzen, F. A. van Abeelen, and B. J. Verhaar, Phys. Rev. Lett. 81, 69 (1998).
- [23] J. L. Roberts, N. R. Claussen, J. P. Burke, Jr., C. H. Greene, E. A. Cornell, and C. E. Wieman, Phys. Rev. Lett. 81, 5109 (1998).
- [24] V. Vuletic, A. Kerman, C. Chin, and S. Chu, Phys. Rev. Lett. 82, 1406 (1999).
- [25] D. S. Hall, M. R. Matthews, C. E. Wieman, and E. A. Cornell, Phys. Rev. Lett. 81, 1543 (1998).
- [26] S. Inouye, A. P. Chikkatur, D. M. Stamper-Kurn, J. Stenger, D. E. Pritchard, and W. Ketterle, Science 285, 571 (1999).
- [27] D. A. Butts and D. S. Rokhsar, Phys. Rev. A 55, 4346 (1997).
- [28] D. S. Rokhsar, Phys. Rev. Lett. **79**, 2164 (1997).
- [29] G. Ferrari, Phys. Rev. A 59, R4125 (1999).
- [30] J. M. K. V. A. Koelman *et al.*, Phys. Rev. Lett. **59**, 676 (1987).
- [31] B. DeMarco and D. S. Jin, Phys. Rev. A 58, R4267 (1998).
- [32] B. DeMarco, S. Papp, and D. Jin, cond-mat/0101445 (2001).
- [33] K. Helmerson, M. Xiao, and D. Pritchard, in International Quantum Electronics Conference 1990 book of abstracts (IEEE, New York, 1990), No. QTHH4.
- [34] J. Ruostekoski and J. Javanainen, Phys. Rev. A 52, 3033 (1995).
- [35] T. Busch, J. R. Anglin, J. Cirac, and P. Zoller, Europhys. Lett. 44, 1 (1998).
- [36] J. Ruostekoski and J. Javanainen, Phys. Rev. Lett. 82, 4741 (1999).
- [37] A. Fetter and J. Walecka, Quantum Theory of Many-Particle Systems (McGraw-Hill, San Fransisco, 1971).

- [38] M. Houbiers, H. T. C. Stoof, W. I. McAlexander, and R. G. Hulet, Phys. Rev. A 57, R1497 (1998).
- [39] M. Houbiers, R. Ferwerda, H. T. C. Stoof, W. I. McAlexander, C. A. Sackett, and R. G. Hulet, Phys. Rev. A 56, 4864 (1997).
- [40] H. T. C. Stoof, M. Houbiers, C. A. Sackett, and R. G. Hulet, Phys. Rev. Lett. 76, 10 (1996).
- [41] R. Côté, A. Dalgarno, H. Wang, and W. Stwalley, Phys. Rev. A 57, R4118 (1998).
- [42] B. DeMarco, J. L. Bohn, J. J. P. Burke, M. Holland, and D. S. Jin, Phys. Rev. Lett. 82, 4208 (1999).
- [43] F. van Abeleen, F. Verhaar, and A. Moerdijk, Phys. Rev. A 55, (1997).
- [44] W. Geist, A. Idrizbegovic, M. Marinescu, T. Kennedy, and L. You, Phys. Rev. A 61, 013506 (2000).
- [45] R. S. Williamson and T. Walker, Jour. Opt. Soc. Am. B 12, 1393 (1995).
- [46] R. Williamson III, Ph.D. thesis, University of Wisconsin Madison, 1997.
- [47] F. S. Cataliotti, E. A. Cornell, C. Fort, M. Inguscio, F. Marin, M. Prevedelli, L. Ricci, and G. M. Tino, Phys. Rev. A 57, 1136 (1998).
- [48] E. Arimondo, M. Inguscio, and P. Violino, Rev. Mod. Phys. 49, 31 (1977).
- [49] G. Modugno, C. Benko, P. Hannaford, G. Roati, and M. Inguscio, Phys. Rev. A 60, R3373 (1999).
- [50] T. Bergeman, G. Erez, and H. Metcalf, Phys. Rev. A 35, 1535 (1987).
- [51] B. DeMarco, H. Rohner, and D. Jin, Rev. Sci. Intrum. 70, 1967 (1999).
- [52] R. Legere and K. Gibble, Phys. Rev. Lett. 81, 5780 (1998).
- [53] B. DeMarco and D. Jin, Science **285**, 1705 (1999).
- [54] M. J. Holland, B. DeMarco, and D. S. Jin, Phys. Rev. A **61**, 053610 (2000).
- [55] D. Jin, B. DeMarco, and S. Papp, in *ICAP 2000 Proceedings* (AIP, College Park, MD, 2001), Vol. 17, p. 414.
- [56] S.-K. Yip and T.-L. Ho, Phys. Rev. A **59**, 4643 (1999).
- [57] T.-L. Ho and V. Shenoy, J. Low Temp. Phys. 111, 937 (1999).
- [58] M. Amoruso, I. Meccoli, A. Minguzzi, and M. P. Tosi, Euro. Phys. J. D 7, 441 (1999).
- [59] M. Amoruso *et al.*, Euro. Phys. J. D 8, 361 (2000).
- [60] G. M. Bruun and C. W. Clark, Phys. Rev. Lett. 83, 5415 (1999).

- [61] G. Bruun and C. Clark, Phys. Rev. A **61**, 061601(R) (2000).
- [62] L. Vichi and S. Stringari, Phys. Rev. A 60, 4734 (1999).
- [63] A. Minguzzi and M. Tosi, Phys. Rev A 63, 023609 (2001).
- [64] J. Schneider and H. Wallis, Phys. Rev. A 57, 1253 (1998).
- [65] G. Bruun and K. Burnett, Phys. Rev. A 58, 2427 (1998).
- [66] M. Mackie, E. Timmermans, R. Cote, and J. Javanainen, Optics Express 8, 118 (2000).
- [67] M. J. Bijlsma, B. A. Heringa, and H. T. C. Stoof, Phys. Rev. A 61, 053601 (2000).
- [68] A. Truscott, K. Strecker, W. Alexander, G. Patridge, and R. Hulet, Science Express Reports 10.1126, 1059318 (2001).
- [69] J. Bohn, Phys. Rev. A **61**, 053409 (2000).
- [70] J. R. Ensher, Ph.D. thesis, University of Colorado Boulder, 1998.
- [71] K. L. Corwin, Z.-T. Lu, C. F. Hand, R. Epstein, and C. Wieman, Applied Optics 37, 3295 (1998).
- [72] M. R. Matthews, D. S. Hall, D. S. Jin, J. R. Ensher, C. E. Wieman, E. A. Cornell, F. Dalfovo, C. Minniti, and S. Stringari, Phys. Rev. Lett. 81, 243 (1998).
- [73] D. M. Stamper-Kurn, M. R. Andrews, A. P. Chikkatur, S. Inouye, H.-J. Miesner, J. Stenger, and W. Ketterle, Phys. Rev. Lett. 80, 2027 (1998).
- [74] K. B. MacAdam, A. Steinbach, and C. Wieman, Am. J. Phys. 60, 1098 (1992).
- [75] J. Marquardt, F. Cruz, M. Stephens, C. Oates, L. Hollberg, J. Bergquist, D. Welch, D. Mehuys, and S. Sanders, in *Proc. SPIE* (SPIE, New York, 1996), Vol. 2834, p. 34.
- [76] D. Mehuys, D. Welch, and L. Goldberg, Elec. Lett. 28, 1944 (1992).
- [77] P. della Porta, C. Emili, and S. Hellier, SAES Technical Report 18, 3.
- [78] L. Hackspill, Helv. Chim. Acta **11**, 1008 (1928).
- [79] R. Honig and D. Kramer, Vapor Pressure of the Elements, 1968.
- [80] D. Sesko, T. Walker, C. Monroe, A. Gallagher, and C. Wieman, Phys. Rev. Lett. 63, 961 (1989).
- [81] W. Ketterle, K. B. Davis, M. A. Joffe, A. Martin, and D. E. Pritchard, Phys. Rev. Lett. 70, 2253 (1993).
- [82] K. Lindquist, M. Stephens, and C. Wieman, Phys. Rev. A 46, 4082 (1992).
- [83] Y. Gott, M. Ioffe, and V. Tel'kovski, Nucl. Fusion 1962 Supplement 3, 1045,1284 (1962).

- [84] D. Pozar, *Microwave Engineering* (John Wiley and Sons, Inc., New York, 1998).
- [85] W. Gerlach and O. Stern, Zeit. für Phys. 8, 110 (1922), *ibid* 9, 349.
- [86] H. M. J. M. Boesten *et al.*, Phys. Rev. A 54, R3726 (1996).
- [87] R. Côté et al., Phys. Rev. A 57, R4118 (1998).
- [88] A. Derevianko, W. Johnson, M. Safronova, and J. Babb, Phys. Rev. Lett. 82, 3589 (1999).
- [89] H. Wang, A. Nikolov, J. Ensher, P. Gould, E. Eyler, W. Stwalley, J. Burke, J. Bohn, C. Greene, E. Tiesinga, C. Williams, and P. Julienne, Phys. Rev. A 62, 052704 (2000).
- [90] J. Sakurai, Modern Quantum Mechanics (Addison-Wesley, New York, 1994).
- [91] J. Burke, Jr., Ph.D. thesis, University of Colorado Boulder, 1999.
- [92] C. R. Monroe, E. A. Cornell, C. A. Sackett, C. J. Myatt, and C. E. Wieman, Phys. Rev. Lett. 70, 414 (1993).
- [93] J. Taylor, Scattering Theory (Wiley, New York, 1972).
- [94] E. Arimondo, M. Inguscio, and P. Violino, Rev. Mod. Phys. 49, 31 (1977).
- [95] J. Burke, C. Greene, J. Bohn, H. Wang, P. Gould, and W. Stwalley, Phys. Rev. A 60, 4417 (1999).
- [96] M. Marinescu, H. Sadeghpour, and A. Dalgarno, Phys. Rev. A 49, 982 (1994).
- [97] J. P. Burke, Jr., J. L. Bohn, B. D. Esry, and C. H. Greene, Phys. Rev. A 55, R2511 (1997).
- [98] H. F. Hess, G. P. Kochanski, J. M. Doyle, N. Masuhara, D. Kleppner, and T. J. Greytak, Phys. Rev. Lett. 59, 672 (1987).
- [99] D. Pritchard et al., in Proceedings of the 11th International Conference on Atomic Physics (World Scientific), Vol. 2834, p. 619.
- [100] N. Masuhara, J. M. Doyle, J. C. Sandberg, D. Kleppner, T. J. Greytak, H. Hess, and G. P. Kochanski, Phys. Rev. Lett. 61, 935 (1988).
- [101] N. Ramsey, *Molecular Beams* (Oxford, 1956).
- [102] K. B. Davis, M.-O. Mewes, and W. Ketterle, Appl. Phys. B 60, 155 (1995).
- [103] H. C. W. Beijerinck, Phys. Rev. A 61, 033606 (2000).
- [104] H. C. W. Beijerinck, Phys. Rev. A 62, 063614 (2000).
- [105] M. E. Gehm, K. M. O'Hara, T. A. Savard, and J. E. Thomas, Phys. Rev. A 58, 3914 (1998).
- [106] P. Valkering, Ph.D. thesis, University of Utrecht, 1999.

- [107] J. Rubbmark, M. Kash, M. Littman, and D. Kleppner, Phys. Rev. A 23, 3107 (1981).
- [108] W. Greiner, L. Neise, and H. Stöcker, *Thermodynamics and Statistical Mechanics* (Springer-Verlag, New York, 1995).
- [109] L. D. Landau and E. M. Lifshitz, *Statistical Physics Part 1* (Pergamon Press, Elmsford, New York, 1958).
- [110] W. Demtröder, Laser Spectroscopy (Springer, Berlin, 1996).
- [111] M. R. Andrews, M.-O. Mewes, N. J. van Druten, D. S. Durfee, D. M. Kurn, and W. Ketterle, Science 273, 84 (1996).
- [112] A. Das and T. Ferbel, Introduction to Nuclear and Particle Physics (John Wiley and Sons, Inc., New York, 1994).
- [113] M. Houbiers and H. T. C. Stoof, in Proceedings of the 21st International Conference on Low Temperature Physics (, 1996), Chap. S1, p. 551.
- [114] F. Reif, Fundamentals of Statistical and Thermal Physics (McGraw-Hill, New York, 1965).

Appendix A

LAB NOTEBOOK INDEX

added preservation coil #16: 84 install AG coil #12: 46 transfer tube electromagnets #9: 112 information on cooling diodes and performance #12: 133-, #14: 56,62,82 calibration of ZFAS-1000 attenuator #4: 5 lifetime, heating with valve shut overnight #12: 33 surface mount component resonant matching of microwave coil #15: 185 recent noise data #16: 138 thermometry data

> last expansion data $m_f = 9/2 \# 17$: 12-13 mixture fit vs FD fit #12: 104 good fugacity data #13: 45, #12: 96, #14: 23,103 angularly averaged cloud profiles #12: 134-135 thermometry systematic vs 7/2 cleanout time #12: 80 mixture fit L dependence on detuning #12: 38-43 FD fit parameters dependence on detuning #15: 18 saturation effects #10: 29 different cuts for energy fits #15: 6 L vs peak OD #10: 115 example of effect of bad probe beam #10: 132 test of image subtraction #15: 32-33

evaporation

study of end of single frequency evaporation #14: 109; #15: 74 two-frequency evaporation performance vs PS voltage #16: 53 best two-frequency evaporation temps (using 8648) #15: 144 check that microwave line spacing make sense #12: 11 two-frequency with 7/2 cleanout #13: 45, 51; #12: 25

data on variation in final T with variation in initial N for evap 1 # 12: 116 evap2 in single-frequency trap #10: 137/2 removal trajectory #12: 109; #15: 44 trajectories with different initial N #12:32removal sweep optimization #12: 13,16 single frequency trajectory #11: 96, 111; #10: 103, 99, 97, 92, 68, 45, 21, 17, 15, 13 microwave lineshapes, checks on $\Delta f \# 8: 124, 144$ very early data on badness of 8656B synth. (and switch to E4420B for evap1) #9: 17 switch to "sheet" coil in evaporation #8:153cooling rate in evap 2 # 15: 88 measure microwave power from removal probability #15: 106, 127, 147 problems caused by synths w/ different duty cycles #15: 130effect of power on end of evap1, varying cut to get 50/50% mixture #15: 141evap2 trajectory with 8648 synth #16:16-17evap2 trajectory with 8647 synth #16:41evap2 trajectory with E4420B synth #16:43implement fast 5/3+3/2 cleanout #16:47 implement 5/2 cleanout during evap 2 # 16: 50 evap1 behavior vs power at fixed cut #16: 109, 112; #17: 76 recent evap2 trajectory #16: 140-142; #17: 64,79 optimize bottom of sweep for evap2 #16: 45imaging optics #7: 73-81

calibration vs PD #5: 109-116 calibration of low magnification #11: 61 check on N #10: 117-125 last focus #14: 96 last/best magnification measurement (and sag measurement) #16: 100-102 absorption lineshapes for different Q fields #15: 25; #16:104,131

MOT behavior

imaging

MOT temp #5: 39; #4: 149; #2: 148 effect of cold MOT stage #6: 37-39,79,138 dependence of catch/loss on repump power #2: 22-26, 42; #6: 20-21 Sci MOT lifetime #2: 3 max fill vs bypass power #3: 35, 42; #7: 10,16; #8: 58 last science MOT fill #17: 19 collection and science MOT behavior vs repump power #12: 144-152 best shims for fill #13: 9 add dark spot #13: 9 last? (detailed) science MOT detuning optimization #14: 87,113; #15: 87

\mathbf{push}

push polarization and detuning dependence #2: 97; 1/12/01push speed #2: 13-14 push duration #2: 15; #3: 86 push speed vs detuning, pulse vs LVIS #1: 152

elastic collision measurement

correct fitting function, reanalysis #6: 53-54 data #5: 60-152 check on scaling with N #5: 66

inelastic rates

looking for $m_f = 5/2 + m_f = 7/2$ Feshbach resonance #11: 58-64

optical pumping and spin composition

better evaporation performance for dark state transition #5: 34-35; #7: 5 fix optical pumping polarization and repumping #6: 135-143 measurement of spin composition #14: 73 fraction transferred into magnetic trap vs $\lambda/4$ angle and detuning #13: 18 optical pumping dependence on power #15: 53 effect of opt pumping on temperature #3: 122-124; #4: 145; #6: 106 compare optical pumping on two different transitions #13: 17-18 repump jump detuning effect on spin composition #14: 123 set up optical pumping beam using reflection #6: 114

magnetic trap

install hall probes for magnetic trap current servos #11: 128-134 frequency response of cloverleaf coils #11: 123-125 explicit magnetic trap wiring diagram #16:150-151magnetic trap sensitivity to air conditioner and drift #13: 81-102 implementation of fast magnetic trap turn-off #14: 45-50paranoia: magnetic trap sensitivity to the door position #15: 97 set up control voltages with 4 resistor dividers #9: 141; #10: 77 bias field short term stability before optimizing servos #11:65-83,108bias field short term stability after optimizing servos #12:7-10drift vs temp of phenolic rods #11: 144 effect of science MOT coils on drift #12: 6effect of magnetization of stuff on bias field #16:71-80magnetic trap coil wiring and plumbing #1:16-17magnetic trap field measurements #1:21-26magnetic trap coil resistances #1: 28-29single frequency trap #10: 9-12water flow #2: 149

mode matching on load #5: 37 drift caused by transfer tube magnets #9: 110 slosh caused by SRS mismatch on ramp #9: 70 effect of turning off alkali dispensers during evaporation #16: 68