

A Tall, Cold Glass of Calcium: Measurement-Based Cooling and Mid-Circuit Readout in
the *omg* Architecture

by

Sean Brudney

A dissertation accepted and approved in partial fulfillment of the
requirements for the degree of
Doctor of Philosophy
in Physics

Dissertation Committee:

Benjamín Aleman, Chair

David Allcock, Advisor & Core Member

Dietrich Belitz, Core Member

Cathy Wong, Institutional Representative

University of Oregon

Spring 2025

© 2025 Sean Brudney
All rights reserved.

DISSERTATION ABSTRACT

Sean Brudney

Doctor of Philosophy in Physics

Title: A Tall, Cold Glass of Calcium: Measurement-Based Cooling and Mid-Circuit Readout in the *omg* Architecture

Trapped ions are a leading candidate for the development of practical quantum computers beyond the Noisy Intermediate-Scale Quantum regime. Metastable states in several ion species provide long-lived qubit subspaces that can be defined for information storage and processing alongside co-trapped ions of the same species used for dissipative operations like laser cooling and fluorescence detection in the ground manifold without disturbing the spectrally-isolated metastable qubits. This architecture lays out a road map for scaling trapped ion quantum computer systems using a single trapped ion species while simultaneously maintaining the functionality of more hardware intensive dual species systems. This work demonstrates dual functionality in a single species chain of $^{40}\text{Ca}^+$. To simplify experimental overhead, we study an alternative to the oven loading method for calcium atoms using laser desorption, which several neutral and ion trap experiments use for other species. After constructing and characterizing the ion trap, as well as building up the laser systems necessary to trap, cool, and manipulate the ions, we herald a two ion configuration with the qubit subspace of one ion defined in the metastable $D_{5/2}$ manifold and the other in the ground $S_{1/2}$ manifold. After heralding, we maintain coherent information in the metastable qubit during sympathetic cooling to the motional ground state via projection cooling and perform non-destructive mid-circuit measurements that will likely be necessary for syndrome extraction in fault-tolerant quantum error correction.

This dissertation includes previously published and unpublished co-authored material.

CURRICULUM VITAE

NAME OF AUTHOR: Sean Brudney

GRADUATE AND UNDERGRADUATE SCHOOLS ATTENDED:

University of Oregon, Eugene, OR, USA
University of Notre Dame, Notre Dame, IN, USA

DEGREES AWARDED:

Doctor of Philosophy, Physics, Spring 2025, University of Oregon
Master of Science, Physics, Winter 2021, University of Oregon
Bachelor of Science, Physics, Spring 2017, University of Notre Dame

AREAS OF SPECIAL INTEREST:

Quantum information and sensing, ion traps

PUBLICATIONS:

Mid-circuit ground-state cooling and ancilla readout in the *omg* architecture.
S. Brudney, G. J. Gregory, E. R. Ritchie, D. J. Wineland, D. T. C. Allcock,
and J. O'Reilly. *In preparation*.

Everybody Loves LIAD: Using non-resonant light to source calcium for ion traps
and other atomic platforms. S. Brudney, J. O'Reilly, and D. T. C. Allcock. *In
preparation*.

Spontaneous Raman scattering out of a metastable atomic qubit. I. D. Moore, A.
Quinn, J. O'Reilly, J. Metzner, S. Brudney, G. J. Gregory, D. J. Wineland, and
D. T. C. Allcock. [arXiv preprint arXiv:2505.04854 \(2025\)](#).

High-fidelity entanglement of metastable trapped-ion qubits with integrated
erasure conversion. A. Quinn, G. J. Gregory, I. D. Moore, S. Brudney, J.
Metzner, E. R. Ritchie, J. O'Reilly, D. J. Wineland, and D. T. C. Allcock.
[arXiv preprint arXiv:2411.12727 \(2024\)](#).

Two-mode squeezing and $SU(1,1)$ interferometry with trapped ions. J. Metzner, A. Quinn, S. Brudney, I. D. Moore, S. C. Burd, D. J. Wineland, and D. T. Allcock. [Phys. Rev. A 110, 022613 \(2024\)](#).

Observing super-quantum correlations across the exceptional point in a single, two-level trapped ion. A. Quinn, J. Metzner, J. E. Muldoon, I. D. Moore, S. Brudney, S. Das, D. T. C. Allcock, and Y. N. Joglekar. [arXiv preprint arXiv:2304.12413 \(2023\)](#).

ACKNOWLEDGEMENTS

I would like to thank my advisor David Allcock and group co-P.I. Dave Wineland for the opportunity to work in the Oregon Ions research group. Thank you to graduated members Alex Quinn, Jeremy Metzner, and Daniel Moore who were the founding graduate students without whom four chapters of this document would not have been possible (or at least would be substantially less interesting). Thank you to current graduate students Gabe Gregory and Evan Ritchie who both assisted me on several occasions to operate the trap and understand important concepts of ion trapping. Thank you to undergraduate member Marcus Polk who is the best mentee ever and for setting up the 729 nm cavity essential to the best parts of my research. Good luck in Indiana! Thank you to my cohort members, especially Matthew Brown and Haidar Esseili without whom I would not have done nearly as well, *i.e.*, failed, in my coursework. Thank you to the Graduate Teaching Fellows Federation for representing graduate employees against the University of Oregon and making it possible for me to afford both rent and health care in an occupation known for its financial hardships. And finally a special thank you to Jameson O'Reilly who, even though he joined our group as a postdoctoral researcher in the final year of my degree, has had an immense impact on my education and training. Thank you for *patiently* repeating your explanations to me, expanding my understanding of ion trapping, quantum information, and physics in general, and for dragging me kicking and screaming to the goal line.

to Scribles Bistro & Deli for their delicious sandwiches which are a shining light in the Eugene food scene

DISCLAIMER

No generative “A.I.” was used in the making of this document or the work presented herein.

TABLE OF CONTENTS

Chapter	Page
I. INTRODUCTION	18
1.1. Dissertation Structure	19
1.2. Previously Published and Unpublished Work	20
II. DESIGNING & BUILDING AN ION TRAP	22
2.1. Trapping Potentials	22
2.2. Rod Traps	26
2.3. Numerical Aperture	27
2.4. Simulations	30
2.5. Laser-Free Control	35
2.5.1. RF Rabi Flopping	35
2.5.2. CVQC Operations	37
2.5.2.1. Parametric Drive Electronics	38
2.6. Future Trap Developments	41
III. LIGHT INDUCED AND THERMAL ATOMIC DESORPTION	47
3.1. Experimental Setup	49
3.1.1. Vacuum chamber	49
3.1.2. Fluorescence Collection	49
3.1.3. Substrates	51
3.1.4. LIAD Lasers	52
3.2. Methods	53
3.2.1. Temperature Measurements	53
3.2.2. Thin Film Measurements	54

Chapter	Page
3.3. Results	55
3.3.1. Fused Silica and Borosilicate Glass	55
3.3.2. Graphene Sheets	57
3.4. Analysis	58
3.4.1. Short Duration Signal	58
3.4.2. Long Duration Signal	59
3.5. Conclusions and Future Work	60
IV. LASERS	62
4.1. Dipole Lasers	62
4.1.1. Ionizing Atoms	63
4.1.2. Fluorescence Detection	63
4.1.3. Cooling	64
4.1.3.1. Doppler Cooling	64
4.1.3.2. EIT Cooling	65
4.1.4. <i>m</i> -Qubit Preparation	66
4.2. Wavemeter Locking Lasers	67
4.2.1. Allan Deviation	69
4.3. Far-detuned Raman Lasers	73
4.3.1. Fiberized Injection Locking	74
4.4. Narrow Linewidth Quadruple Transition	76
4.4.1. PDH Lock	78
4.4.2. Polarization, Shifts, and Rabi Frequencies	83
4.4.3. Off-Resonant Driving	86
V. MID-CIRCUIT SYMPATHETIC COOLING	88
5.1. <i>mg</i> Crystal Preparation and Measurement	90

Chapter	Page
5.1.1. Heralding mg Crystals	90
5.1.2. mg Readout Scheme	92
5.1.2.1. Cabinet Shelving and Bucket Deshelving	96
5.1.3. mg Crystal SPAM	97
5.1.3.1. mg Crystal SPAM Error Sources	98
5.1.4. mg RF Control	99
5.2. Heating and Sympathetic Re-Cooling	100
5.2.1. Measuring Fock State Distributions via Sideband and Carrier Driving	101
5.2.2. Idle Heating	104
5.2.3. Displacement Heating	104
5.2.4. Fluorescence Heating	105
5.2.5. Doppler Cooling	107
5.2.6. Electromagnetically Induced Transparency Cooling	108
5.3. m -Qubit Compatibility	110
5.3.1. Lifetime Measurements	111
5.3.2. Coherence Times	111
5.3.3. m - m Protected Subspace	113
VI. PROJECTION COOLING	119
6.1. Heralding Motional Ground States	120
6.1.1. Passing Fluorescence Checks	122
6.1.2. Temperature Limits	123
6.2. Minimum Temperature Measurements of mg Crystals	124
6.3. m -Qubit Compatibility	126
6.4. Non-Destructive Mid-Circuit Measurements	127
6.4.1. Optical Qubit Ancilla Measurement Results	129

Chapter	Page
VII. SUMMARY AND OUTLOOK	134
7.1. Projection Cooling Limits	134
7.2. Mid-Circuit Measurement Improvements	135

LIST OF FIGURES

Figure	Page
2.1. View of our linear rod trap CAD model.	23
2.2. Assembled rod trap and vacuum chamber.	24
2.3. Diagram and definitions of geometry for determining NA_{eff}	29
2.4. Plot of the DC needle voltage.	31
2.5. Electrical circuit diagram of the trapping RF, DC bias, and parametric driving RF signals.	32
2.6. Radial secular frequency and RF frequency as a function of the RF voltage.	33
2.7. Plot of the necessary V_{RF} to trap $^{40}\text{Ca}^+$ ions.	34
2.8. Predicted values of axial frequency plotted with measured values of the axial frequency.	35
2.9. Measured values of the radial frequency and corresponding q_x values as a function of V_{RF}	36
2.10. Electrical testing of RF driving sources through the squeeze box.	39
2.11. Definitions for angle θ between the compensation rods and the trapping rods and the mode orientation ϕ	42
2.12. Compensation rod voltages, RF bias voltages, mode splitting values, and mode orientation values plots.	44
2.13. Simulations of mode splittings and mode orientations with compensation rod voltage and RF bias voltage	45
2.14. Pseudopotential depths for different trap designs.	46
3.1. Schematic of laser and calcium interactions inside the vacuum chamber. . . .	50
3.2. Imaging system schematic.	51
3.3. Comparison of signals from oven and laser desorption.	54

Figure	Page
3.4. Amount of transmitted light passing through a thin film of calcium metal on a glass substrate.	55
3.5. Fluorescence response from 450 nm laser on fused silica glass.	56
3.6. Sustained signals of desorbed calcium off fused silica glass.	57
3.7. Fluorescence response from moving the 450 nm laser beam spot across a graphene sheet.	58
4.1. Diagram of ionization for neutral calcium and electronic states of $^{40}\text{Ca}^+$ and relevant transitions.	63
4.2. Zeeman sublevels in the $S_{1/2}$, $P_{1/2}$, $P_{3/2}$, $D_{3/2}$, and $P_{5/2}$ manifolds. Designation of the m - and g -qubit.	67
4.3. The wavemeter inside the temperature control box.	69
4.4. Example spectroscopy and fit of the 397 nm laser frequency detuning against the $^{40}\text{Ca}^+$ transition for $S_{1/2} \leftrightarrow P_{1/2}$	70
4.5. Measurement of cooling lasers and deviations of the laser laboratory environment conditions.	71
4.6. Allan deviation of the 397 nm laser with and without interpolated dead time compared against an ideal laser.	73
4.7. Allan deviation of the 397 nm laser with wavemeter lock on and off.	74
4.8. Raman transitions for driving between the m -qubit states using the far-detuned 976 nm lasers.	75
4.9. Schematic diagram of the fiberized injection lock setup of the 976 nm Raman lasers as well as the beat note measurement.	76
4.10. Frequency spectrum of the beat note between 976 nm Raman beams.	77
4.11. Schematic diagram of PDH locking setup.	79
4.12. Diagram of the frequency spectrum of the incident laser beam to the cavity in the PDH setup.	81
4.13. Ring-down measurement and setup diagram.	82
4.14. Frequency scan of 729 nm carrier transition with an interrogation time of 10 ms and lowered 729 nm power.	83

Figure	Page
4.15. All ten allowed transitions for $S_{1/2} \leftrightarrow D_{5/2}$	84
4.16. Geometry of the 729 nm beam in the trap setup with parameterized relative coupling strength of $S_{1/2} \leftrightarrow D_{5/2}$ transitions.	85
4.17. Population transfer from off-resonant driving overlaid on the trace of the frequency spectrum of the reflected cavity signal.	87
5.1. Real images of two-ion configurations in our trap corresponding to preparation in the $D_{5/2}$ manifold and $S_{1/2}$ manifold.	93
5.2. Diagrams of the mg crystal readout scheme.	94
5.3. Individual RF Rabi flopping of m and g -qubits.	100
5.4. Data and fit for readout on BSB and carrier.	103
5.5. Average occupation state $\bar{n}$ for various idle times.	104
5.6. Readout of displacement $ \alpha $ due to RF shuttling.	105
5.7. Measured $\bar{n}$ from heating an mg crystal using the fluorescing lasers.	107
5.8. mg crystal Doppler cooled for variable time fitted to Eq. (5.14).	109
5.9. mg crystal EIT cooled for variable time fitted to exponential decay.	110
5.10. Lifetime measurements with various lasers turned on.	112
5.11. Ramsey contrast and pulse sequence for spin-echo experiments.	117
5.12. Parity Π measurements of the protected subspace for two m qubits and pulse sequence.	118
6.1. Diagram of projection cooling.	121
6.2. Theory model for projection cooling with variable BSB pulse times and numbers.	123
6.3. 976 nm RSB driving calibration of projection cooling.	125
6.4. Experiments for testing projection cooling compatibility with m -qubits.	131
6.5. Comparison of the readout of the m qubit with readout via the o -qubit.	132

LIST OF TABLES

Table	Page
3.1. LIAD laser table	52
4.1. $S_{1/2} \leftrightarrow D_{5/2}$ magnetic field sensitivities, Zeeman shifts, and relative couplings.	86
5.1. Probability table for detecting each type of two-ion configuration.	92
5.2. Corresponding detection result for a given sequence of bright ion detections in mg crystals.	96
5.3. SPAM fidelity of each of the four qubit state combinations for an mg crystal.	98
5.4. Simulation of error due to π -pulse time uncertainty on readout of an mg crystal.	98
5.5. Simulation of error due to unequal illumination of the ions by 729 nm light on readout for an mg crystal.	99
5.6. Variables used to calculate fluorescence heating rate for our trap setup.	106
6.1. Corresponding state for o -qubit ancilla readout scheme.	129
6.2. Results of ancilla readout using the o -qubit to determine the m -qubit.	133

CHAPTER I

INTRODUCTION

The idea that quantum computers instead of classical computers could be used to simulate quantum systems was first suggested by Yuri Manin in 1980 [1]. Since that time, the development of technologies to measure and manipulate quantum systems has provided a laundry list of platforms to study [2]. In the last thirty years especially, the industry for quantum information, sensing, and computing (and the resources available to academia [3]) have boomed. Thanks to this investment, we have entered what is referred to as the Noisy Intermediate-Scale Quantum (NISQ) era [4, 5] in which quantum computers contain on the order of dozens of qubits that are limited in utility due to significant noise.

Trapped ion systems have been a leading platform for the development of quantum computers due to their naturally identical internal structures, long coherence times [6, 7], and the ability to isolate them from their surrounding environment [8]. Ion traps also currently hold the highest fidelities for one [9, 10] and two qubit gates [11, 12] among all quantum computing platforms. Already, large systems using arrays of ion traps have been demonstrated, known as the quantum charge-coupled device (QCCD) architecture [13, 14] and the field of error-correction is finding new, more efficient ways to achieve fault-tolerance in a generic trapped ion system [15].¹ Simulations of small quantum systems are already possible in traps with < 10 qubits [16], with fully fault-tolerant systems potentially able to efficiently simulate complex molecular interactions in chemistry and biology [17], quantum phase transitions and high temperature superconducting materials in condensed matter physics [18], and the Standard Model and

¹“Oh sure, the Frinkiac 7 looks impressive — don’t touch it! But I predict that within 100 years, computers will be twice as powerful, ten thousand times larger, and so expensive that only the five richest kings of Europe will own them.” — Professor Frink, *The Simpsons* S7E23, “Much Apu About Nothing”

effective field theories in high energy physics [19]. Development of trapped ion systems is also benefiting the field of quantum sensing [20], with ions being highly sensitive to small displacements [21, 22] and suggestions of using trapped ions for dark matter detection [23].

Along with long circuits and error correction comes the need for separation of functionality between qubits, which is possible using multiple species of ion within the same trap [24–26]. However, dual species traps come with the drawback of doubling the number of atom sources and laser systems as well as increased engineering efforts for maintenance and control, which can limit scalability both physically and financially. The so-called *omg* architecture [27], also referred to as ‘dual-type’ qubits [28], provides the ability to keep dual species functionality without the overhead of trapping and controlling multiple species. Using laser addressing, the *omg* architecture designates the functionality of ions by encoding some qubits in the long-lived metastable electronic manifolds where lasers that address ground qubits will be far off-resonance. The work discussed in this dissertation demonstrates the utility of operating within this framework and provides new methods for ground state cooling and mid-circuit measurements. In particular, projection cooling provides a ground state cooling operation that does not disturb metastable qubits, which allows for sympathetic cooling of single species ion chains and improves the fidelity of non-destructive mid-circuit measurements, which are crucial for error correction in long circuits [29, 30].

1.1 Dissertation Structure

Chapter II starts with an introduction to the basic physics of ion trapping and provides details of the design and construction of the Oregon Ions’ first (and currently only) ion trap. It ends with details describing a potential future trap design that could address limitations of already published work.

Chapter III describes an alternative to oven loading of ion traps as well as other atomic physics experiments.

Chapter IV is a description of all the laser systems used for our ion trap. This includes our dipole lasers, 976 nm far-detuned injection-locked Raman lasers, and 729 nm cavity-locked quadrupole laser.

Chapter V describes an implementation of the *omg* architecture with the heralding of *mg* crystals as well as dual species functionality with sympathetic cooling using the *g*-qubit. Additionally, we explain how typical ground state motional cooling techniques are incompatible with *m*-qubits.

Chapter VI presents projection cooling as a method for heralding motional ground states that is compatible with the *omg* architecture. Testing shows temperatures below the EIT cooling limit. Results from non-destructive ancilla measurements after projection cooling *mg* crystals are reported.

Chapter VII provides a summary and outlook for future directions.

1.2 Previously Published and Unpublished Work

Chapters II and IV are descriptions of the ion trap (the primary experimental apparatus of the research group) and surrounding systems and includes work and details previously described in the dissertations of Drs. Daniel Moore [31], Jeremy Metzner [32], and Alex Quinn [33]. Design, construction, and testing of the ion trap and surrounding systems was done by these three and myself under supervision of our advisor, Prof. David Allcock. Additional upgrades to the data acquisition system were done by Gabe Gregory. Chapter II includes details of previously published work describing the experimental apparatus (the ‘squeeze box’) to perform motional squeezing and beamsplitter operations on trapped ions (see Section 2.5.2.1) [22]. Construction and testing of the squeeze box was primarily done by myself and the experiment and analysis described in the paper

were performed by Drs. Metzner and Quinn. Chapter II contains analysis of simulations using COMSOL Multiphysics™ software package done by Dr. Quinn with the analysis performed by myself. Chapter III is currently unpublished work being prepared for pre-print under the supervision of Prof. Allcock. Chapter IV contains a description of the 976 nm Raman injection laser system used in previously published work (see Section 4.3) [34]. Setup and testing of the injection laser system was performed by myself while the experiment described in the paper was performed by Dr. Quinn with simulations and error budget calculated by Gabe Gregory. Chapters V and VI describe currently unpublished work to be published performed by myself and Dr. Jameson O'Reilly. The *mg* crystal SPAM error tables (Section 5.1.3.1) were produced using simulations performed by Gabe Gregory. The simulations of projection cooling (Section 6.2) were performed by myself with assistance from Connor Burns. This dissertation was written and edited under the supervision of Dr. O'Reilly and Prof. Allcock.

CHAPTER II

DESIGNING & BUILDING AN ION TRAP

The design, construction, and characterization of the ion trap was performed by Drs. Daniel Moore [31], Jeremy Metzner [32], Alex Quinn [33], Prof. David Allcock, and myself. The building and characterization of the ‘squeeze box’ was primarily performed by myself as well as Dr. Metzner, the description of which is published in J. Metzner, *et al.*, *Two-mode squeezing and $SU(1,1)$ interferometry with trapped ions*, [Phys. Rev. A **110**, 022613 \(2024\)](#). All simulations involving COMSOL were performed by Dr. Quinn.

Since their introduction as atomic quadrupole mass filters in 1953 by Wolfgang Paul [35, 36], ion traps have developed as a tool with extensive use in physics and science in general. They are useful for a variety of applications including quantum sensing and information [37]. Many types of ion traps exist today: 3D rod traps [22, 34], 2D surface-electrode traps [38], which exist in many forms including rings [39] and race tracks [40]; traps with integrated fibers [41, 42], cavities [43], and half-sphere mirrors [44]; and many others. For this dissertation,¹ the trap used is the Oregon Ions linear 3D rod trap with calcium-40 as the trapped atomic species.

2.1 Trapping Potentials

The linear ion trap is a form of the generic quadrupole or Paul trap [36, 45]. According to Earnshaw’s theorem [46], a single charged particle in free space cannot be trapped in three dimensions using only static electric potentials. This conclusion can be derived from Gauss’s law, which states that the divergence of an electric field in free space is zero. For a charged particle to remain in a stable equilibrium, all force field lines must point inward to a point, which is a negative divergence; therefore, static electric potentials alone are insufficient. Static electric fields can produce a saddle point, but

¹“Ketchup... Catsup... Ketchup... Catsup... Cats... K-k... Huh, I’m in way over my head.”
— Mr. Burns, *The Simpsons* S8E21, “The Old Man and Lisa”

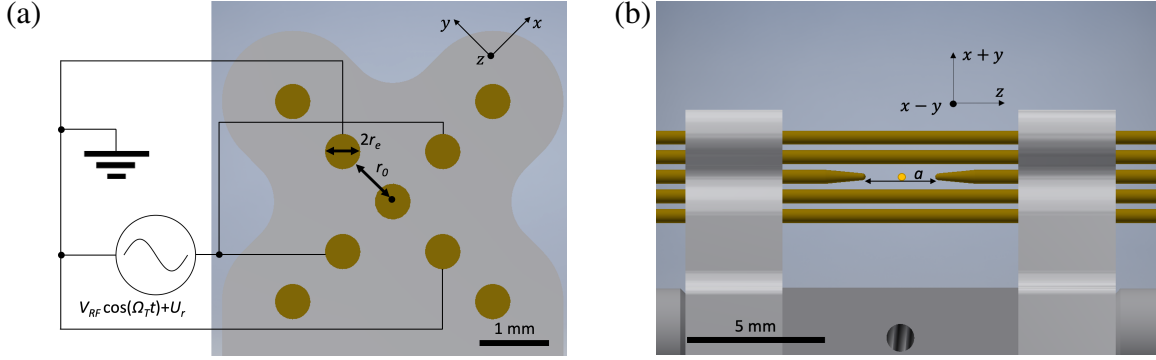


Figure 2.1. (a) Axial view of our linear rod trap CAD model. The center circle is the center rod or needle, the four middle circles are the trapping rods, and the four outer circles are the compensation rods. r_e is the electrode radius and r_0 is the trap axis to electrode surface distance or trap radius. Two opposing inner rods are connected to the oscillating voltage V_{RF} at frequency Ω_T with offset $U_{r,RF}$ and the other two inner rods are held at ground. (b) Side view of a linear rod trap CAD model where a is the tip-to-tip distance between the needles and the yellow dot is an exaggerated ion giving its approximate location in the trap. There is an $(x + y)$ axis towards the top of the page and $(x - y)$ axis out of the page in (b) since the x and y axes align with the ‘X’ shape of the electrodes.

small perturbations in the anti-confining direction will result in a charged particle being pushed away. By using oscillating electric fields, a saddle point’s confining and anti-confining axes are switched back-and-forth, which for a homogeneous field time averages to zero net forces and no stable confinement. In a periodic inhomogeneous field, the time average gives a non-zero net force that under certain conditions can be used to stably confine a charged particle. A linear Paul trap generates a periodic inhomogeneous field, a quadrupole field, from voltages on electrodes to produce this time-averaged potential or ‘pseudopotential’ in two dimensions and static fields to confine in the third dimension. We use rod-shaped electrodes arranged as shown in Figs. 2.1 and 2.2.

Near the center of a linear Paul trap, the radial potential around the center z -axis (defined in Fig. 2.1) is

$$\Phi = \mu \frac{(V_{RF} \cos \Omega_T t + U_r)}{2} \left(1 + \frac{x^2 - y^2}{r_0^2} \right), \quad (2.1)$$

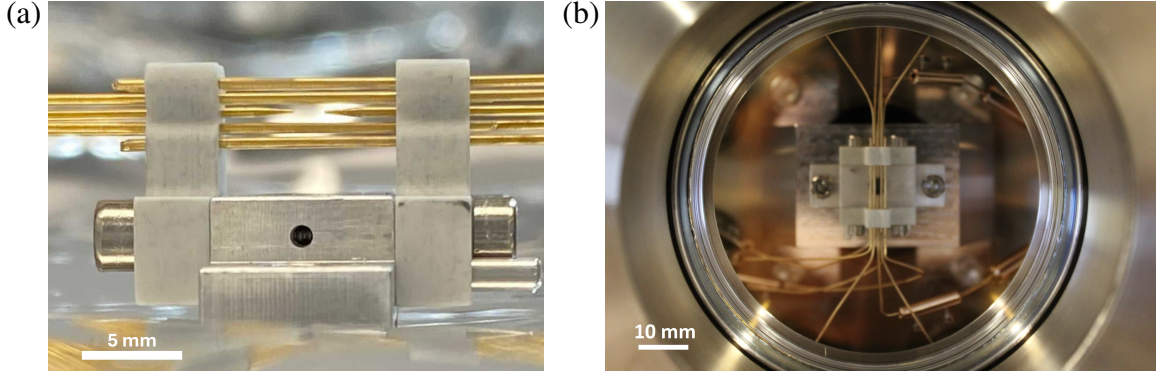


Figure 2.2. (a) Assembled rod trap attached to the two plates and screwed to the center block. (b) Assembled vacuum chamber with the rod trap visible through the front viewport for the high NA imaging system.

and the static harmonic well in the z direction is

$$\Phi_s = \kappa U_0 \left[z^2 - \frac{1}{2} (x^2 + y^2) \right] = \frac{m}{2q} \omega_z^2 \left[z^2 - \frac{1}{2} (x^2 + y^2) \right], \quad (2.2)$$

which is anti-confining in the x and y directions. V_{RF} is the amplitude of the radio frequency (RF) voltage, Ω_T is the RF frequency, U_r is a DC quadrupole voltage, x and y are the radial positions, z is the center axis position along the z -direction (see Fig. 2.1b), r_0 is the distance from the center of the trap to the electrode surface, μ and κ are geometric factors specific to the trap design, U_0 is the static voltage on the center endcap electrodes or ‘needles’ (the pointed electrodes in Fig. 2.1b), m and q are the mass and charge of the particle, and the axial secular frequency

$$\omega_z = \sqrt{\frac{2\kappa q U_0}{m}} \quad (2.3)$$

is the frequency of the oscillation along the trap axis [45]. Eq. (2.1) is the potential that results in the ponderomotive force that keeps the ions trapped in the radial dimensions.

The transverse equations of motion for particles near the center axis of an ideal quadrupole potential are the Mathieu equation with solutions to first order in a_i and

second order in q_i :

$$u_i(t) = A_i \left(\cos(\omega_i t + \phi_i) \left[1 + \frac{q_i}{2} \cos(\Omega_T t) + \frac{q_i^2}{32} \cos(2\Omega_T t) \right] + \beta_i \frac{q_i}{2} \sin(\omega_i t + \phi_i) \sin(\Omega_T t) \right), \quad (2.4)$$

where u_i is the x or y coordinate, A_i is determined by initial conditions, ϕ_i is a phase offset,

$$\beta_i \approx \sqrt{a_i + \frac{q_i^2}{2}}, \quad a_x = \left(\frac{4q}{m\Omega_T^2} \right) \left(\frac{U_r}{r_0^2} - \kappa \frac{U_0}{z_0^2} \right), \quad a_y = - \left(\frac{4q}{m\Omega_T^2} \right) \left(\frac{U_r}{r_0^2} + \kappa \frac{U_0}{z_0^2} \right),$$

$q_x = -q_y = 2qV_{RF}/(\Omega_T^2 m r_0^2)$ is the stability parameter, and

$$\omega_i = \beta_i \frac{\Omega_T}{2} \quad (2.5)$$

is the frequency of the oscillation in the pseudopotential, also known as the secular motion. In addition, there is a component of the overall motion synchronous with the oscillating electric field called the ‘micromotion’ [47]. It is represented in Eq. (2.4) by the oscillating terms with frequency Ω_T .

When $a_i \ll q_i^2 \ll 1$, $U_r \approx 0$, and we neglect the micromotion, $\beta_i \approx q_i \Omega_T / (2\sqrt{2})$ and the time-averaged pseudopotential has approximately the harmonic form

$$q\Phi_p = \frac{1}{2} m \omega_i^2 (x^2 + y^2), \quad (2.6)$$

and the radial secular frequency [48]

$$\omega_i \approx \mu \frac{q_i \Omega_T}{2\sqrt{2}} = \mu \frac{q V_{RF}}{\sqrt{2} \Omega_T m r_0^2}. \quad (2.7)$$

The potentials from the four additional, outer rods, the compensation rods, provide offsets to the trapping field and compensate for stray DC fields that would otherwise move the ion out of the RF null.

2.2 Rod Traps

Real traps approximate ‘ideal’ traps which would use perfectly hyperbolic electrodes extending to infinity that perfectly match the shape of the equipotential lines of a quadrupole trapping field. We use a rod trap because it provides more optical access and is easier to fabricate than hyperbolic traps. When calculating the parameters of non-ideal traps, the geometric factor μ in Eq. (2.1) accounts for the inefficiencies of the non-hyperbolic equipotential lines that trap the ions. It is calculated as the ratio of the actual radial potential and the ideal radial potential at a given trap frequency [49]. The geometric factor κ in Eq. (2.2) accounts for the shape and separation of the needles. The trap parameters, which were chosen based off of work described in Sections 2.3 and 2.4, determine the geometric factors. The geometric factors were calculated after construction of the trap using COMSOL Multiphysics™ software and verified by measuring the axial and radial frequencies of trapped ions as described in Section 2.4. The constructed rod trap can be seen in Fig. 2.2a² and its placement in the vacuum chamber can be seen in Fig. 2.2b.³

Beryllium copper was chosen as the material for both the rods and needles due to its stiffness, thermal conductivity, and lack of magnetism. The rods and needles are gold plated to prevent oxidization, which prevents the rods from building excess charge. A layer of palladium separates the gold and beryllium copper to prevent the gold from diffusing into the rods.

To maintain the spacing and alignment of the rods and needles, two plates made of an electrically nonconducting, ceramic material called Shapal™ are held around a

²A dowel pin is sticking out in the bottom right of the image below the bolt holding the sides and center pieces together. This is because I made the tolerances too small in the CAD file before sending it out to be manufactured. Oops!

³Lab tours available upon request.

center aluminum block. It is made of a combination of aluminum nitride and boron nitride. Shapal™ was chosen for being readily machinable and having relatively high thermal conductivity, low dielectric loss, and high dielectric strength compared to similar materials.⁴ The dissipation of thermal energy from the rods is important to reduce heating of the ions caused by noisy electromagnetic fields [50], known as anomalous heating [51]. Optical access for the imaging system was made via a hole in the center aluminum block, which is visible behind the electrodes in the center of Fig. 2.2b.

The trap itself must be in ultra-high vacuum (UHV) (on the order of 10^{-11} Torr) to reduce the likelihood of ions colliding with background gas particles, which will heat and melt Coulomb crystals formed by the ions [52]. This is done by placing the trap in a stainless steel chamber under active pumping using an ion pump and non-evaporable getter (NEG) pump. The constructed system was baked in a large oven at 190 °C for two weeks to remove water and organic molecules. Fused silica viewports around the chamber allow for optical access.

2.3 Numerical Aperture

Imaging systems are used to capture light emitted by a trapped ion. Captured photons can be used for qubit readout and diagnostic imaging of trapped ions as well as remote qubit entanglement [53]. An imaging system uses an objective, which is a system of lenses, to focus and direct light to other components. There are currently two objectives available for the ion trap. The smaller one was constructed in our lab and is for qubit readout and diagnostic imaging (see A. Quinn’s dissertation [33] for more details on how we image and count trapped ions). It covers $\sim 4\%$ of the total solid angle around the trapped ion and is currently in use on our trap. The larger one that has yet to be installed on the trap was purchased from the company Photongear and covers 10% of

⁴[Product page for Shapal™ via Precision Ceramics.](#)

the total solid angle. It will allow for high efficiency photon collection and can be used to study ion-photon entanglement and quantum network applications such as quantum key distribution, quantum teleportation, and secure quantum computing [54–56].

A useful metric for imaging is the numerical aperture (NA), which determines the fraction of photons emitted by the trapped ions that are captured by the imaging system. The NA is defined as $NA = n \sin \theta$, where n is the index of refraction of the medium the light travels through before the lens (1 for the UHV of the chamber and ~ 1 for the air outside the chamber) and θ is the half-angle of the maximal cone of light that can enter the lens and be directed through the optical system. The larger objective has an $NA = 0.6$ and the smaller one has an $NA = 0.41$.

This simple definition of the NA does not account for the fact that the electrodes may effectively reduce the NA. The limiting factors for determining the effective NA are the objective NA and the parameters r_0 and r_e of the trap itself (see Fig. 2.1a). When looking from the side of the trap as in Fig. 2.3b, the rods above and below the trap center will block optical access within the lens NA. To account for this, the effective NA, NA_{eff} , is defined by finding the unblocked area of the base cone of light due to the rods such that

$$NA_{eff} = \sin \left[\arctan \left(\frac{\sqrt{A\gamma/\pi}}{\cos \phi} \right) \right], \quad (2.8)$$

where A is the area of the projected base of the cone with half angle θ in a system $NA = \sin(\theta)$ and there is an unobstructed line from the tip of the cone at the ion location to the objective at angle θ' as diagrammed in Fig. 2.3b. Because NA is related to the solid-angle covered by the objective lens, *i.e.*, the *curved* area of the sphere centered at the ion, calculating the blocked area of the projected *flat* base of the cone with half-angle θ only approximates the unblocked solid angle. However, the angle difference $\Delta\theta = \theta - \theta'$ is less than one degree for the range of trap parameters being considered making this definition of NA_{eff} a useful approximation. γ represents the ratio of the area of the base

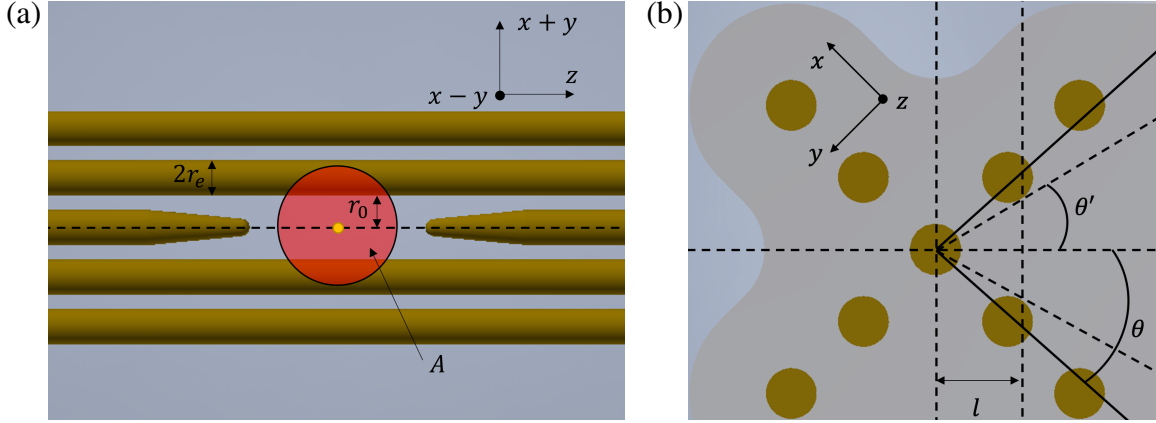


Figure 2.3. Diagram and definitions of geometry for determining NA_{eff} . (a) View of the rod trap from perspective of the objective lens. (b) Side view of the rod trap.

that is unobscured by the rods to the full area A and ϕ is defined using trap parameters.

Using the trap parameters r_e and r_0 as well as the objective NA:

$$\begin{aligned}
 A \gamma &= \pi R^2 - 2 \left[R^2 \arccos \left(\frac{\sin \phi}{R} \right) - \sin \phi \sqrt{R^2 - \sin^2 \phi} \right], \\
 R &= \cos \phi \cdot \frac{NA}{\sqrt{1 - NA^2}}, \\
 \text{and } \phi &= \left(\frac{\pi}{4} - \arcsin \frac{r_e}{r_e + r_0} \right).
 \end{aligned} \tag{2.9}$$

Increasing the trap radius r_0 as well as decreasing the electrode radius r_e will increase the NA_{eff} to a maximum of NA by decreasing the blocked area of the projection of the cone of light at the rod positions.

Choices of beryllium copper rod sizes available for purchase limits the selection for r_e . Too small of a radius reduces the rigidity of the rod and requires it to be kept under tension to maintain the trap shape and requires higher voltages to create equivalent trapping potentials. Too large will block more of the NA. Balancing these constraints, we use $r_e = 0.25$ mm. The trap radius r_0 similarly must balance voltage limitations with NA_{eff} , which will be discussed in the following section.

2.4 Simulations

Limitations on the RF and DC voltages (less than 1000 V for each to prevent damage on vacuum feedthrough dielectric materials) and optical access were used to inform what the trap parameters should be. Boundary element simulations of the trapping potential and ion trap in COMSOL were used to determine the optimal trap dimension values given our optical access needs and available voltages. The trap dimensions were chosen based on the analysis in Section 2.3 with trap radius $r_0 = 0.75$ mm and electrode radius $r_e = 0.25$ mm.

The rods and needles are held in place by friction by the holes in the trap plates. The needles tip-to-tip spacing a can therefore be adjusted to any length less than the plate spacing of 8.5 mm. Even though calcium ions sit microns apart in the trap, the needles cannot be brought arbitrarily close together. If a is too small, the needles will shield the ions from the potential produced by the electrode rods. The needles also cannot be too far apart, as the voltage required to maintain a trapping frequency in the axial direction increases with a^2 . The DC voltages needed to produce various axial frequencies were calculated using COMSOL and are plotted in Fig. 2.4. Based on these calculations and considerations, $a = 3.0$ mm was chosen as it requires reasonable DC voltages for a lab setting.

The needles were secured into the plates by hand and measured to be $a = 3.1 \pm 0.1$ mm apart (see Fig. 2.2a) using a known manufactured value (the width of the plates) as a reference with uncertainty due to the resolution and off centered-ness of the image. Direct measurement of the needle spacing was not done due to the fragility of the needle tips and risk of bending the beryllium copper rods.

In general, $q_i < 0.4$ is desirable for trapping as the harmonic approximation of the trapping potential matches closely to the exact solution in this regime. The trapping

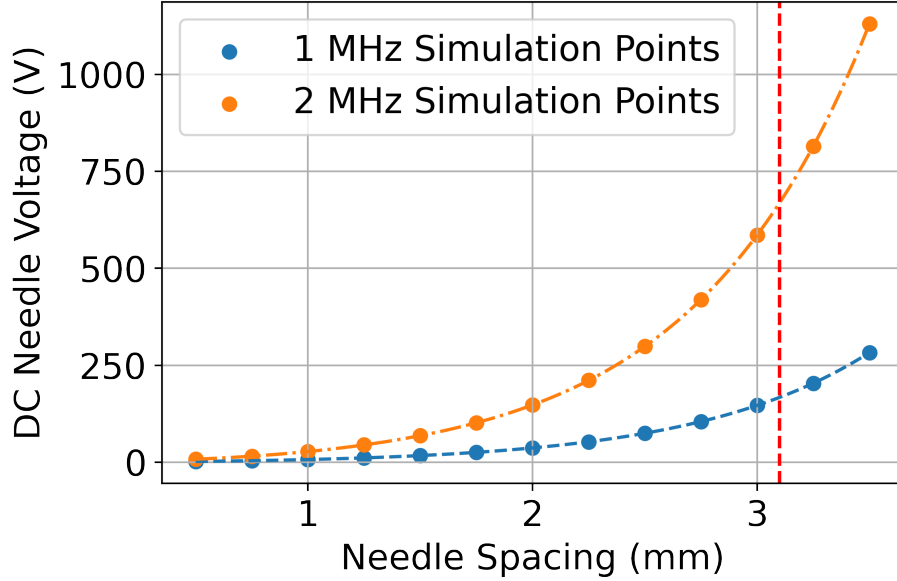


Figure 2.4. Plot of the DC needle voltage needed to produce the labeled oscillation frequencies as a function of the tip-to-tip needle spacing. The vertical dashed line indicates the measured value for the needle spacing.

RF is delivered to the trap electrodes by stepping up the $50\ \Omega$ RF source to high voltage via a helical resonator acting as a transformer [32], as shown in Fig. 2.5. This RF drive is amplitude stabilized by the “Squareatron” [57] to improve the motional coherence of the ion’s radial modes. A range of values for the radial frequency is presented in Fig. 2.6 as a function of the electrode voltage V_{RF} and a given value of q_x . Corresponding values of the RF frequency Ω_T are given as a function of V_{RF} and q_x as well. While both the RF voltage and frequency effect the secular frequency (Eq. (2.7)), because our helical resonator is resonant at 14.13 MHz and designed to maintain high voltages [58], it is preferable to adjust the RF voltage to change the secular frequency.

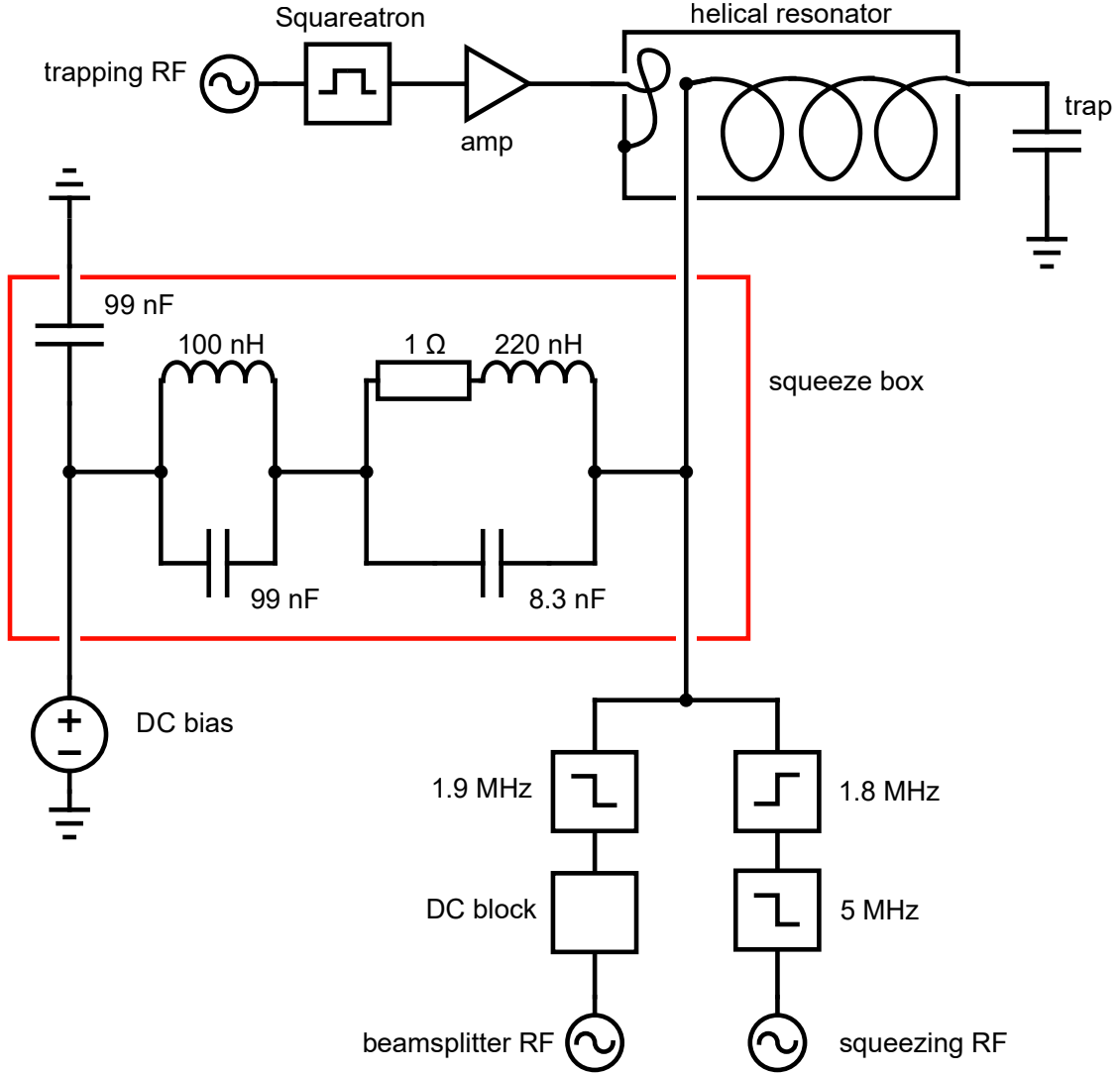


Figure 2.5. Electrical circuit diagram of the trapping RF, DC bias, and parametric driving RF signals. The beam splitter RF is filtered by a Mini-Circuits SLP-1.9+ MHz low-pass and a DC block BLK-18W-S+; the squeezing RF is filtered by a Mini-Circuits SLP-5+ MHz low-pass filter and a custom made 5th order elliptic 1.8 MHz high-pass filter.

Fig. 2.7 plots the relationship between r_0 and the V_{RF} required to maintain a given secular frequency. The value of the geometric factor μ changes as $(r_0/r_e)^2$ [49] and is included in the calculations. As the rods spread farther apart, the voltage must increase

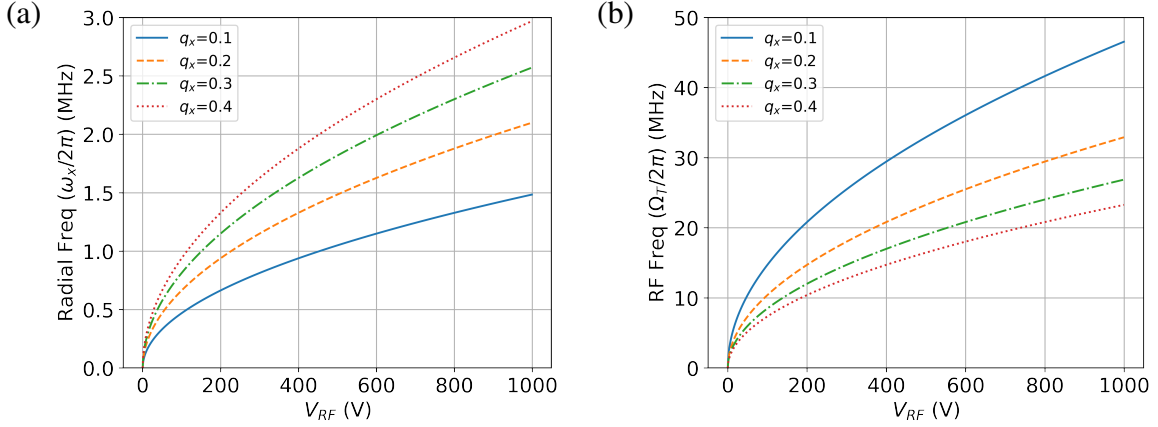


Figure 2.6. (a) A calcium-40 ion's radial secular frequency as a function of the RF voltage applied on the electrodes for fixed q_x values and floating RF frequency. (b) RF frequency Ω_r for a given q_x as a function of the RF voltage. Parameters used: $r_0 = 0.75$ mm, $r_e = 0.25$ mm, $m = 40$ amu (calcium), and $\mu = 0.90$.

to maintain the same trap potential. Beyond 1 mm, the voltage reaches values that would be difficult to produce in a lab setting. The second vertical axis of Fig. 2.7 gives values of the NA_{eff} for the larger NA objective for the corresponding r_0 values. At 1.5 mm, the electrodes no longer block the NA. At 0.75 mm, only 7% of the light is blocked that would otherwise enter the imaging system using the large NA system ($NA_{eff} = 0.58$) while still maintaining a reasonable voltage below 500 V. This is equivalent to covering over $\sim 9\%$ of the total solid angle surrounding the trapped ion. The smaller NA system is not blocked by the rods with these trap dimensions, preserving the full $NA = 0.41$.

With the trap parameters chosen and the rod trap constructed, simulations were performed to calculate the geometric factors μ and κ . Using Eq. (2.3) with the simulated axial frequency from our COMSOL model and the measured trap needle spacing of $a = 3.1 \pm 0.1$ mm, the axial geometric factor was calculated to be $\kappa = 4.8 \pm 0.7 \times 10^4 \text{ m}^{-2}$. After trapping a single ion, we measured the axial frequency by applying an RF voltage to the needles and performing spectroscopy at different voltages amplitudes. Fitting this frequency data, plotted in Fig. 2.8, gave a value of $\kappa = 5.2 \pm 0.2 \times 10^4 \text{ m}^{-2}$. The axial

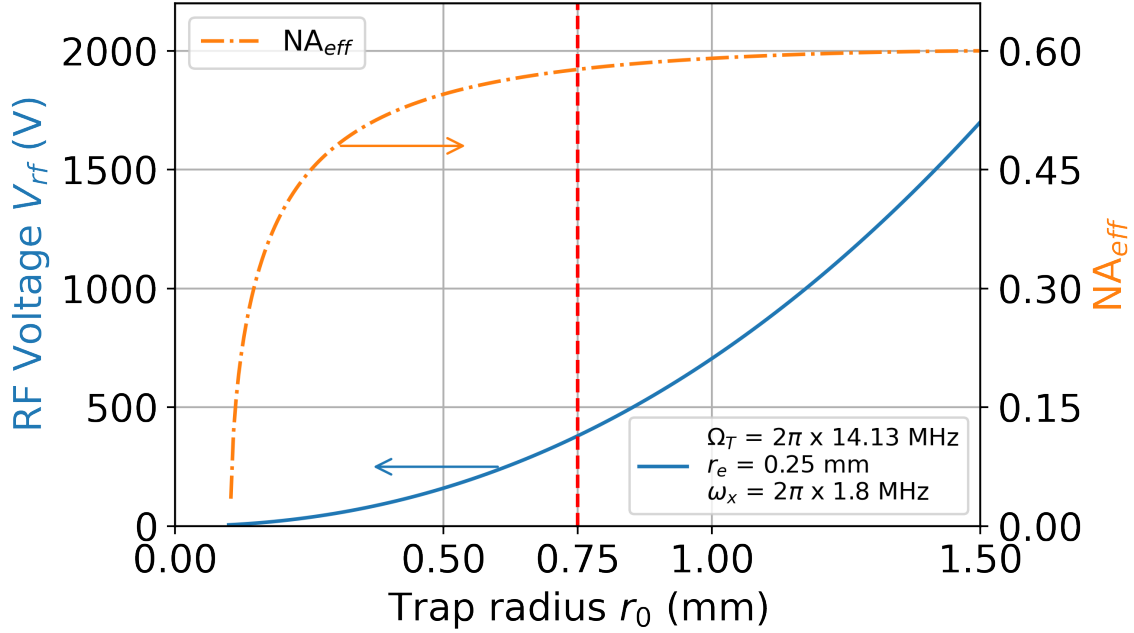


Figure 2.7. Plot of the necessary V_{RF} to trap $^{40}\text{Ca}^+$ ions with the displayed values as a function of the trap radius and corresponding trap NA_{eff} . The dashed vertical line indicates the chosen value of r_0 . Arrows indicate vertical axis associated with lines.

frequency data remains within the uncertainty of κ due to the uncertainty of the measured value of a . The fit has a small, non-zero frequency at 0 V of 8 kHz due to stray fields in the system that keep the ions loosely trapped in the axial direction.

Fig. 2.9 compares the predicted radial frequencies as a function of the RF voltage with collected data. The simulated value of the geometric factor $\mu = 0.90$ was used in Figs. 2.6 and 2.9. Testing produced data consistent at lower voltages and low trapping parameter ($q_x < 0.4$) with a μ of 0.90 and the expected radial frequencies using Eq. (2.7). The first five data points were then fitted to give a step up of 96.0 ± 0.5 of the trapping RF voltage by the RF resonator as shown in the plot.

The threshold for trapping is $q_x \leq 0.9$ [36], where above this value the motion of the ions is unstable. We do not intend to perform experiments for $q_x > 0.4$ as the solutions for our equations for the radial frequencies (Eq. (2.7)) assumes $q_x \ll 1.0$, and

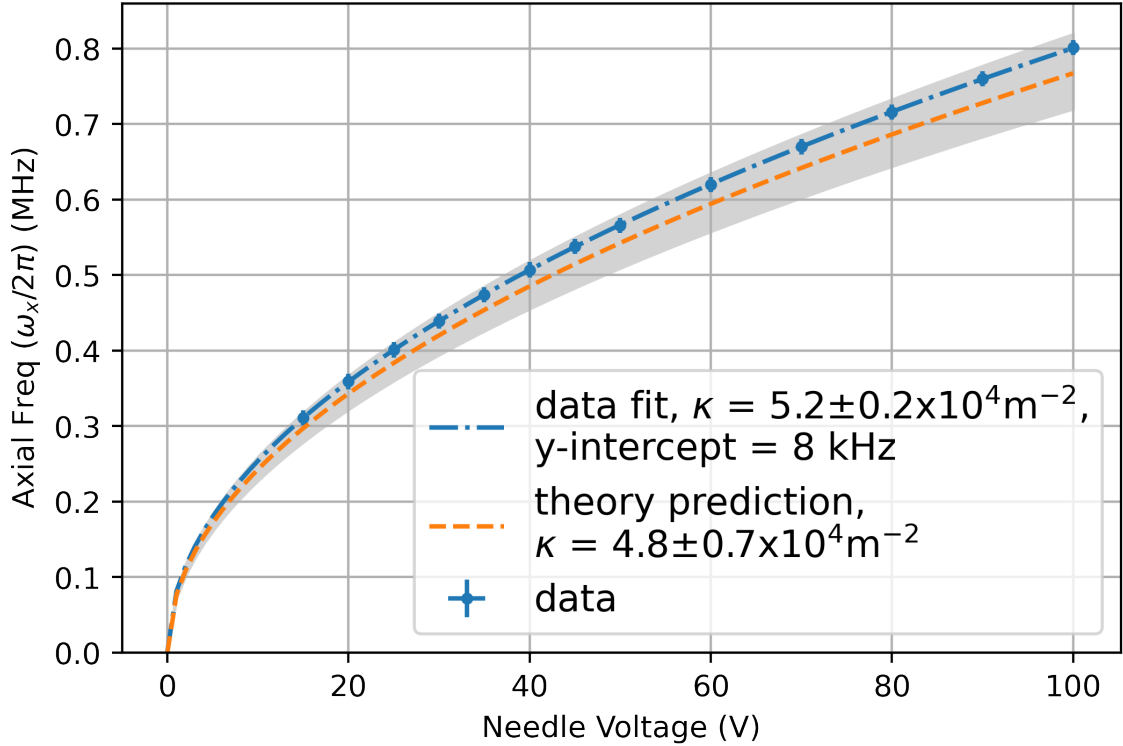


Figure 2.8. Predicted values of axial frequency for the predicted value of κ plotted with measured values of the axial frequency with fitted value of κ . The gray area is due to the uncertainty in the needle distance which places an uncertainty on κ . The frequency error is on the order of 1 kHz and the DC voltage error is 0.0025%. Both are smaller than the data marker. Parameters used: $r_0 = 0.75$ mm, $r_e = 0.25$ mm, $m = 40$ amu (calcium), and $\mu = 0.90$.

as q_x approaches 0.4 the predicted radial frequencies diverge from the measured value.

The radial frequency reaches zero at a non-zero RF voltage due to the interaction with the DC voltages on the needles.

2.5 Laser-Free Control

Many experimental processes involving our trapped ions require the use of laser systems, which are discussed in detail in Chapter IV. There are many processes, however, that can be performed fully electronically.

2.5.1 RF Rabi Flopping. We can induce precession of the ion's magnetic moment with an oscillating magnetic field via an additional RF currents applied to one

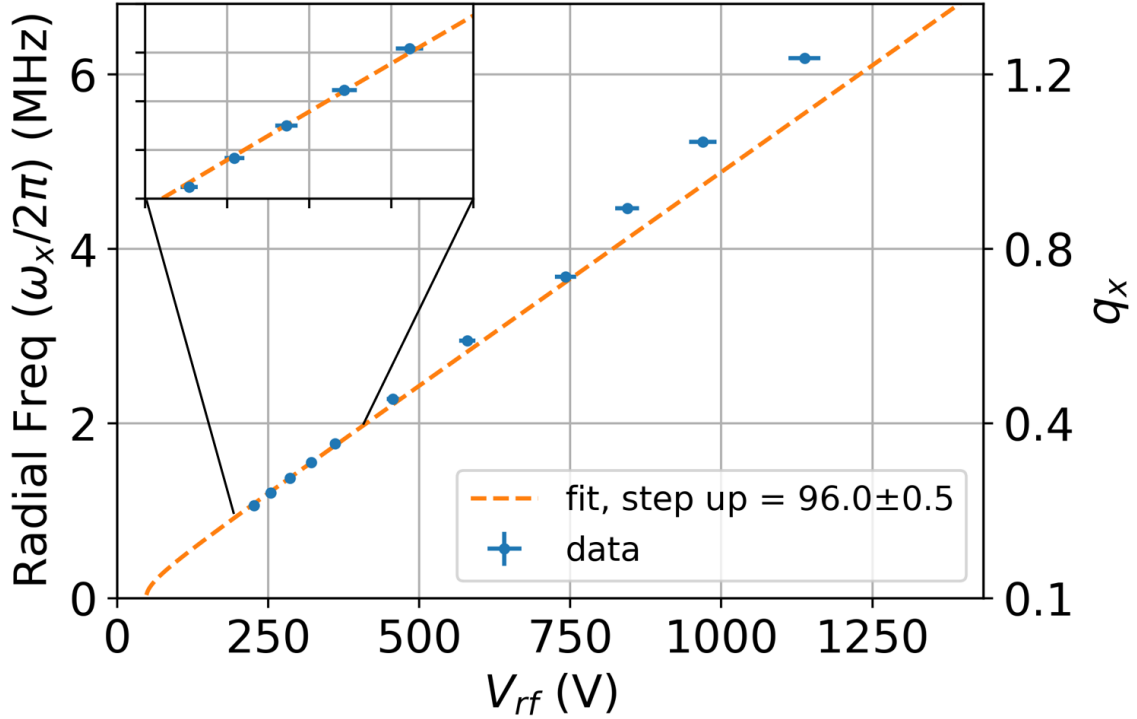


Figure 2.9. Measured values of the radial frequency and corresponding q_x values as a function of V_{RF} . The fit floated the value of the step up of the trapping RF voltage by the RF resonator using the five lowest voltage values (see inset). The frequency error is on the order of 1 kHz and smaller than the data marker. Error on the voltage comes from the non-linearity of the power meter of ± 0.1 dB. Parameters used: $r_0 = 0.75$ mm, $r_e = 0.25$ mm, $m = 40$ amu (calcium), and $\mu = 0.90$.

of the trapping rods. A two-level (effective spin-1/2) system with a magnetic moment aligned or anti-aligned to the quantization magnetic field B_z with an AC magnetic field oscillating at frequency ω perpendicular to the quantization field in the x -direction B_x has the Hamiltonian

$$H = \frac{\hbar\mu B_z}{2}\sigma_z + \frac{\hbar\mu B_x}{2}\cos(\omega t)\sigma_x \quad (2.10)$$

where μ is the magnetic moment of the ion, ω_z is the qubit splitting, and σ_x , σ_y , and σ_z are the Pauli operators. Switching to the interaction picture and applying the rotating-

wave approximation [59], we can get an interaction Hamiltonian

$$\bar{H} \approx \frac{\hbar\Omega}{2} \begin{bmatrix} 0 & e^{i\delta t} \\ e^{-\delta t} & 0 \end{bmatrix}, \quad (2.11)$$

where $\Omega = \mu B_x/2$ is the Rabi frequency and $\delta = \omega - \omega_0$ is the frequency difference between the oscillating AC magnetic field ω and qubit resonance ω_0 .

The state of the system can be written as $|\psi\rangle = c_0|0\rangle + c_1|1\rangle$. In the interaction frame, the coefficients co-rotate with the freely evolving qubit states $\bar{c}_n = c_n e^{i\omega_n t}$, where the populations of each qubit state are $|\bar{c}_n|^2$. Using Schrödinger's equation, the equations of motion can be solved to give the likelihood of finding population in state n at a given time:

$$\begin{aligned} P_0(t) &= \cos^2\left(\frac{\chi t}{2}\right) + \frac{\delta^2}{\chi^2} \sin^2\left(\frac{\chi t}{2}\right) \\ P_1(t) &= \frac{\Omega^2}{\chi^2} \sin^2\left(\frac{\chi t}{2}\right), \end{aligned} \quad (2.12)$$

where $\chi = \sqrt{\Omega^2 + \delta^2}$ is the generalized Rabi frequency for off-resonant driving [59, 60]. This treatment of RF-driven qubits can be applied to laser-driven coherent transitions as well.

2.5.2 CVQC Operations. While discrete quantum computers use qubits such as the electronic states of an atom to store and perform quantum logic gates, a continuous variable quantum computer (CVQC) uses physical quantum systems that can exist as a continuum of values. Moving from discrete computation also represents a move from a finite-dimensional to a theoretically infinite-dimensional Hilbert space. While a CV quantum computer cannot perform any computation that its discrete counterpart cannot, it may be able to perform some more efficiently [61].

Ion traps provide effective systems for CVQC due to their long-lived vibrational modes [61, 62]. These vibrational modes can be coupled to the electronic states via a Jaynes-Cummings interaction [63], which allows for information to be stored as

continuous variables [64, 65]. Information can be stored in ion traps by utilizing the three motional modes of each ion instead of only considering electronic states through the creation of superpositions of the quantized energy eigenstates of the harmonic oscillator in the phase-space of position and momentum [64]. Quantum computation can also then be performed on a finite-dimensional code space embedded within an infinite-dimensional Hilbert space, such as ‘cat’ or ‘GKP’ codes [65, 66]. These codes can be created by using a combination of linear, Gaussian processes [65, 67, 68] and non-Gaussian processes. For universal logic on a continuous variable platform, at least one non-Gaussian operation is required [61, 64, 66, 67], which can be performed using laser fields [69].

A Gaussian process is characterized only by its mean and standard deviation. These include single-mode squeezing, two-mode squeezing, and beamsplitting, all of which can be implemented in ion traps by modulating voltages on the trap electrodes [67, 70]. We apply the voltages required to generate the quadrupole potential for motional squeezing and beamsplitter operations [22, 32] directly to the RF trapping electrodes through the use of our parametric drive electronics (see below). Using the RF trapping electrodes ensures the quadrupole is centered on ions located on the axis of our linear trap.

2.5.2.1 Parametric Drive Electronics. The trapping RF and beamsplitter RF are produced by a Rigol DG1022Z arbitrary waveform generator (AWG) and the squeezing RF is produced by an Urukul DDS card [71], part of our ARTIQ Sinara experimental control system [72]. The voltages are applied directly to the ‘back’ of the helical resonator’s coil (see Fig. 2.5) via bandpass filters that ensure the two sources are isolated from each other. Because the tones are applied to the coil connected directly to the trap and are well below the resonant frequency of the helical resonator, they

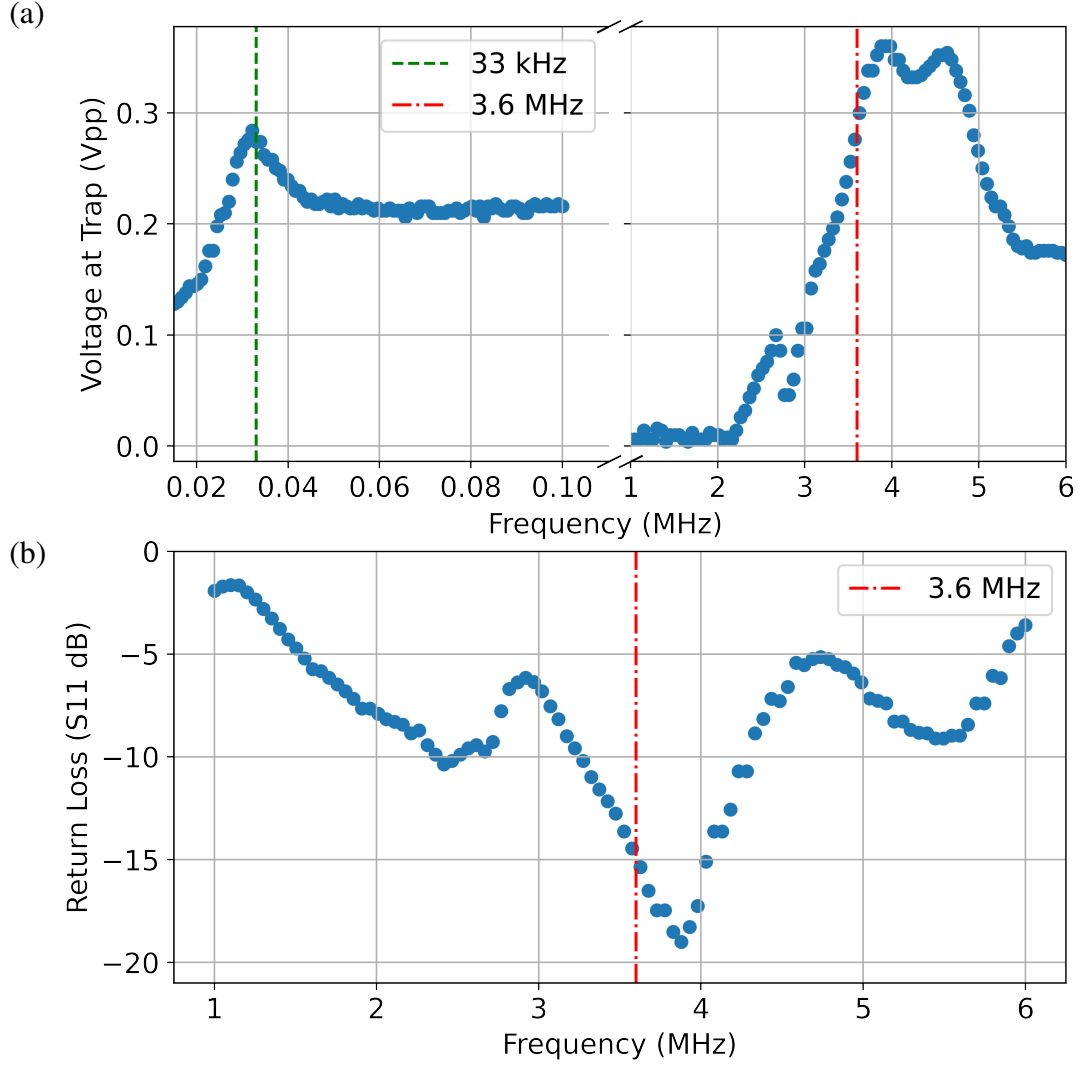


Figure 2.10. Electrical testing of RF driving sources through the squeeze box. **(a)** Voltage measured at the ion trap vacuum feedthrough with beamsplitter and squeezing RF sources outputting 0 dBm of power. Vertical dashed green and dash-dot red lines indicate the beam splitter and squeezing frequencies, 33 kHz and 3.6 MHz, respectively. **(b)** Return loss for the squeezing drive input showing good impedance matching.

pass through to the trap electrodes relatively unimpeded without any coupling to the Squaretron and trapping RF signal source path.

Also connected to the back of the resonator coil is the “squeeze box” containing a pair of LC circuits, which provides three functions:

1. It provides a DC bias path to allow control of the mode rotation and splitting on the trapping rods.
2. It provides a low impedance path to ground at the RF trapping frequency, ensuring there is no effect on trapping conditions. Compared to a direct ground connection, the “squeeze box” decreases the quality factor ($Q = f_0/\text{FWHM}$) of the helical resonator modestly, from 310 to 295, and shifts its resonant frequency of 14.5 MHz by approximately 6 kHz.
3. It provides an impedance match for the two parametric drive tones by having the LC circuits in the box tuned to match the beamsplitter and squeeze drive frequencies of 33 kHz and 3.6 MHz respectively.

Electrical testing shows we get voltages on the RF electrodes close to what we expect from simulations of the RF electronics, as well as low return loss at the squeezing drive frequency (see Fig. 2.10). While the squeeze box allows for a DC bias path on the RF trapping rods, we use the compensation rods to apply DC bias voltages since they are symmetrically aligned to the inner trap rods.

The potential applied at the trap electrodes determines the parametric coupling strength

$$g = \mu \frac{q V_p}{m r_0^2 \omega_i}, \quad (2.13)$$

where μ is the geometric factor of the trap, q is the charge of the ion, V_p is the voltage amplitude of the modulating quadratic potential, m is the mass of the ion, r_0 is the trap

radius, and ω_i is the secular frequency of the motional mode being coupled to. For measured voltages on the feedthrough for the trap RF electrode of 0.26 V_{pp} for the beam splitter and 0.58 V_{pp} for both single and two-mode squeezing we get couplings of $g_{BS} = 2\pi \times 0.66 \text{ kHz}$, $g_{SMS} = 2\pi \times 3.68 \text{ kHz}$, and $g_{TMS} = 2\pi \times 1.09 \text{ kHz}$, respectively. These coupling values are calculated from theory using the above voltages and agree reasonably with the values calculated from experiments presented in [22].

2.6 Future Trap Developments

In the previous section, setup work for our study of two-mode squeezed states in the motional modes of trapped ions [22] was presented. However, further research on the subject of squeezed modes will require larger mode splittings. To meet this need, a second iteration of the ‘Oregon style’ linear Paul trap was designed. A key feature of the first trap is the alignment of the trapping and compensation rods in the ‘X’ formation. This design, without additional voltages, would result in degenerate radial secular mode frequencies in the absence of trap imperfections. Non-degenerate secular mode frequencies are important for driving single and two-mode squeezing independently. The frequency for driving single-mode squeezing is at twice the secular motional frequency $\omega_{SMS} = 2\omega_i$ and the frequency for driving two-mode squeezing is the sum of the two modes’ secular frequencies $\omega_{TMS} = \omega_x + \omega_y$ [22]. The mode splitting, $|\omega_x - \omega_y|$, informs how separable the single and two-mode squeezing operations are from each other. As the driving frequency approaches the resonant frequency of each operation, the greater the parametric coupling strengths (of the same form as Eq. (2.13)) and the greater magnitude of the squeezing operation measured with the squeezing parameters $r_{SMS} = gt$ and $r_{TMS} = 2gt$ where t is the time the driving field is applied. The coupling strengths of each operation are π -periodic with the mode orientation ϕ such that $g_{SMS} \propto \cos(2\phi)$ and $g_{TMS} \propto \sin(2\phi)$. This is the angle of the ions’ secular motion relative to the x and y axes (see Fig. 2.11).

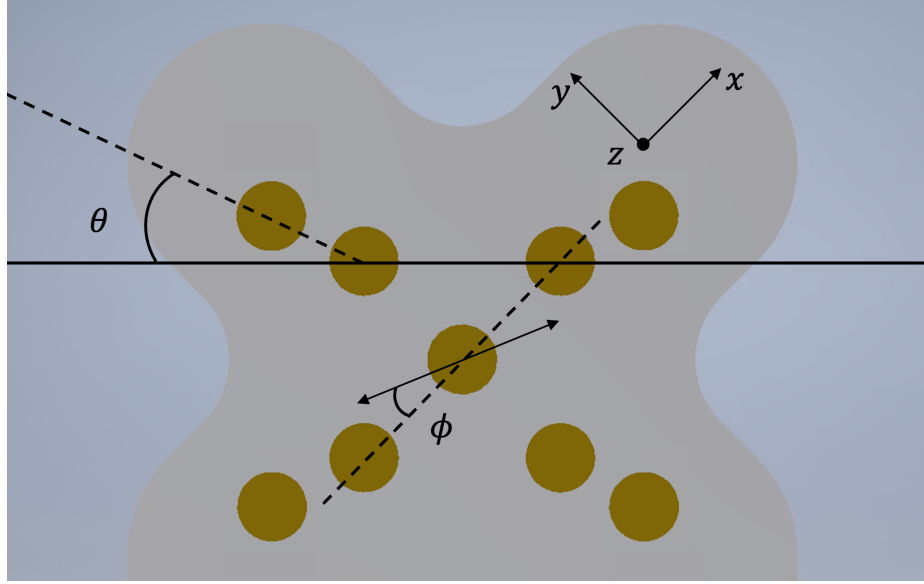


Figure 2.11. Definitions for angle θ between the compensation rods and the trapping rods and the mode orientation ϕ , which is the angle between the direction of the secular frequency (double-headed arrow) and the direction defined by the RF trapping rods in the $x - y$ plane. See Fig. 2.1 for the pattern of voltages applied to the rods.

At $\phi = 0^\circ$ the ions' secular motion will be aligned with the x and y axes and at $\phi = 45^\circ$ anti-aligned. Single-mode squeezing coupling is maximized when aligned with the rods at $\phi = 0$ and two-mode squeezing coupling is maximized when anti-aligned at $\phi = 45^\circ$. In [22], the mode orientation was set to $\phi = 11^\circ$. This allowed for both squeezing operations to be performed.

If the mode splitting is too small, the maximum squeezing magnitude achievable of one operation performed independent of the other will be limited. Adding a bias voltage to the RF rods or creating a DC quadrupole with the compensation rods can increase the mode splitting, but as the voltage increases, the secular motion of the ions are further radially aligned with the axes of the rods. To access both operations in practice with large coupling rates, $\phi = 22.5^\circ$ is preferred. By constraining our parameter space to $\phi \in [22^\circ, 23^\circ]$ our current trap design is limited to a mode splitting of less than 100 kHz.

As can be found in [22], our single-mode squeezing, two-mode squeezing, and beamsplitter operations were successfully achieved and studied. However, further study as well as related research topics, such as spin-squeezed states for quantum metrology [73], will benefit from a greater splitting between our secular mode frequencies without aligning the modes to the trap rods. To stop this alignment, we can change the position of the compensation rods via the ‘trap angle’ θ (see Fig. 2.11). Instead of the ‘X’ formation, the angle of the rods can be adjusted to flatten out the overall shape of the cross-section of the trap, which moves the DC quadrupole out of angular alignment with the RF electrode bias. The current trap design has $\theta = 45^\circ$. Simulation data is limited to mode splittings > 100 kHz, $\phi = 22\text{--}23^\circ$, and RF bias voltage of 0–15 V. The numerical simulation performed gave discrete values when calculating ϕ with too few values at exactly 22.5° . In order to understand trends of the parameter space, values in the range of $22\text{--}23^\circ$ were analyzed. Simulations were done in steps of 5° from 0 to 45° , see Fig. 2.12. There are no presented results for $\theta \geq 30^\circ$ because there were no simulated values for our compensation rod voltages and RF bias voltages that resulted in mode splittings ≥ 100 kHz and ϕ between $22\text{--}23^\circ$.

Simulations reveal that all θ values have clear minimums for mode splitting and a point where the mode orientation ϕ can be adjusted with arbitrarily small changes. See Fig. 2.13 for simulation results for $\theta = 0, 35$, and 45° . Additional simulations were made to verify the geometry of the pseudopotential would not be significantly altered due to the adjusted position of the compensation rods. Fig. 2.14 shows that the geometries are comparable, especially near the center of the trap.

While construction of the second Oregon style trap has yet to begin, a compensation rod angle $\theta \sim 0^\circ$ is the likely design due to the wide range of achievable mode splitting frequencies for $\phi = 22\text{--}23^\circ$.

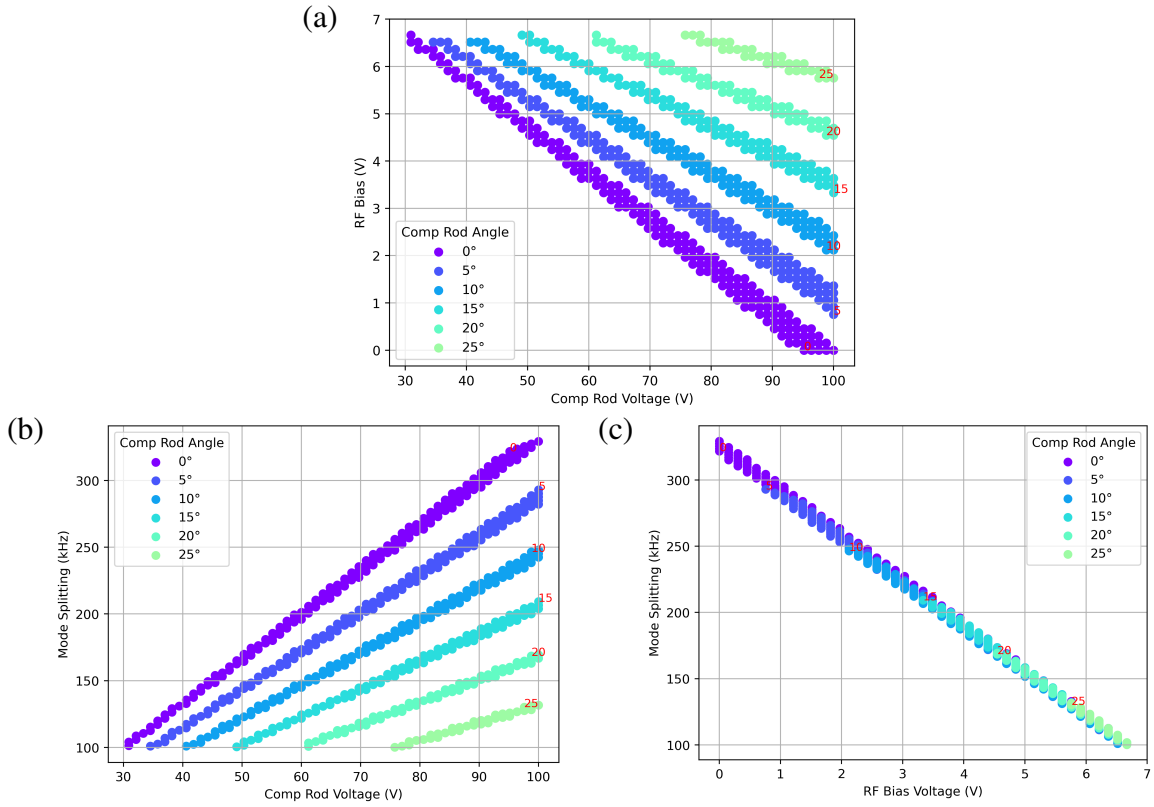


Figure 2.12. (a) Combinations of different compensation rod voltages and RF bias voltages where mode splitting values ≥ 100 kHz and orientation of 22–23°. (b) Mode splitting values as a function of compensation rod voltage for RF bias between 0–15 V and mode orientation of 22–23°. (c) Mode splitting values as a function of RF bias voltage and mode orientation of 22–23°.

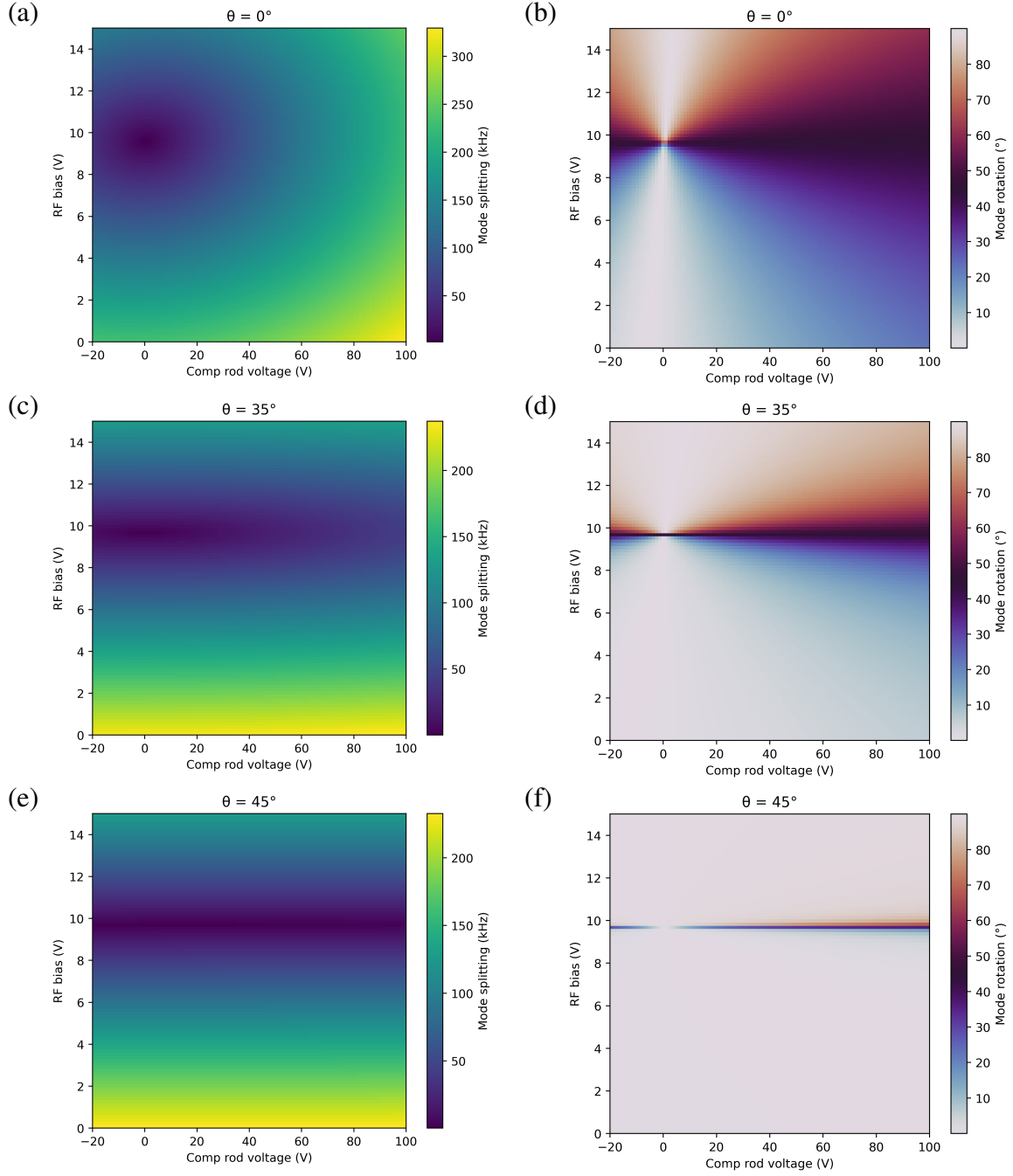


Figure 2.13. Simulations of mode splittings and mode orientations as a function of compensation rod voltage and RF bias voltage for various compensation rod angles.

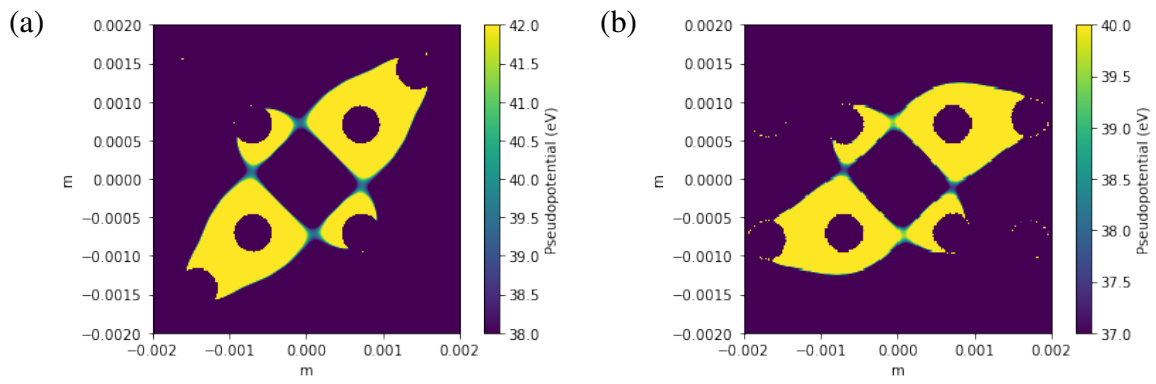


Figure 2.14. Pseudopotential depths for different trap designs. (a) $\theta = 45^\circ$. (b) $\theta = 5^\circ$.

CHAPTER III

LIGHT INDUCED AND THERMAL ATOMIC DESORPTION

The work contained in this chapter was performed by myself under the supervision of Prof. David Allcock, with all writing done by myself with editorial assistance from Dr. Jameson O'Reilly.

The most common way of producing free atoms for an experiment is a thermal source, where bulk metal is heated until the desired vapor pressure is achieved.¹ Typically, an electrical current is used but heating with a CW laser has also been demonstrated [74]. The resulting vapor can either be used directly or collimated into a beam. Experiments can be performed using the beam itself or individual atoms can be loaded into an optical, electromagnetic, or magneto-optical trap. Ion traps typically load directly from the beam as they are deep ($\sim 100\text{--}1000\text{s K}$) and the atoms can be ionized at the bottom of the trapping potential. In contrast, neutral atom traps often require a pre-cooling stage such as a Zeeman slower.

In laser ablation loading, a laser pulse with energy on the order of 10 mJ and length $< 10\text{ ns}$, often provided by an Nd:YAG laser, is directed onto a material that can be either refined metal, an alloy, or a compound that contains the target species [75]. The laser pulse vaporizes or sublimates the top layer of the material, including the target atoms [76–78], which tend to escape in a plume perpendicular to the substrate surface.

While both methods reliably produce free atoms, most will have too high of a kinetic energy to be trapped by the shallow pseudopotential of surface traps commonly used for trapped ion quantum computing ($< 100\text{ K}$) [79]. Thermal sources produce atoms with a Maxwell-Boltzmann velocity distribution centered at a temperature as low as $\sim 100\text{ }^{\circ}\text{C}$, while laser ablation produces atoms with velocity distribution much higher,

¹“Until now, this was the only way to get juice out of an orange.” — Troy McClure, *The Simpsons* S4E21 “Marge in Chains”

anywhere from 100s to 1000s of degrees hotter [80]. As the peak temperature of the freed atoms increases, the percentage of untrapped atoms increases and more atoms will need to be released into the pseudopotential to consistently trap. Additionally, ablation lasers cost on the order of tens of thousands of dollars and increase the experimental complexity. Ovens take time to warm up and cool down and release background gases such as hydrogen. The heat load can also cause problems in cryogenic systems [81].

Light induced atomic desorption (LIAD) is an alternative method for free atom production that may overcome these issues. LIAD works by taking advantage of an interesting feature of non-resonant light interacting with bulk materials; light above a minimum intensity can increase the vapor pressure of a bulk solid in vacuum [82]. Interestingly, this may partially explain how sunlight causes the moon to have an atmosphere [83]. LIAD produces atoms at room temperature [84], which reduces the number of atoms that need to be liberated for trapping. Another form of desorption is thermal desorption, which also requires a high power CW laser incident on bulk material but produces free atoms at temperatures similar to oven sources. While LIAD, thermal desorption, and ablation all require an additional laser over an electrical oven, a desorption laser can be cheaper and more compact with a 100 mW CW diode laser costing \$100 being sufficient.

LIAD has already been shown to work for alkali metals [85–88], ytterbium [89], and even heavy, radioactive elements used to improve the measurement of the atomic mass of uranium-238 by two orders of magnitude [90] in an electron beam ion trap [91], but laser desorption has not been explored extensively for calcium. One study on desorbing calcium looked at photo-ejecting a thin calcium layer on a polymer substrate called polydimethylsiloxane (PDMS) [92], but this method has limited utility for many experiments as PDMS is not ultra-high vacuum (UHV) compatible.

One alkaline earth metal that has been studied in the context of LIAD is strontium. A bulk piece of SrO, as well as a paste made from mixing SrO with acetone, was placed in a vacuum chamber with a laser directed at them. Both targets produced a vapor pressure and measurable fluorescing signal [82]. The following parallels this work to extend these results to calcium and further studies thermal desorption as an option for calcium-based atomic experiments.

We tested several substrate surfaces along with several wavelengths of light to eliminate candidates and find a best combination. We also attempted to determine the temperature of the desorbed calcium, but our method for extracting the temperature from the Gaussian broadened fluorescence signal was not sensitive or consistent enough.

3.1 Experimental Setup

3.1.1 Vacuum chamber. A steel vacuum chamber was kept under active vacuum pumping with an attached turbo pump that achieved pressures on the order of 10^{-8} mbar. At the center of the chamber, an oven filled with calcium metal was placed over an aluminum shelf with a forked end that supports a flat substrate, see Fig. 3.1. The oven was used to deposit calcium films on the substrate below and to provide a reference atomic beam for alignment checks. The desorption beam was directed into the chamber through a fused silica viewport from above so that it passed through the substrate and between the fork prongs through the bottom viewport of the chamber to a photodiode (PD).

3.1.2 Fluorescence Collection. A 423 nm beam aimed between the oven and substrate drove the optical cycling $4^1S_0 \leftrightarrow 4^1P_1$ transition in any calcium atom emanating either from the oven or the substrate surface. Perpendicular to the path of the fluorescence beam was an imaging system (pictured in Fig. 3.2) with a high effective NA of 0.3 ending

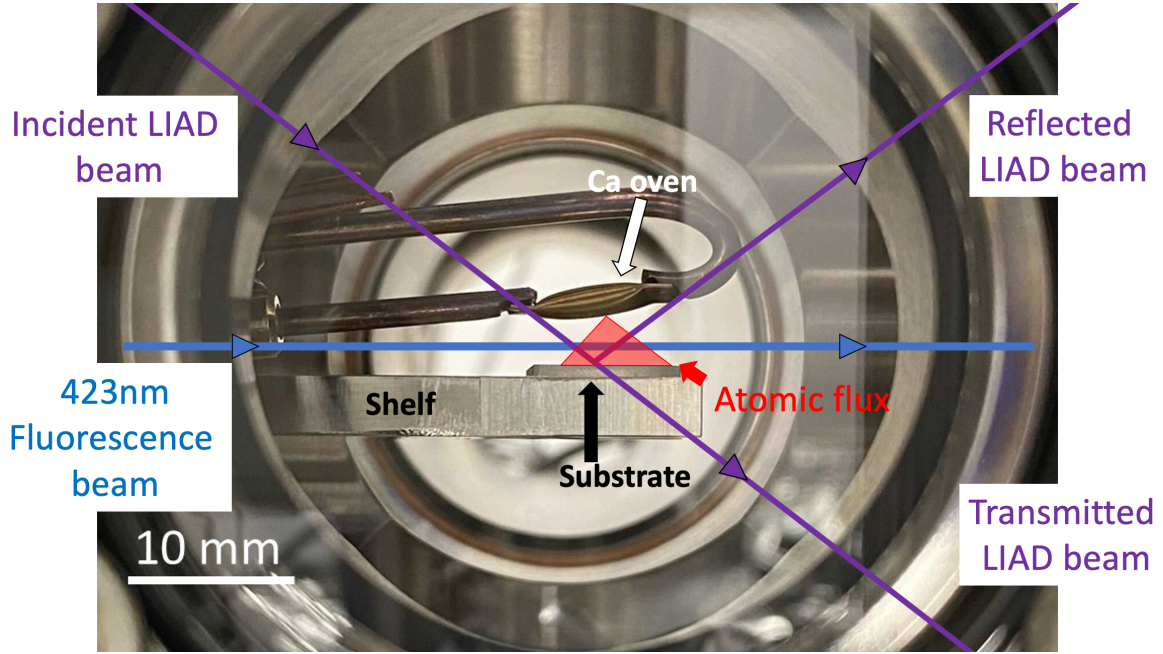


Figure 3.1. Schematic of laser and calcium interactions inside the vacuum chamber. The incident and transmitted laser beam powers were measured by PDs outside the chamber to determine the energy absorbed by the calcium and substrate.

with a photomultiplier tube (PMT). The intersection of the field-of-view (FOV) of the imaging system and the 423 nm laser is referred to as the interaction zone.

The density of atoms in the interaction zone due to a heated oven [93] is $N = \frac{N_0 a}{4\pi d^2}$, where the atom number density inside the oven $N_0 = \frac{p}{k_B T}$ with p being the partial pressure of the chamber, k_B the Boltzmann constant, T the temperature of the oven, a the area of the oven orifice, and d the distance from the orifice to the interaction region. This density can be multiplied by the FOV of the imaging system and the 423 nm beam width to determine the number of atoms N_e in view of the imaging system. Assuming beam intensity below saturation, the corresponding count rate of photons is $S = \eta N_e A \rho$, where η is the imaging system efficiency, $A = 2.23 \times 10^8 \text{ s}^{-1}$ is the Einstein coefficient of the fluorescing transition, and $\rho = 3/2$ is a factor to take into account polarisation effects in the detection system and the angular distribution of the scattered light [74, 94].

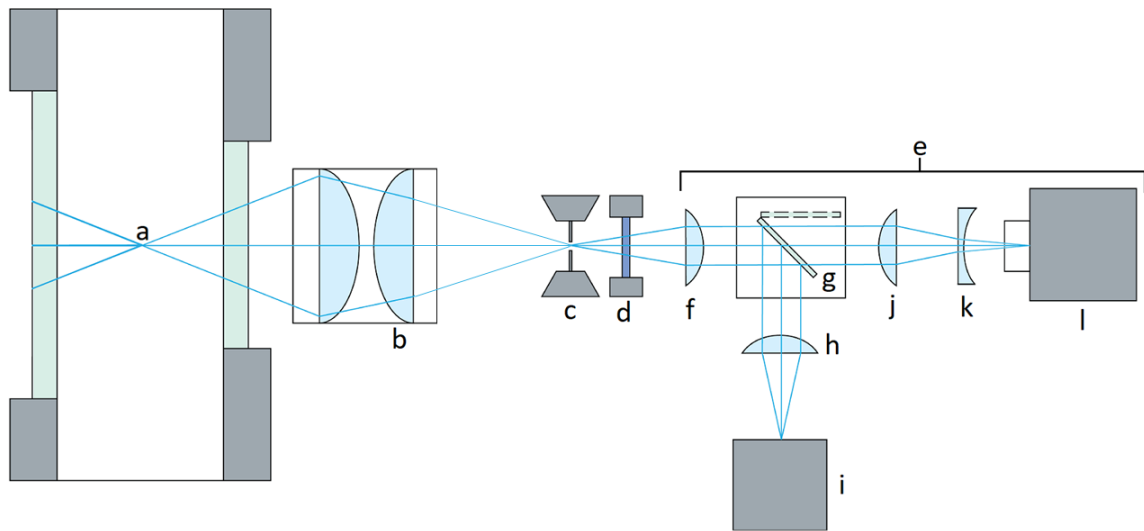


Figure 3.2. Imaging system schematic. (a) Interaction zone inside the vacuum chamber. (b) High NA objective lens. (c-h, j, and k) Second-stage optical components for directing photons to the (i) PMT and (l) CMOS camera.

Using the ion trap in the Oregon Ions lab and comparing setups in [93] and [74], we determined that a loading time of approximately one second using a 375 nm ionizing laser intensity of 22 W/cm^2 ($700 \mu\text{W}$ of power and beam waist of $32 \mu\text{m}$) requires a density of 100 atoms/cm^3 in the interaction zone. This density corresponds to as little as 500 counts/s above background using the PMT. Such a small signal is difficult to distinguish in this setup, so a higher threshold of 5,000 counts/s was set to confirm the presence of calcium in the interaction zone.

3.1.3 Substrates. The choice of substrates was constrained by UHV compatibility. We performed measurements on 0.1 mm and 1 mm thick borosilicate glass, 0.1 mm and 1 mm thick fused silica glass, 0.5 mm thick C plane oriented sapphire, 1 mm thick (111) crystal orientated CaF_2 , 100 nm thick gold film on 1 mm thick soda lime glass, 1/16" thick MacorTM ceramic, 1 mm thick isomolded graphite plate, and $25 \mu\text{m}$ thick graphene sheets. Only testing with borosilicate glass, fused silica glass, and graphene

Table 3.1. LIAD laser table

Wavelength (nm)	Max Power (W)	Max Intensity (W/cm ²)
405	0.16	5809
450	1.28	4074
808	2.4	105

sheets resulted in any measurable data distinguishable from background so these will be the only substrates discussed moving forward.

We also looked at desorbing Ca atoms from calcium oxide (CaO) due to the success reported by Kock with SrO [82]. CaO was made by leaving refined calcium metal in open air to oxidize for several days. The CaO was then crushed with a mortar and pestle and acetone was mixed in until a ‘slurry’ with the consistency of crushed ice drink or frozen margarita was formed. The slurry was then spread across a sample substrate and placed on a hot plate to evaporate the acetone. The substrate with the dried CaO was then placed into the test chamber. Attempts at extending the SrO work described in [82] to CaO were unsuccessful and none of the results presented here involved desorbed Ca from CaO.

3.1.4 LIAD Lasers. Seeing as one goal of laser desorption for loading atomic systems is to use cheaper and easier to use lasers, the choice of wavelengths tested were determined by availability of simple diode lasers that cost only a few hundred dollars and produce powers in the range of 100 mW to 3 W. This included wavelengths of 405, 450, and 808 nm. See Table 3.1 for LIAD laser details.

In previous works, wavelengths beyond green did not result in desorption due to a LIAD effect [95], with the majority of work being done in the near-UV, blue, and green wavelengths [82, 85–88]. This suggests that any signal coming from red wavelengths is a thermal effect and can be used to compare against the shorter wavelength lasers.

3.2 Methods

3.2.1 Temperature Measurements. Desorption signals were compared to signals produced by the oven while depositing calcium onto the target substrate in the testing chamber. The oven was heated by sending a current of up to 10 A through electrical feedthroughs into the chamber. When fluorescing calcium we expect a frequency spectrum fitting a Voigt profile: a convolution of the Lorentzian spectral line shape with linewidth γ and Gaussian line shape of width σ due to Doppler broadening from the high temperature of the atoms.

To perform spectroscopy, the 423 nm laser diode was scanned ± 3 GHz around the transition frequency between the 1S_0 and 1P_1 state of neutral calcium-40 [96]. The number of photons detected as a function of the scanned laser frequency is graphed in Fig. 3.3. Similarly, the desorption signal was measured while using the non-resonant laser light to desorb calcium deposited on the substrate and allowing it to pass through the interaction zone.

Theoretically, the temperature of the fluorescing calcium could be calculated from the Gaussian linewidth

$$\sigma = \sqrt{\frac{k_B T}{m c^2}} f_0, \quad (3.1)$$

where m is the mass of a calcium atom, c is the speed of light in vacuum, and f_0 is the frequency of the emitted photon, i.e., the transition frequency from the 1S_0 to 1P_1 states. In practice, the Gaussian widths extracted from these signals did not correspond to consistent temperatures. Often the signals from the oven had similar or smaller σ values and temperatures than the desorption signals which is not expected for a LIAD signal. The oven signals also spanned a nonsensical range of temperatures including below 0 °C, which puts into question the implementation of this method of extracting temperature. See Fig. 3.3 for a comparison of oven and desorption signals to each other. The shifts in

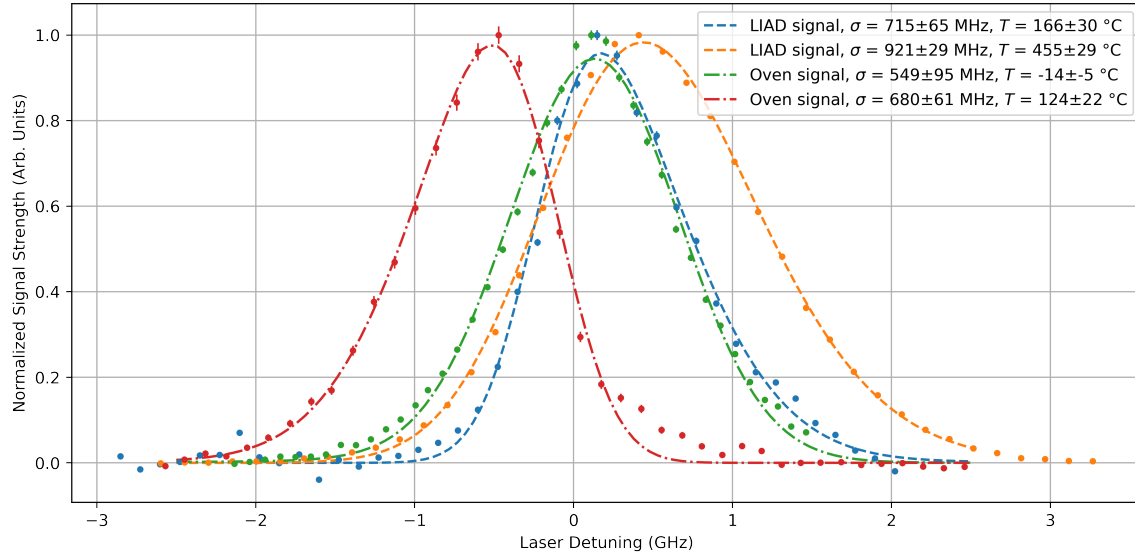


Figure 3.3. Comparison of signals from oven and laser desorption (LIAD) using 450 nm laser light and fitted to Voigt profiles with Gaussian width, σ .

the peaks 100s of MHz apart suggests an error in these measurements for which we do not have an explanation.

3.2.2 Thin Film Measurements. To determine how much material is removed from substrates, the transmitted power of the LIAD beam through the calcium films and transparent substrates was measured. The opening in the center of the shelf supporting the substrate in the chamber allowed for the transmitted beam to pass fully through the vacuum chamber to a photodiode. Measuring the thickness of deposited calcium where the desorption signal was detected also would help determine how important the defect site theory is [88]. The uneven sputtering of calcium atoms out of the oven onto the substrates results in some positions having thicker layers of calcium than others. Relative thickness of the calcium film over the course of an experiment could be measured to change as there is a nearly exponential relationship [97] between the thickness of a thin metal film and transmitted incident light through an air-metal-substrate-air interface close to $T = e^{-\alpha d}$, where T is the percent of transmitted light, α is

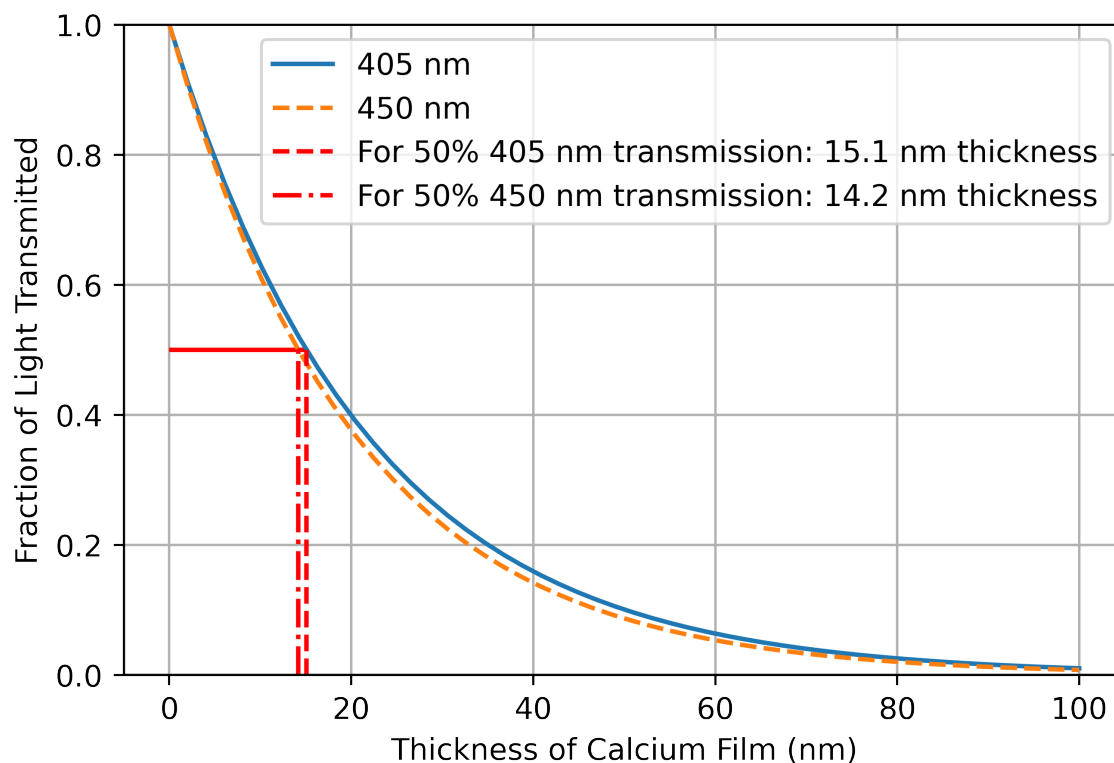


Figure 3.4. Amount of transmitted light passing through a thin film of calcium metal on a glass substrate. Amount of light transmitted is dependent on the thickness of the film and the wavelength of the light which determines the extinction coefficient, refractive index, and reflective index of the film. At 100 nm thickness, all wavelengths of light are approximately fully reflected or absorbed by the metal.

the absorption co-efficient for the metal at a specific wavelength, and d is the thickness.

Fig. 3.4 plots the approximate fraction of transmitted light through an air-metal-substrate-air interface for both 405 and 450 nm light and fused silica as the substrate.

3.3 Results

3.3.1 Fused Silica and Borosilicate Glass. Testing on fused silica and borosilicate glass gave nearly identical results for signals that were detected for a time on the order of < 10 s. Testing was also performed on borosilicate glass slides that had been sanded to a rough surface as defects in the surface were predicted in result in more

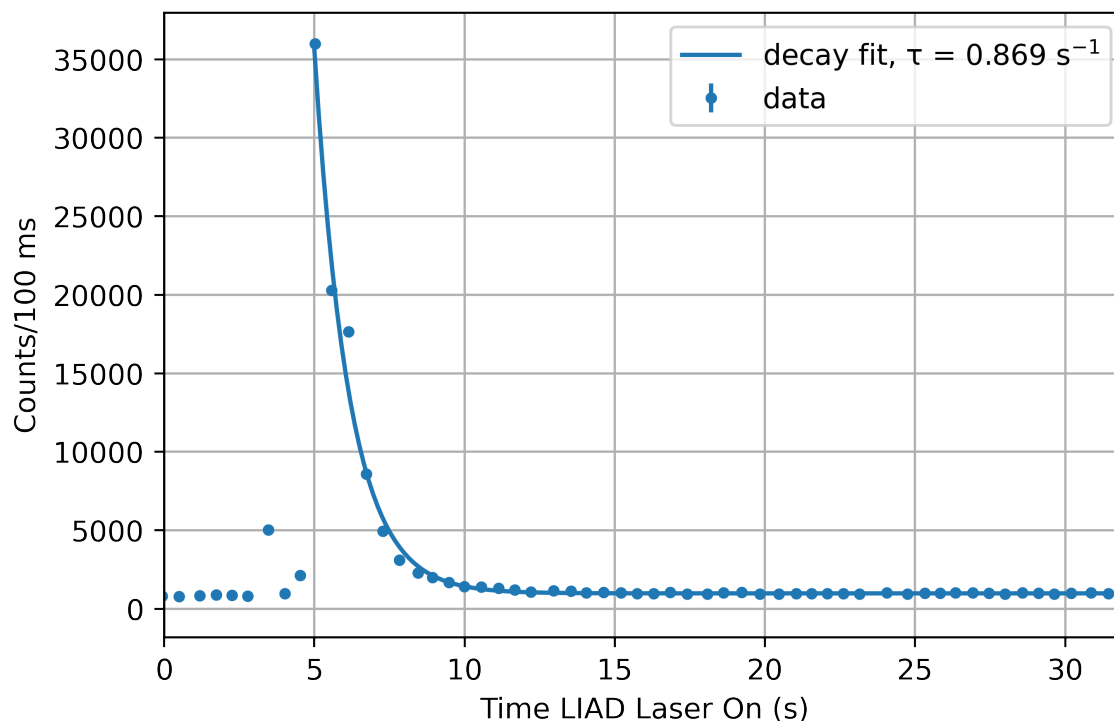


Figure 3.5. Fluorescence response from 450 nm laser on fused silica glass. Signal is fitted to a exponential decay function. The 450 nm laser was on for 5 s before the spike in signal, which in general was the amount of time need to see a signal of this type.

consistent LIAD signals. Testing showed, however, no apparent difference from the smooth surface.

Fused silica in combination with the 450 nm laser resulted in two detectable types of signal. First is the short-decay signal shown in Fig. 3.5, which could also be replicated with borosilicate glass. The 450 nm beam was directed at a position on the glass where a signal was detected after 5 s of exposure using the 423 nm laser at the expected resonant frequency for the cycling transition of ^{40}Ca . The peak of the signal and its decay are fitted to an exponential function with characteristic decay value $\tau = 0.869 \text{ s}^{-1}$. Use of the 405 nm laser resulted in similar signals with the same exposure time.

The second type of signal detected with fused silica and the 450 nm LIAD laser is the long duration signal shown in Fig. 3.6a. This particular signal was observed

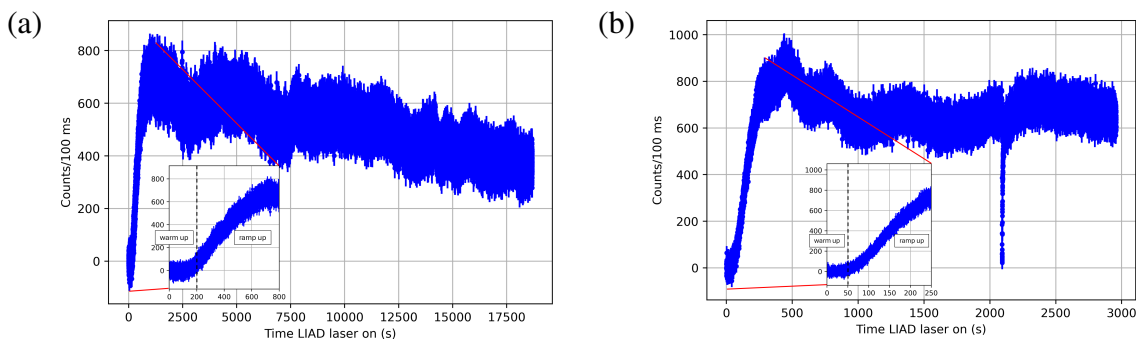


Figure 3.6. (a) Sustained signal of desorbed calcium off fused silica glass by a 450 nm laser beam that lasted nearly five hours. Inset of first 800 s with ‘warm-up’ and ‘ramp up’ times labeled. A signal like this was repeated several times over two weeks of testing at the same position. The fluorescence 423 nm laser was tuned to maximize the signal. (b) Sustained signal of desorbed calcium off of fused silica glass by an 808 nm laser beam that lasted nearly an hour. The fluorescence 423 nm laser was tuned to maximize the signal. Inset of first 250 s with ‘warm-up’ and ‘ramp up’ times labeled. The dip in signal after 2000 s was due to turning off the LIAD laser to see how quickly the signal would recover. The signal dropped to background in 1 s and recovered within 10 s of turning the LIAD laser back on.

continuously for nearly five hours with the same spot producing similar signals on repeated measurements over the course of two weeks of testing. This type of signal has been seen in other studies [86] where the decrease in signal can be attributed to the depletion of material being desorbed at the spot location.

A similar sustained signal was produced by shining the 808 nm LIAD laser on fused silica glass at different positions. An example is shown in Fig. 3.6b. To check the short time response of the signal with the LIAD laser, the beam was removed after 2,000 s and the signal went to background in 1 s. After the signal dropped below background, the LIAD beam was returned to the same position and the signal fully recovered after 10 s.

While testing with borosilicate glass gave similar short time-scale results with fused silica such as Fig. 3.5, longer lasting signals like Fig. 3.6 could not be repeated.

3.3.2 Graphene Sheets. Desorbing calcium from graphene sheets deposited by an oven gave short bursts of signals using both the 405 and 450 nm LIAD lasers. An example of signals found by slowly moving the 450 nm spot across a graphene sheet until

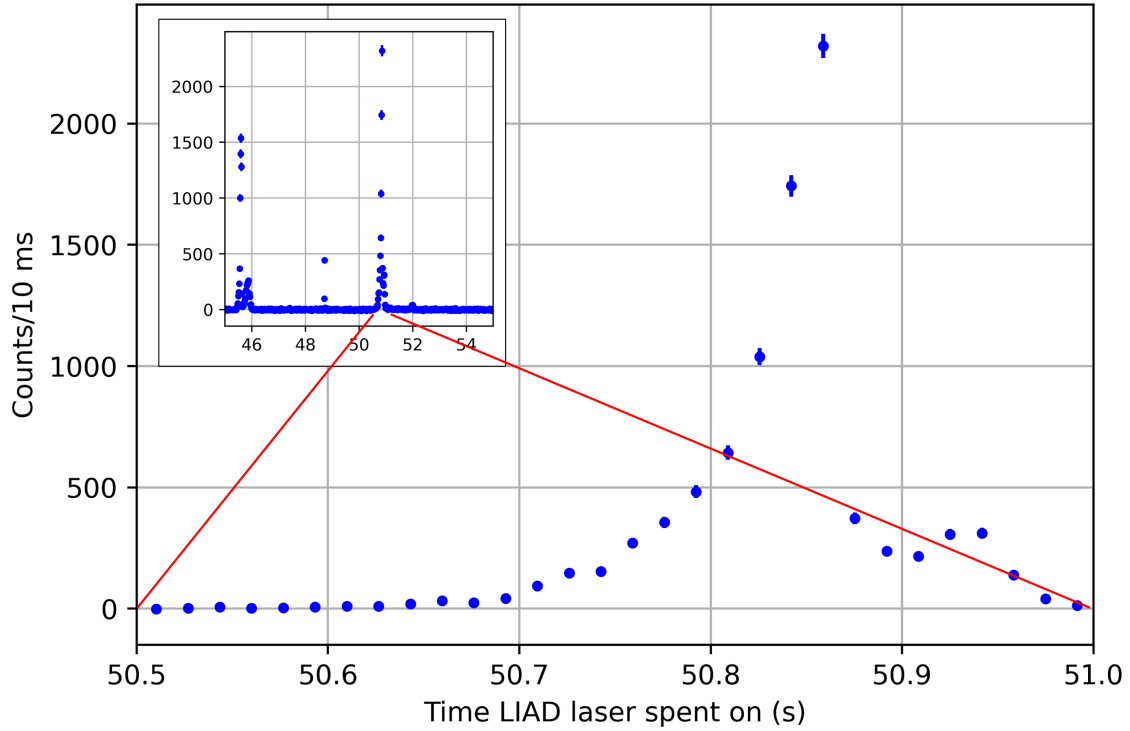


Figure 3.7. Fluorescence response from moving the 450 nm laser beam spot across a graphene sheet. Inset is of longer time showing other, smaller signals at nearby spots on the substrate surface.

detecting signals of approximately 0.5 s at a time is shown in Fig. 3.7. These signals were difficult to repeat and found at random intervals as the beam spot was spatially scanned across the graphene sheet. Similar signals were detected using the 405 nm laser. Signals longer than 1 s, like those discussed in Section 3.3.1, were never detected and, therefore, no attempts at measuring temperature from desorbed signal off of graphene via the method described in Section 3.2.1 were made.

3.4 Analysis

3.4.1 Short Duration Signal. Measuring the thickness of the calcium film at short duration signal sites lasting ~ 0.5 s, such as Fig. 3.7, revealed that the transmitted power of the LIAD beams did not increase significantly, indicating that this type of signal is due to defect sites [88]. These defects sites are most likely explained by defects in the

substrate structure, by dust that was on the surface of the substrate before calcium was deposited from the oven, or non-calcium material sputtered from the oven during the depositing process. Previous studies of LIAD show near-immediate response times on the order of less than 1 s, indicating that the burst signals measured here are LIAD events.

The longer-time signals on the scale of a few seconds, such as Fig. 3.5, corresponded to an increase in transmitted LIAD laser power as the fluorescence signal decayed. This type of signal is likely due to the depletion of available desorbing material and can be found in other LIAD studies [86]. The signal in Fig. 3.5 has an initial LIAD beam exposure time of 5 s before the signal appears. Due to the expected immediate response time like that seen in burst signals for LIAD, this signal could be due to a thermal effect where the energy from the incident beam must overcome the dissipation of heat through the film and substrate before the calcium is hot enough to increase the vapor pressure.

3.4.2 Long Duration Signal. Signals lasting several hours, such as those in Fig. 3.6, produce enough calcium atoms in the interaction zone to fluoresce at levels similar to load ion traps. Based on the thickness of the calcium spot determined from the transmitted LIAD beam, the size of the LIAD beam, and the rate at which atoms were being detected in the interaction zone, these substrate locations could each load an ion trap similar to the Paul trap in our lab (Chapter II) tens of thousands of times; far more than most traps need over a lifetime of use. This was determined by taking the fluorescence rate at which the ion trap can load an ion and appropriately scaling the dimensions of the oven and imaging areas to the LIAD chamber setup (Sections 3.1.1 and 3.1.2) to get an equivalent fluorescence rate.

These long duration events always had a warm up time before the signal began ramping up. The 450 nm desorption had a warm up of 200 s before signal began with a

ramp up to peak signal time of 600 s. The 808 nm signal appeared at shorter time scale with only 50 s warm up and 200 s ramp up. Because previous work did not see LIAD effects at wavelengths longer than green, the 808 nm signal is likely a thermal effect. It is likely the 450 nm signal was also a thermal effect as shorter wavelengths are more likely to produce a LIAD effect on the scale of $\ll 1$ s [86] but this signal had a warm up time of 200 s, which implies energy was being dissipated from the substrate to the fork and chamber.

3.5 Conclusions and Future Work

Based on the results presented here, cheap laser systems can produce atoms at rates consistent with oven and ablation loading systems in trapped ion experiments. We produced desorbed calcium atoms off of fused silica glass for up to hours at a time in the same location several days in a row using 450 and 808 nm lasers. Definitive direct measurements of the temperature of desorbed ions were not achieved. However, long duration signal behavior requiring several 10s of seconds to establish constant desorption indicates that a thermal effect is likely taking place. Fortunately, there are other methods and techniques, such as using a two-beam geometry for measuring the shift in the fluorescence spectrum due to the change in laser beam angle relative to the atomic beam [98], that may be able to more accurately measure temperature in future work. Regardless of temperature, the long duration signal producing combinations of fused silica glass and 450 and 808 nm light do have potential utility as consistent, long-lasting atom sources. The simplicity of the oven and desorption setup could be used as a less disruptive and cheaper alternative to popular oven and ablation methods.

The short burst signals produced by the combination of graphene sheets and 450 nm light did not exhibit any warm up behavior indicating they are a result of LIAD. Because of the short time scales of these signals, a temperature measurement was not

possible. Further study on the use of graphene sheets as a substrate for LIAD using calcium should be pursued.

One outstanding issue is the inconsistent performance of the substrate and metal. Once the calcium in a spot on the substrate had been depleted, a new location was needed to continue experiments. Finding a new spot in practice required slowly moving the targeted spot until a new signal was detected. If all calcium was depleted or a new spot could not be found, the oven would need to layer more material on the substrate, albeit not often if there are long lasting spots. In experimental platforms, this could cause disruptions on par with how typical thermal sources are used. Some experimental designs have been tested and proposed [74], but these included a magnetic actuator to block the atomic beam that would disrupt experiments sensitive to stray fields. Useful LIAD will require further development for robust performance in platforms like ion traps. Additionally, measurements of short burst signal temperatures were not possible here. Maintaining these signals for longer periods or being able to accurately determine and repeat conditions for these signals would allow for study of the temperature.

Future studies of LIAD and thermal desorption could include UV laser or LED light sources, more sophisticated platforms for accurately determining desorbed atom temperatures, and other UHV-compatible substrates. Additionally, surface material scientists [99] could study the physical mechanism behind why desorption occurs to make it clearer what combinations of wavelengths, substrates, and atomic species combinations should be most effective. These improvements along with the continued study of alkaline earth metal-based platforms will enable more useful and scalable applications in quantum information, computing, and sensing.

CHAPTER IV

LASERS

The setup of the of the dipole laser systems was performed by Drs. Daniel Moore [31], Jeremy Metzner [32], and Alex Quinn [33]. The experimental data for the characterization of the wavemeter lock with the Allan deviation was collected by myself and Dr. Quinn. The 976 nm fiberized injection locking setup was installed by myself with assistance from Dr. Quinn, the description of which is published in A. Quinn, *et al.*, *High-fidelity entanglement of metastable trapped-ion qubits with integrated erasure conversion*, [arXiv preprint arXiv:2411.12727 \(2024\)](https://arxiv.org/abs/2411.12727). The construction of the high-finesse cavity system for the 729 nm laser was performed by Marcus Polk and Prof. Allcock and the alignment by Marcus Polk with my assistance. Characterization of the cavity lock for the 729 nm laser was performed by myself and Dr. Jameson O'Reilly.

When trapped in a pseudopotential in a vacuum housing, an ion is isolated from other matter. To interact with the ions and elicit changes in either the motional or electronic states we use oscillating electromagnetic fields. We oscillate electromagnetic fields two ways: radio frequency (RF) magnetic fields (Section 2.5.1) and focused laser beams.¹ This chapter describes the lasers we employ and how we track and maintain appropriate frequencies.

4.1 Dipole Lasers

There are six dipole laser wavelengths needed to perform the basic processes of ionizing, cooling, fluorescing, and m -qubit preparation: 375, 393, 397, 423, 854, and 866 nm. For details on these laser systems themselves, see [31–33]. The m -qubit can also be prepared using the 729 nm quadrupole laser discussed in Section 4.4. The 375 and 423 nm lasers' corresponding transitions in neutral ^{40}Ca are diagrammed in Fig. 4.1a. The

¹“Before, my energy was all over the place. Now, it’s concentrated like a laser beam.” — Bart Simpson, *The Simpsons* S11E2, “Brother’s Little Helper”

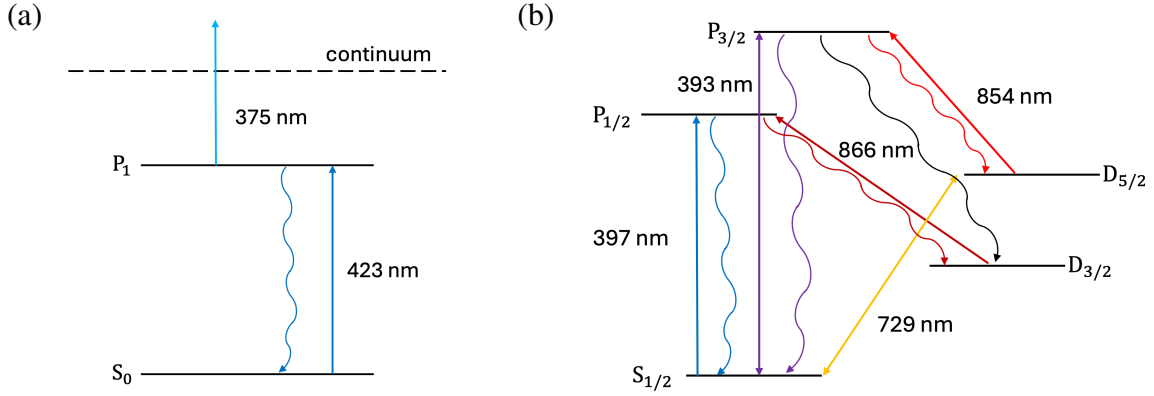


Figure 4.1. (a) Diagram of the ionization process for neutral calcium with relevant electronic states and transitions. (b) Electronic states of $^{40}\text{Ca}^+$ and relevant transitions for fluorescing, cooling, and prepping. Straight arrows represent transitions driven by laser interactions and squiggly arrows represent spontaneous emission. The black squiggly arrow from $P_{3/2}$ to $D_{3/2}$ does not have a corresponding laser to drive population in the opposite direction.

393, 397, 729, 854, and 866 nm transitions in $^{40}\text{Ca}^+$ are diagrammed in Fig. 4.1b. For conciseness, spectroscopic notation will be abbreviated such that the principal quantum numbers and multiplicity of the valance electron's energy levels will be omitted:

$$4^2S_{1/2} \equiv S_{1/2}, 4^2P_{1/2} \equiv P_{1/2}, 4^2P_{3/2} \equiv P_{3/2}, 3^2D_{3/2} \equiv D_{3/2}, \text{ and } 3^2D_{5/2} \equiv D_{5/2}.$$

Additionally, the magnetic quantum number $m \equiv m_j$ will always refer to the total angular momentum of the system.

4.1.1 Ionizing Atoms. For neutral ^{40}Ca , the 423 nm laser brings population to the excited P_1 state where it decays with a lifetime of 4.42 ns [100]. 375 nm light can be co-propagated with the 423 nm beam path and results in the excited electron being sent into the 'continuum' (Fig. 4.1a) where it dissociates from the calcium atom, resulting in a positively charged, trappable ion.

4.1.2 Fluorescence Detection. The fluorescing (*i.e.*, cycling) transition of ionized calcium $S_{1/2} \leftrightarrow P_{1/2}$ is at 397 nm. The excited state $P_{1/2}$ has a lifetime of 6.904(26) ns [101] and can be driven coherently with a sufficiently intense (typically pulsed) laser, which we do not have. Instead, our continuous-wave laser excites

population from the ground state and is followed by a spontaneously emitted photon as the state decays. Looking for scattered photons as the short-lived $P_{1/2}$ state decays is how we determine when our population is in the ground state. This is called fluorescence detection and it can be used to determine population in our qubit by only driving the $|1\rangle$ state to $S_{1/2}$, for example. See Section 4.1.4.

Fluorescence detection is not done with 397 nm light alone. The branching ratio for the decay to the metastable state $D_{3/2}$ (lifetime of 1.196(11) s) is 1:14 [101, 102] meaning population spontaneously decays to this state instead of returning to the ground state after an average of 14 cycles. To prevent population from building up in $D_{3/2}$ and resulting in an undetected ion, we apply the 866 nm beam, which drives population, or repumps, back to $P_{1/2}$, simultaneously with the 397 nm beam.

4.1.3 Cooling. To cool the ions, we utilize Doppler and electromagnetically induced transparency (EIT) cooling [103, 104]. Doppler and EIT cooling both use the 397 and 866 nm lasers. The 866 nm laser is used to repump population out of $D_{3/2}$ in both schemes.

4.1.3.1 Doppler Cooling. Doppler cooling relies on the Doppler effect to reduce the energy in the motional modes of trapped ions [105]. When an ion moves perpendicular to a laser beam path, the shift in the laser frequency is zero in the ion's frame of reference. When moving parallel (anti-parallel) to the laser beam, the laser frequency decreases (increases) in the ion's frame of reference. In the lab frame, this means the ion's motion contributes to the energy supplied by the laser photons by increasing the energy needed to absorb a photon when moving with the beam and decreasing the energy needed when moving against the beam. By decreasing the frequency of the laser in the lab frame to compensate for the motion of the ion, absorption of light on the appropriate cycling transition will be biased to when the ion is moving the

opposite direction of the laser beam path. For our setup, we only need one laser beam for all the motional modes of our ions because of the angle the beam is directed through our trap allows for a non-zero projection of every motional mode on to the direction of the beam propagation. Every emission event will ‘kick’ the ion momentum an amount $\hbar\vec{k}$, where $\vec{k}$ is the wavevector, that will average to zero over many events, but on average add energy $(\hbar\vec{k})^2/2m$. Because of the motionally dependent absorption and with the frequency of the laser detuned below resonance, the average absorption event will result in a decrease in the momentum of the ion. A steady state energy will be reached when the rate the photon kicks add energy matches the rate the absorption process subtracts energy.

A Doppler-cooled ion’s motional mode temperature follows the equation

$$k_B T = \frac{\hbar\Gamma}{4} \frac{1 + (2\Delta/\Gamma)^2}{2\Delta/\Gamma}, \quad (4.1)$$

where k_B is Boltzmann’s constant, T is the mode temperature, Γ is the linewidth of the cycling transition used for cooling, and Δ is the red detuning from the resonant frequency of the cycling transition. Eq. (4.1) is minimized for $\Delta = \Gamma/2$, making $T_{min} = \hbar\Gamma/2k_B$. This minimum temperature is limited by the cycling transition linewidth. The mode temperature is related to the average phonon number $\bar{n}$ such that [106]

$$\bar{n} = \frac{1}{e^{\hbar\omega_m/k_B T} - 1} = \frac{1}{e^{2\omega_m/\Gamma} - 1}. \quad (4.2)$$

4.1.3.2 EIT Cooling. Doppler cooling is limited to above ground state cooling because it utilizes the $S_{1/2} \leftrightarrow P_{1/2}$ transition which is too broad to resolve the motional sidebands where a single phonon can be added or subtracted. EIT cooling suppresses the carrier transition and phonon-adding motional sideband [103, 104] through destructive interference and allows the phonon-subtracting motional sideband to dominate through a Fano-like absorption profile [107]. Access to the phonon-subtracting motional sideband

is critical to common ground state cooling methods and allows for rapid cooling over a broad range of motional modes at once.

EIT cooling uses a three-level Λ configuration of electronic states where two ground states $|g\rangle$ and $|f\rangle$ are coupled via a short-lived state $|e\rangle$ [103, 104]. A Fano-like absorption profile [107] is created by driving $|g\rangle \leftrightarrow |e\rangle$ with Rabi frequency Ω_σ at detuning Δ and $|f\rangle \leftrightarrow |e\rangle$ with Rabi frequency $\Omega_\pi \ll \Omega_\sigma$ at detuning Δ_π . The driving field light shifts the dressed states by an amount δ :

$$\delta = \frac{1}{2} \left(\sqrt{\Omega_\sigma^2 + \Delta} - |\Delta| \right). \quad (4.3)$$

For $\Delta_\pi = \Delta$ when an ion is in a harmonic trap of frequency ω , the vibrational sidebands will absorb light on the red sideband transition (phonon subtracting) with maximal probability when $\delta = \omega$. Every absorption event by the red sideband results in the loss of one phonon. With carrier absorption suppressed, the blue sideband absorption can be reduced relative to the red sideband resulting in a net loss of phonons.

4.1.4 m -Qubit Preparation. We designate a qubit in the $D_{5/2}$ manifold with Zeeman levels $m = +5/2$ as $|0\rangle$ and $m = +3/2$ as $|1\rangle$. We call this the m -qubit because it is in the metastable manifold [27]. Likewise, we designate our ground (g) qubit in $S_{1/2}$ with Zeeman levels $m = +1/2$ as $|\uparrow\rangle$ and $m = -1/2$ as $|\downarrow\rangle$. See Fig. 4.2. To prepare $|0\rangle$, we use a set of three laser pulses with specific polarizations. The polarization of laser light can be described as a linear combination of three types of polarization with respect to the magnetic field $\vec{B}$ in our trap: π light drives transitions with $\Delta m = 0$, σ_+ light drives transitions with $\Delta m = +1$, and σ_- light drives transitions with $\Delta m = -1$. The three sets of laser pulses are as follows:

1. 397 nm σ_+ and 866 nm beams pump population to $S_{1/2}, m = +1/2$,
2. 393 nm σ_+ shelves population in $D_{5/2}, m \geq +1/2$, and

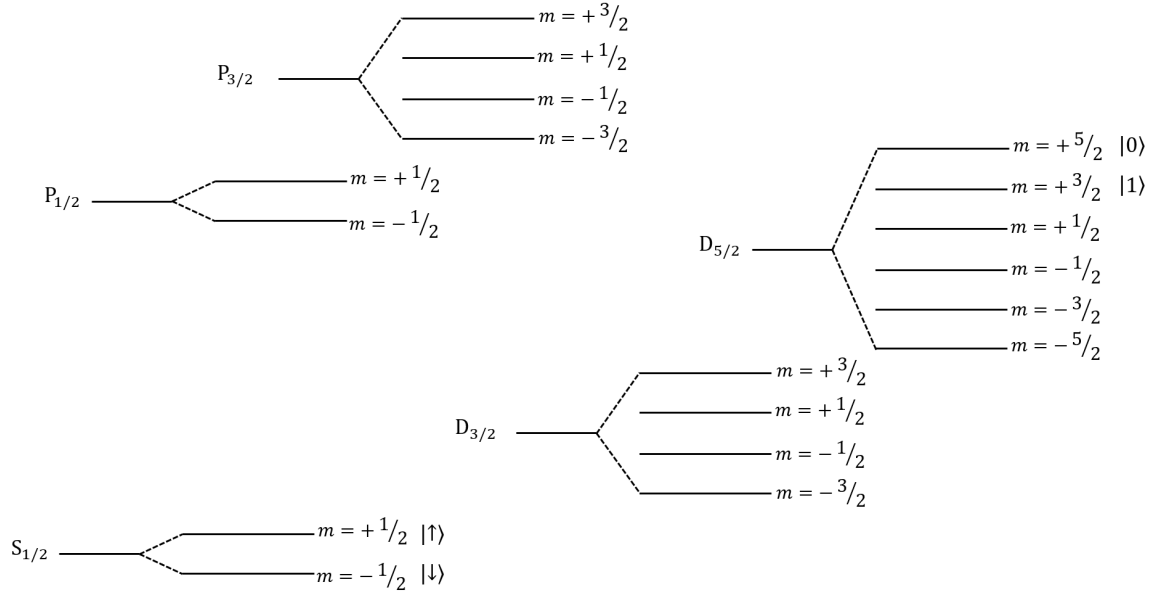


Figure 4.2. Zeeman sublevels in the $S_{1/2}$, $P_{1/2}$, $P_{3/2}$, $D_{3/2}$, and $P_{5/2}$ manifolds. Designation of the m -qubit is in $D_{5/2}$ with $|0\rangle$ and $|1\rangle$ and the g -qubit is in $S_{1/2}$ with $|\uparrow\rangle$ and $|\downarrow\rangle$.

3. 854 nm π deshelves population out of $D_{5/2}, m < +5/2$ back to the ground state.

Given the branching ratios for $P_{3/2} \leftrightarrow D_{5/2}$, one iteration of this pulse sequence prepares the m -qubit in $|0\rangle$ at a rate of 67%. By repeating this preparation sequence (typically $N=7$), we can achieve over 99% successful qubit preparation. For more details on the m -qubit including preparation fidelity and state detection, see A. Quinn's dissertation [33].

4.2 Wavemeter Locking Lasers

To make sure these basic functions perform well, the laser frequencies must be monitored for drifts and corrected as necessary. The device we use to measure the frequency of lasers with a maximum resolution of 200 kHz is called a wavemeter (WS8-10 from HighFinesse). A small percentage of light from each laser system is picked off with beamsplitters and coupled into multichannel fiber switches which directs one laser beam at time into the wavemeter [108, 109]. The wavemeter then returns the measured frequency of the light. A wavemeter needs to be calibrated against some reference in

order to accurately measure frequency. Common references include highly stable laser systems [110–112], atomic transitions [113], and optical cavities [114]. Our wavemeter setup relies on our trapped ions as an absolute frequency reference using each laser’s corresponding atomic transition in $^{40}\text{Ca}^+$ and neutral calcium.

The main source of drift comes from environmental changes including temperature (as well as humidity and pressure). The effects of these factors can be minimized by isolating the wavemeter device from the environment [108, 112, 114]. We achieve this by enclosing the wavemeter in a box surrounded by insulating foam and reflective material and stabilizing the temperature of the box with thermoelectric Peltier modules and heatsinks (see Fig. 4.3). Additionally, the laser systems and wavemeter are placed in a temperature controlled room only accessed for maintenance and not regular experimental work. The room temperature is stabilized to typically better than 0.4°C over a day and changes in the wavemeter box are $\sim 10\%$ the size of room temperature changes.

Based on the wavemeter measurement, the control computer provides the proportional-integral-derivative (PID) control on the laser diodes’ cavity piezo voltages and current to adjust the frequency to match a lock frequency. The PID control adjusts once every 100–1000 ms and can keep the measured laser within ~ 1 MHz of the lock frequency. A study was performed over the course of a week to show the stability of the center frequency of our cooling lasers using these methods. Spectroscopy experiments were performed with four of our lasers addressing the relevant transitions in $^{40}\text{Ca}^+$ (Fig. 4.1b) and the data was fitted to a power-broadened Lorentzian curve every work day for eight days. See Fig. 4.4 as an example of the fitting for the 397 nm laser spectroscopy. The center frequency of each scan is plotted in Fig. 4.5 along with the environmental conditions of the laser laboratory room. Changes in center frequencies range from

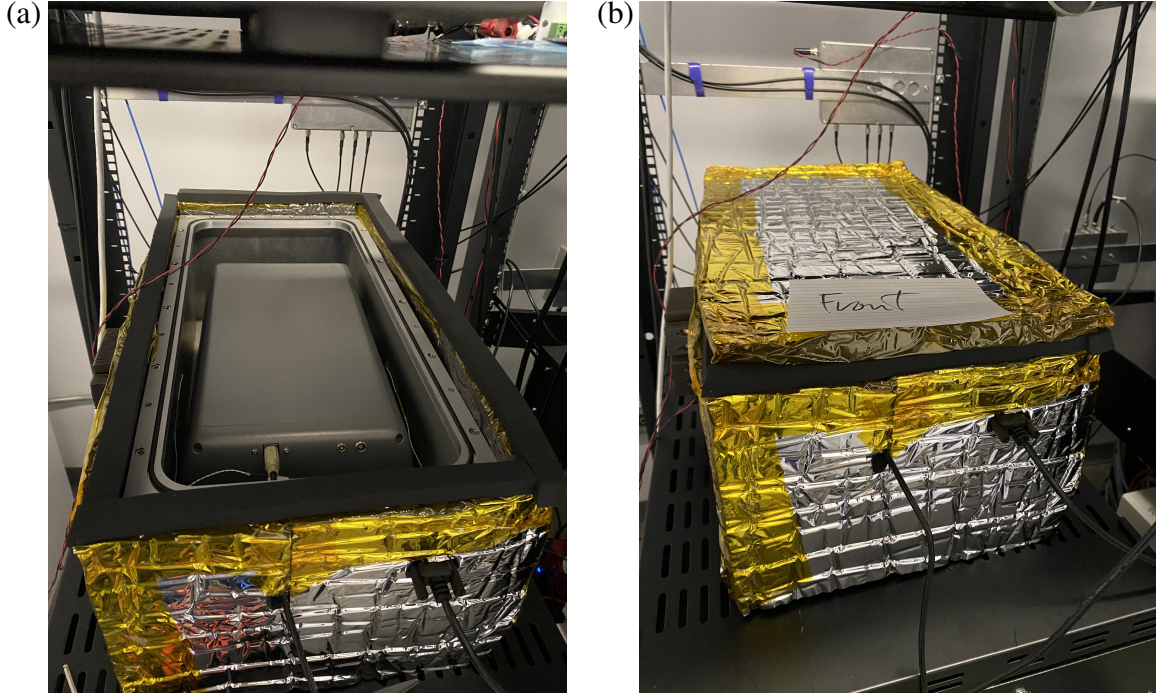


Figure 4.3. (a) The wavemeter visible inside the temperature control box. (b) The wavemeter fully enclosed in the temperature control box. Two thermoelectric Peltier modules are placed on either side of the box next to temperature sensors to provide PID control on the box temperature. A third sensor is placed near the ‘front’ of the box as an out-of-loop monitor. Feedthrough holes are cut through the box and insulating material for fiber optics, power, sensors, and communication cables.

0–5 MHz. These measurements were done on a daily basis and not on the scale at which typical experiments run on the order of a few seconds or minutes. Corrections to frequency drifts at shorter timescales are discussed in the following section using Allan deviations.

4.2.1 Allan Deviation. Allan variance measures the fractional frequency fluctuations in oscillators and can distinguish between several types of noise over a wide range of time domains [115]. The Allan variance σ_y^2 is defined as

$$\sigma_y^2(\tau) = \frac{1}{2(M-1)} \sum_{i=1}^{M-1} [y_{i+1} - y_i]^2, \quad (4.4)$$

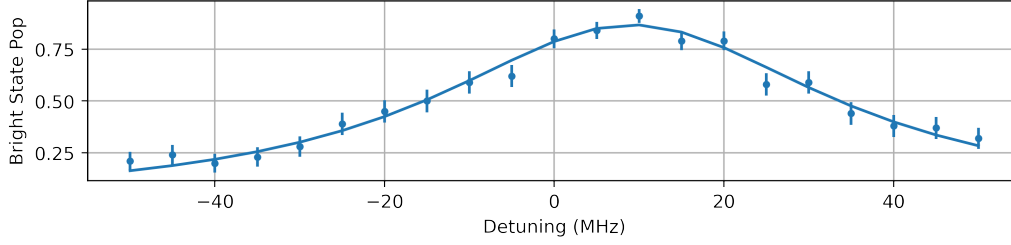


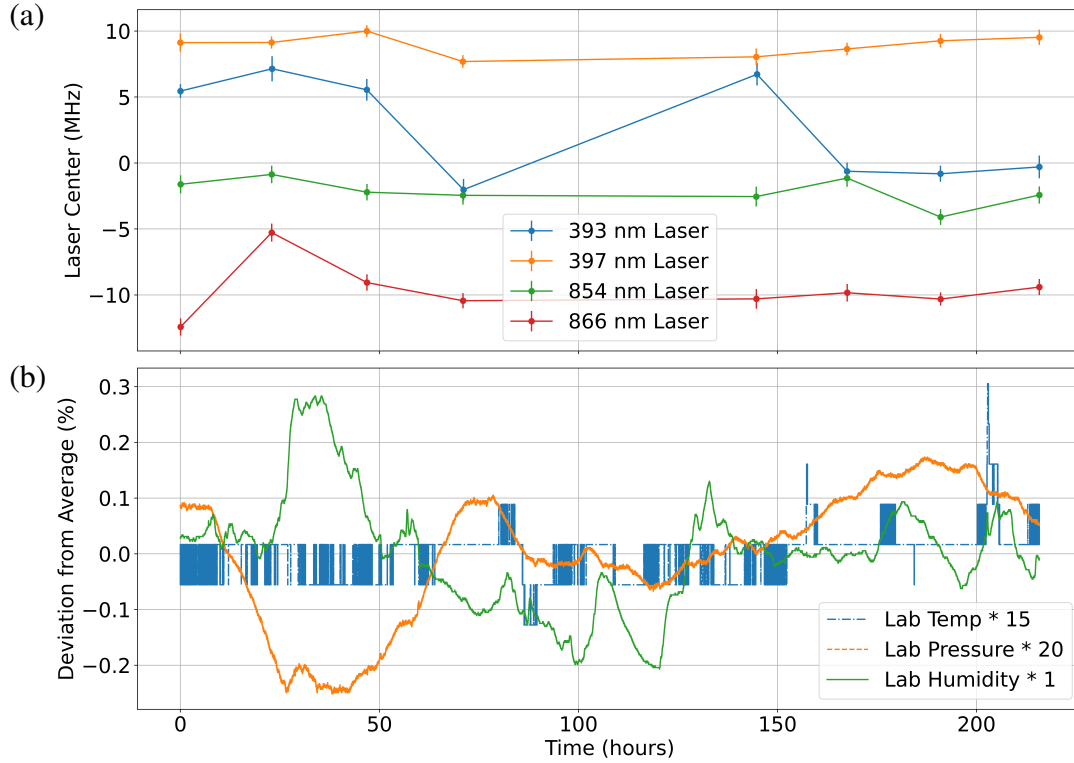
Figure 4.4. Example spectroscopy and fit of the 397 nm laser frequency detuning against the $^{40}\text{Ca}^+$ transition for $S_{1/2} \leftrightarrow P_{1/2}$.

where

$$y_i = \frac{x(iT + \tau) - x(iT)}{\tau} \quad (4.5)$$

is the n th fractional frequency average over the observation time τ , T is the time between each frequency sample, M is the number of clock reading intervals, t is the reference clock, and $x(t)$ is the average fractional frequency time series. Allan variance is often plotted as the Allan deviation, which is simply the square-root of the variance. For an oscillator with only white noise, *e.g.*, a perfect laser with no external noise and only fluctuations according to its linewidth, the Allan deviation follows the power law $\sigma_y \sim \tau^{-0.5}$. When τ approaches the time scale at which a different noise source occurs, σ_y will stop following a linear trajectory on a log-scaled plot and the slope will increase and/or oscillate.

To study noise on a range of time-scales, the 397 nm laser stability was measured over several minutes. The laser frequency was parked near the half-width at half-maximum from resonance of the $S_{1/2} \leftrightarrow P_{1/2}$ transition for $^{40}\text{Ca}^+$. The slope of the Lorentzian shape of the linewidth is maximal at this point meaning changes in the fluorescence of the ions will be maximally sensitive to changes in the laser frequency. Photons were collected and time-tagged in periods of 3 s with 0.5 s of dead time between measuring periods. In the analysis of the collected data, the 3 s measuring periods are further divided into smaller bins according to the time-scale denoted by τ .



*Figure 4.5. (a) Measurement of cooling lasers fitted center frequencies detuned from atomic resonance. (b) Deviations of the laser laboratory environment conditions from average over the same time period. The laboratory temperature and pressure deviations are scaled by factors of 15 and 20 respectively to be visible on same plot as the humidity which is left unscaled. The horizontal axis is shared between **a** and **b**.*

Dead time refers to the gaps in data collection where the data acquisition control (DAC) system must process information before continuing. For our ARTIQ system, we were limited to 3 s of continuous data collection before data processing was required. Dead time can be dealt with in two ways without limiting the experiment lengths to timescales the DAC can handle without stopping. First, the dead time can be accounted for in the calculation of the Allan variance through a bias function [115, 116]. This process is easily done in open source software [117, 118]. Second, the dead time can be filled in with simulated data generated by interpolating neighboring data points. Subtle differences in the results of these two methods can occur [116], so it is best to examine both methods.

Fig. 4.6 is a plot of the Allan deviation for the 397 nm laser with and without interpolated dead time compared to a simulated ‘perfect’ laser with only white noise. The data for $\tau < 0.1$ s follows the slope of the perfect laser, meaning white noise dominates. The lift-off around 0.1 s shows where the instability of the laser become noticeable. Our experiments occur on the scale of $\ll 0.1$ s with most pulses taking at most 1 ms and pulse sequences < 10 ms. In the previous section, it was determined that the wavemeter lock for the dipole lasers maintained the laser frequencies within 5 MHz of the target value. The linewidth of the dipole transitions are 10–20 MHz meaning a 5 MHz error is sufficient for dipole operations. Interestingly, the interpolated dead time suppresses fluctuations of the resulting Allan deviation while still following the shape of the biased dead time. Oscillations due to the dead time in the data without interpolation stop at 0.5 s, corresponding to the dead time interval. The interpolated data smooths out these oscillations.

Photons were collected while the lock was on as well as off to show the differences in the fractional frequency fluctuations. In Fig. 4.7, the overlap of the two experiments show that the wavemeter does not impact the stability of the laser at short time scales of $\ll 0.1$ s where the majority of our individual operations are < 1 ms and full experiments < 10 ms. The fractional frequency fluctuations are the same until ~ 0.4 s where the experiment with the lock engaged begins to diverge from the experiment with the lock off. At ~ 30 s, the fractional frequency of experiment with the lock on is dropping when it intersects with the experiment with the lock off as it increases. This is the time scale at which the wavemeter is preventing the frequency from drifting from the lock point, decreasing the fractional frequency. On the time-scales of 1–30 s, the lock on experiment has a larger fractional frequency likely due to sub-optimized PID parameters making the corrections too sensitive in this range.

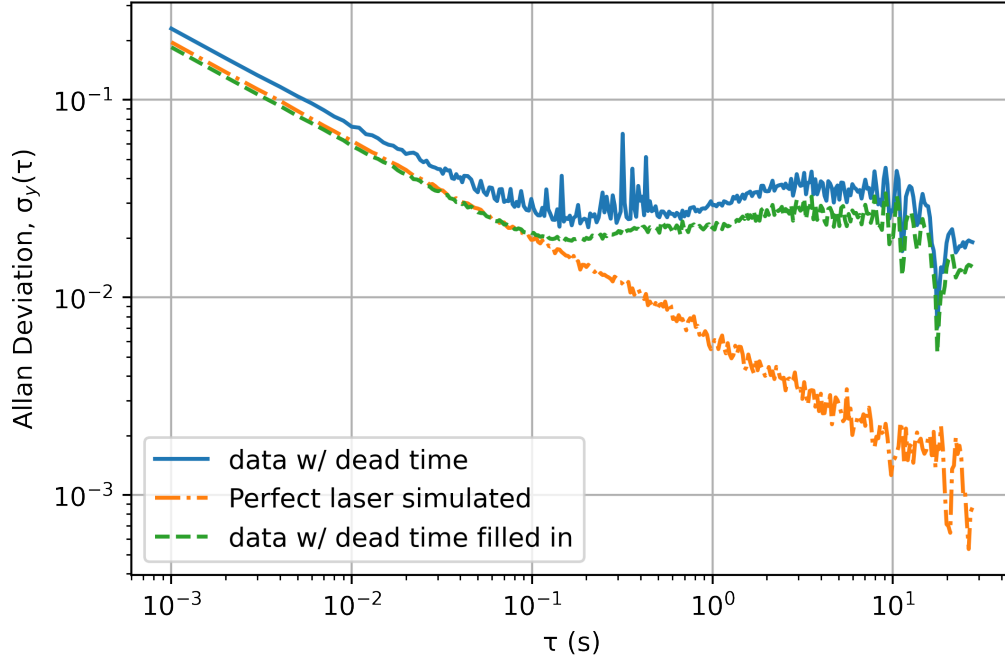


Figure 4.6. Allan deviation of the 397 nm laser with and without interpolated dead time compared against an ideal laser.

4.3 Far-detuned Raman Lasers

Once the m -qubit is prepared, we can perform operations to manipulate the state followed by a readout sequence (see Sections 4.1.2 and 5.1). Our designated qubit sublevels are separated by 2.63 MHz due to the linear Zeeman effect with a B-field of 1.57 G. This transition can be driven via oscillating currents applied to the one of the trapping rods, which produces an oscillating magnetic field as described in Section 2.5.1. The transitions between the sublevels of $D_{5/2}$ are degenerate, so to prevent population from being driven between all six sublevels, we apply light from a different 854 nm laser, the “854 light shift (LS)” laser, detuned -21 GHz from the $P_{3/2} \leftrightarrow D_{5/2}$ transition. This laser shifts all but the $m = +5/2$ and $+3/2$ levels. Instead of driving population between the qubit states with RF, we also employ a two-photon Raman process [119] using 976 nm lasers -44 THz detuned from the $D_{5/2}$ to $P_{3/2}$ [34] transition, as depicted in Fig. 4.8.

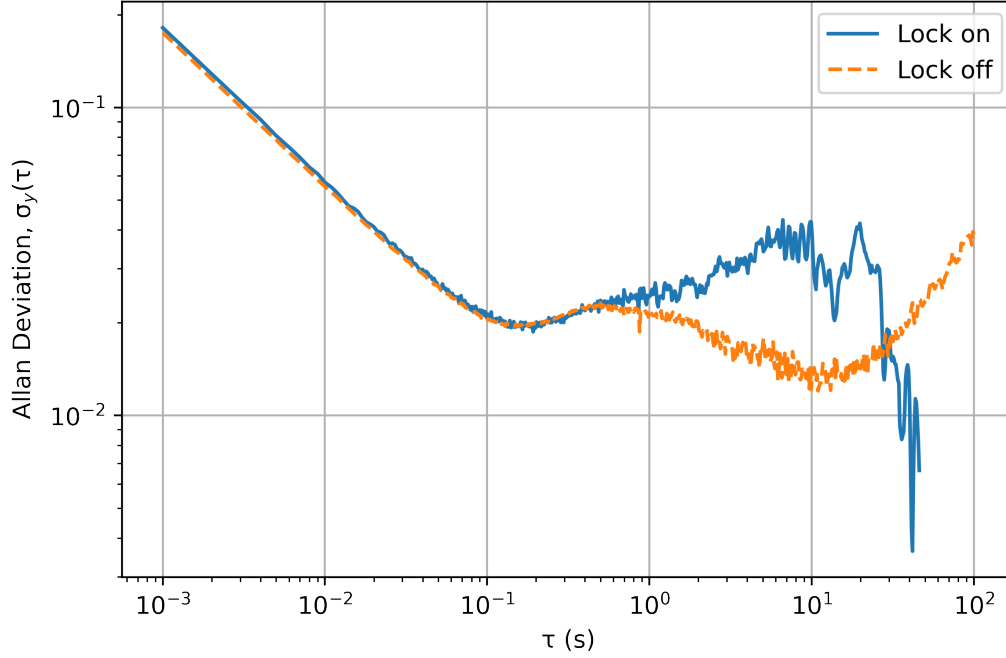


Figure 4.7. Allan deviation of the 397 nm laser with wavemeter lock on and off.

4.3.1 Fiberized Injection Locking. To maintain their relative frequency offset and increase the optical power available in each beam, we implemented an injection-locked diode system based on [120]. The layout of the fiber-injection system is diagrammed in Fig. 4.9. Light from a single volume-holographic-grating (VHG)-based free-space 976 nm 700 mW hybrid external cavity laser (HECL) (the ‘seed’ laser, 10976SB0700B from Innovative Photonic Solutions) passes through an optical isolator to prevent feedback. Then, it passes through a beamsplitter to pick off light for diagnostics before it is split into two beam paths which are each coupled into Port 1 of two fiberized optical circulators (PMOC-F-3-98-5-PM980-90-05-FA from DK Photonics). About 50 mW of light from the seed laser then exits each circulator from Port 2 and is injected into two fiber-pigtailed 976 nm 1 W diodes (the ‘amp’ lasers, 976LD-2-0-0 from AeroDIODE). The amp laser light re-enters the circulators from Port 2 and is directed out through Port 3 into fiberized acousto-optic modulators (AOMs) (model number CAFA-

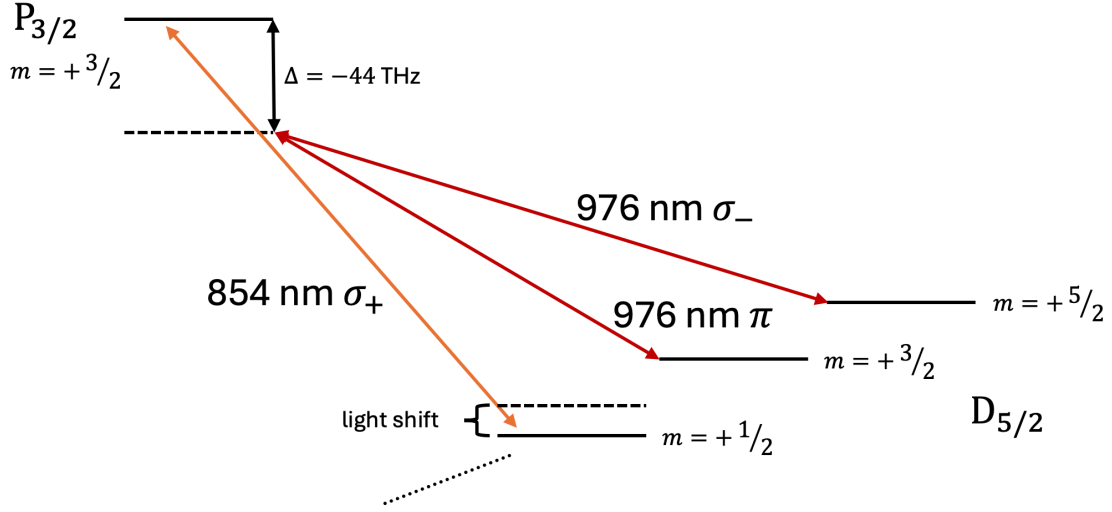


Figure 4.8. Raman transitions for driving between the m -qubit states using the far-detuned 976 nm lasers. The 854 nm σ_+ beam provides the light shift that isolates the $m = +5/2, +3/2$ qubit states from the other four sublevels in $D_{5/2}$.

100-2-TE2F-980-AF-A87 from CASTECH, Inc.) for switching and frequency tuning.

The circulators prevent amplified light from exiting Port 1 and returning to the seed laser.

At the vacuum chamber, each beam is collimated with a beam waist diameter of ~ 17 mm by 60FC-L-4-M40-54 Schafer+Kirchhoff collimators. Beam samplers pick off $\sim 2.5\%$ of the light in each beam before the chamber and direct it to photodiodes (PDs) for intensity stabilization. After the pickoffs, each beam is measured to have over 200 mW, which is focused by lenses with focal length 150 mm to a waist of $\sim 40 \mu\text{m}$ at the center of the ion trap.

The difference in the two Raman beams' frequency, the beat note, is the driving frequency of the Raman transition for the m -qubit. Control of the beat note is important to prevent off-resonant driving. To make a beat note measurement, a second pickoff of each beam directs light into a 50:50 fiber splitter that combines the light into a PD connected to a spectrum analyzer that measures the spectral power density (Fig. 4.9). The beat note of two tones separated by 1 MHz between the two AOMs is measured to have a full width at

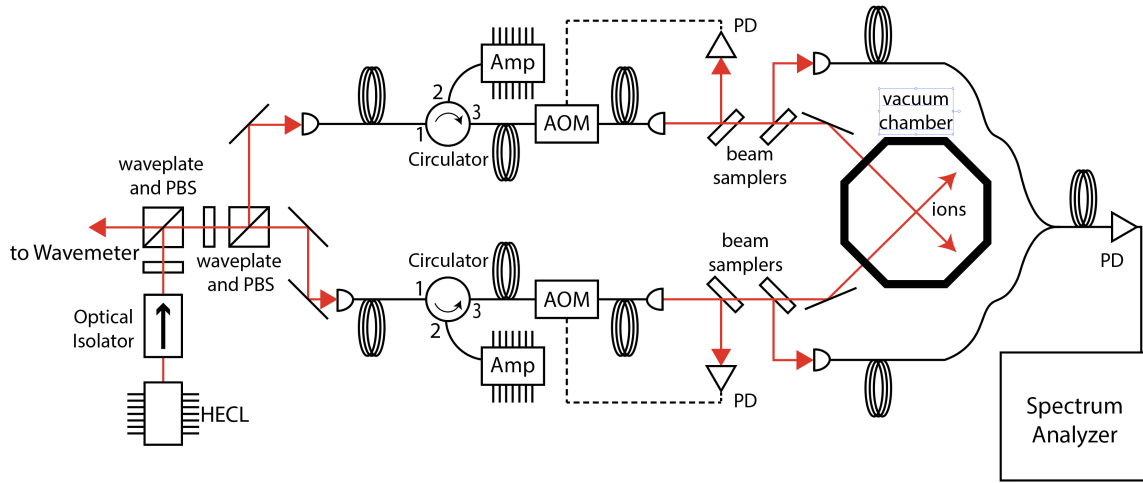


Figure 4.9. Schematic diagram of the fiberized injection lock setup of the 976 nm Raman lasers as well as the beat note measurement.

half maximum of 4 Hz (Fig. 4.10), indicating precise control of the driving frequency for the m -qubit.

Compared to our previous system that used a single 700 mW free space diode laser (details in [33]), this setup triples the optical power in each beam. In turn, this triples the rate at which the Raman beams entangle our ions [34]. Another major advantage is the cost-effectiveness of using small, butterfly-packaged high power diode lasers that were developed for the telecommunications industry as pump diodes for erbium-doped fiber amplifiers.

4.4 Narrow Linewidth Quadruple Transition

Access to the metastable manifold directly from the ground manifold in alkaline earth metals is useful for a variety of reasons, including those discussed in Chapter V. The electric quadrupole transition, *i.e.*, dipole forbidden transition, for $^{40}\text{Ca}^+$ between the $S_{1/2}$ and $D_{5/2}$ manifolds is at 729 nm (see Fig. 4.1b). ‘Forbidden’ [121] transitions are named so because they violate dipole selection rules, which do not apply to quadrupole transitions because the quadrupole interaction can impart more angular momentum

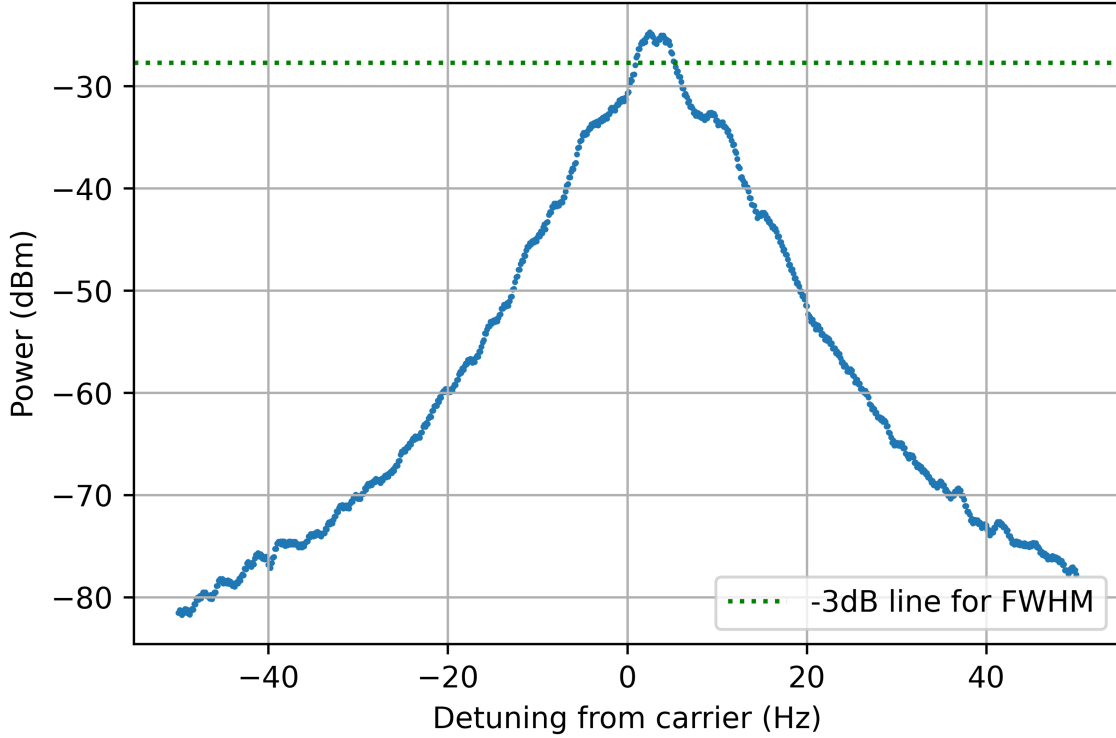


Figure 4.10. Frequency spectrum of the beat note between 976 nm Raman beams.

($\Delta m > 1$). Electric quadrupole transitions are relatively weak compared to permitted transitions that obey dipole selection rules,² *i.e.*, electric dipole transitions, because of much weaker matrix elements. In particular, $S_{1/2} \leftrightarrow D_{5/2}$ transitions violate the dipole selection rule $\Delta l = \pm 1$ and some violate the rule $\Delta m = 0, \pm 1$. This violation is possible because the dipole interaction is just the first order interaction and does not consider higher order interactions such as the electric quadrupole interaction.

In $^{40}\text{Ca}^+$, $S_{1/2} \leftrightarrow D_{5/2}$ has a natural linewidth $\Gamma = 0.136(1)$ Hz [102] related to its long lifetime τ of 1.168(9) s with $\tau = 1/(2\pi\Gamma)$. In order for a transition to be efficiently driven by a laser beam, there needs to be sufficient spectral power in the narrow frequency band matching the natural linewidth. We utilize a 729 nm 40 mW Toptica ECDL with

²Interestingly, the first observations of forbidden transitions (including several in calcium) were from stellar emission lines performed by astronomers and astrophysicists [122].

a specified free running linewidth of < 60 kHz controlled via a Toptica Digital Laser Controller (DLC) pro. The wavemeter described in Section 4.2 has a maximum resolution of 200 kHz and cannot lock the 729 nm laser to a narrow enough linewidth to coherently drive the optical transition. Additional equipment is necessary to lock the laser to a linewidth below 10 kHz where driving of this weakly allowed transition is possible.

4.4.1 PDH Lock. A tried and true method for producing a narrow-linewidth laser is by locking it to a reference cavity with the Pound-Drever-Hall (PDH) [123–127] technique (see Fig. 4.11). Laser light passes through an electro-optic modulator (EOM) (model number PM730 from JENOPTIK) to modulate its phase, resulting in the creation of sidebands, and is then directed into a high-finesse cavity. The electric field E_0 of the incident laser beam after phase modulation becomes

$$E_{inc} = E_0 e^{i(\omega t + \beta \sin \Omega t)} \quad (4.6)$$

and can be expanded using Bessel functions to

$$\begin{aligned} E_{inc} &\approx E_0 [J_0(\beta) + 2iJ_1(\beta) \sin \Omega t] e^{i\omega t} \\ &= E_0 \left[J_0(\beta) e^{i\omega t} + J_1(\beta) e^{i(\omega + \Omega)t} - J_1(\beta) e^{i(\omega - \Omega)t} \right] \end{aligned} \quad (4.7)$$

where J_n are the Bessel functions of the first kind, β is the modulation depth, ω is the carrier frequency, and Ω is the modulation frequency. Higher order terms of the expansion have been dropped. The modulation creates two first order sidebands at frequencies $\omega \pm \Omega$. The total power of the incident beam $P_0 \equiv |E_0|^2$ is divided between the carrier power P_c and sideband powers P_s maximally when $\beta < 1$ such that $P_0 \approx P_c + 2P_s = J_0^2(\beta) P_0 + J_1^2(\beta) P_0$ [123, 127].

When the frequency of the laser matches the resonance of the cavity, the light passes through the cavity and the reflected signal (measured by a PD) off the front of the cavity dips. When off resonance, the reflected signal will increase. To know which direction to adjust the frequency we must know the phase of the reflected light. There are

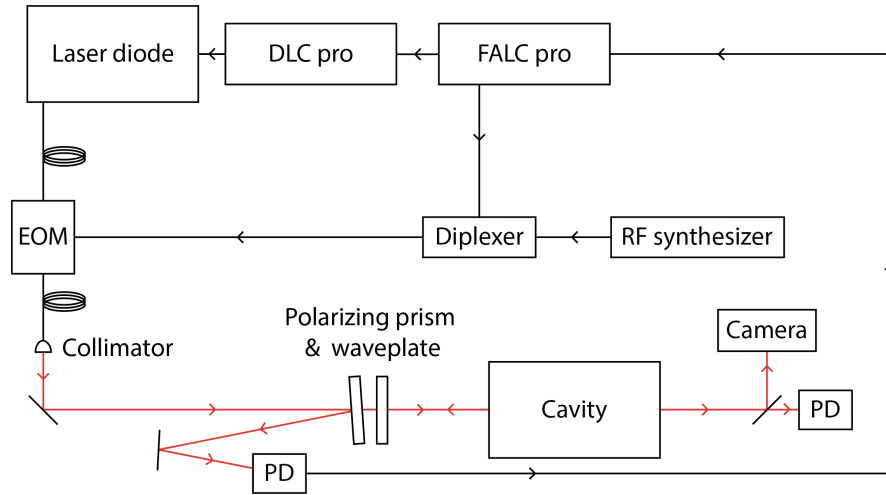


Figure 4.11. Schematic diagram of PDH locking setup. Black lines with loops represent fiber optics; black lines with directional arrows and without loops represent electrical cables; red lines with directional arrows represent laser beam paths.

no electronics that can directly measure the phase of light via the electric field, but we can indirectly determine it. The sidebands generated by the EOM have a specific phase relation with the carrier, so by interfering the reflected signal with the modulation signal we generate a beat pattern from which we can measure the phase. This beat pattern is then mixed with the EOM RF driving signal to create an error signal. The error signal feeds into a PID controller that adjusts the diode laser current to put the laser frequency back in resonance with the cavity.

Much of the electronics for our PDH lock is contained within a Fast Analog Linewidth Control (FALC pro) unit, also from Toptica. It provides the RF signal for the EOM to produce sidebands on the laser, mixes the reflected cavity voltage signal with this RF signal to produce an error signal, and provides the PID feedback voltage for the diode laser current modulation port.

Additionally, an RF synthesizer produces a signal that is multiplexed with the FALC pro RF signal via a diplexer. The synthesizer allows us to apply an additional

sideband to the laser carrier frequency (which will have sidebands from the FALC pro RF signal) at a variable frequency, which is required because the cavity resonance is not related to the ion resonance. By applying the synthesizer RF only to the laser beam incident on the cavity, the vast majority of the laser power, which is directed into the ion trap as a separate beam, can remain on the ion transition frequency. The beam with the applied RF tones can then be used to PDH lock the laser system to the cavity's resonant frequency far from the ion's resonant transition frequency, typically 10s–100s of MHz. The set frequency of the RF synthesizer is determined by finding the resonant frequency of the ion transition starting from the literature value [128] and an uncalibrated measurement from the wavemeter. Exposing the ion to the aligned 729 nm beam, the set frequency of the RF synthesizer can then be adjusted tens of MHz at a time until quantum jumps are maximized while fluorescing the cycling transition [129]. With light on the ion confirmed by observing quantum jumps, the AOM can be used to finely determine the frequency offset needed to calibrate the wavemeter measurement through spectroscopy experiments. The use of an AOM allows us to detune the frequency of the beam directed at the ion an additional ± 20 MHz as necessary. See Fig. 4.12 for a plot of the frequency spectrum of the incident beam on the cavity setup. A camera and PD on the transmission side of the cavity are used for alignment as well as to provide convenient checks for maintaining the lock.

The cavity we use comes from Stable Lasers Systems, LLC. It consists of two 1" diameter, 0.25" thick mirrors separated 100 mm. The mirrors are polished and have an inner coating with surface transmission value of $2\text{E}-4\%$ and outer coating with surface reflection value of 0.05% according to the product specifications, both measured at 729 nm. The cavity is encased in a vacuum chamber maintained by an ion pump and temperature controller to keep the cavity length stable. The chamber contains

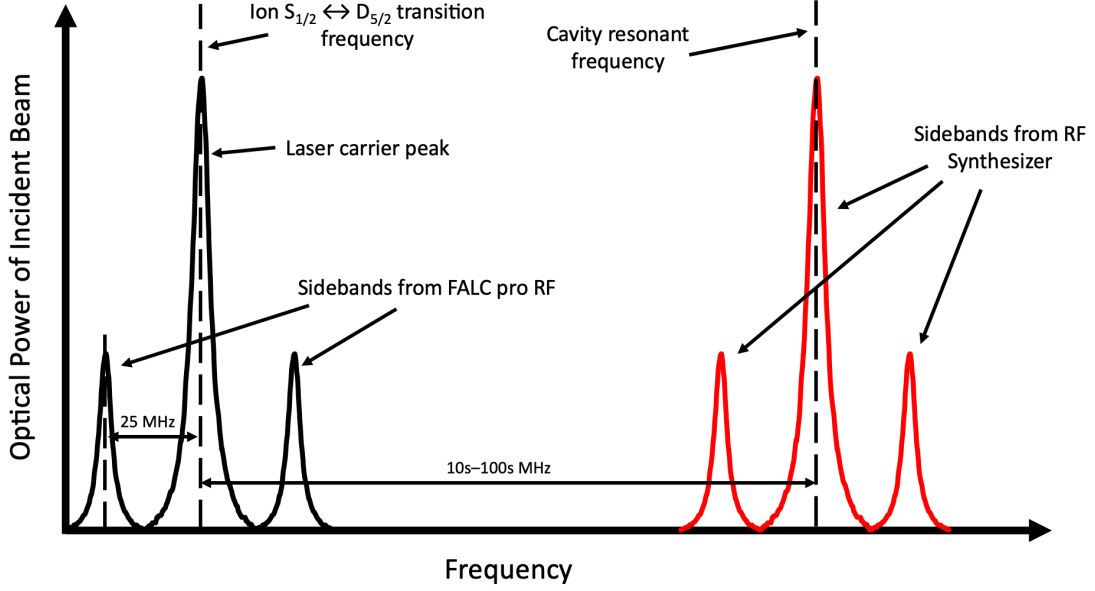


Figure 4.12. Diagram of the frequency spectrum of the incident laser beam to the cavity in the PDH setup. The FALC pro applies sidebands on either side of the laser carrier frequency (black curves). The RF synthesizer applies sidebands of the laser carrier and FALC pro RF sidebands with only the higher of the two sidebands depicted here (red curves).

multiple thermal shields with heaters and thermistors to that keep temperature changes to < 1 mK/day after stabilizing according to the product specifications. Vibrations are minimized by placing the cavity vacuum chamber on an active stabilization platform (model number TS-150 from Herzan, LCC). Additionally, the table sits on a 3" thick, 150 lb. slab of granite with four 2.5" diameter Sorbothane™ hemisphere rubber feet to prevent slipping on the room floor. According to the product specifications, the cavity free spectral range (FSR) is $\nu_{FSR} = 1.5$ GHz and the finesse is $\mathcal{F} > 250,000$, corresponding to a cavity linewidth $\delta\nu < 6$ kHz from the relationship [127]

$$\delta\nu = \frac{\nu_{FSR}}{\mathcal{F}}. \quad (4.8)$$

We measured the linewidth of our cavity using a ring-down measurement [127]. With the laser locked to the cavity, the RF signal from the synthesizer that multiplexes with

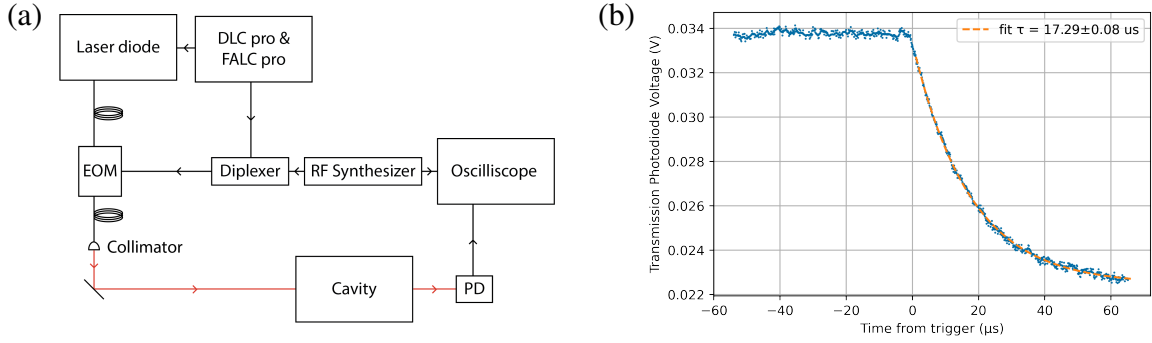


Figure 4.13. (a) Ring-down measurement setup diagram. Components relevant to locking but not the ring-down measurement omitted. See Fig. 4.11 for the full PDH lock diagram. **(b)** Ring down measurement gives a fineness $\mathcal{F} = 163,000 \pm 800$.

the FALC pro RF signal was interrupted, which triggered an oscilloscope to measure the voltage of the PD on the transmission side of the cavity. Without the synthesizer providing the frequency shift on the order of 10s to 100s of MHz that matches the laser frequency to the cavity resonance, the laser light is unable to transmit through the cavity. The only light that exits the cavity once the shift turns off is light that was already resonating within the cavity and slowly leaking out. The exponential decay is recorded on the oscilloscope and quantified by the cavity lifetime τ where

$$\tau = \frac{1}{2\pi\delta\nu}. \quad (4.9)$$

See Fig. 4.13 for a diagram of the measurement setup and measurement result. Measuring τ gives a cavity linewidth of $9.21 \pm 0.04 \text{ kHz}$ and a finesse of $163,000 \pm 800$ (Fig. 4.13b). The measured finesse is lower than expected, which can be a result of a variety of undetermined issues including imperfections on the cavity mirrors or misalignment of the cavity mirrors during installation. Given the sensitive nature of the equipment construction, this was not studied any further as the measured linewidth is sufficient to lock the laser frequency to a narrow linewidth [130].

To determine the quality of the lock, we scanned the 729 nm laser across the carrier transition with a lowered 729 nm beam power to remove the effect of

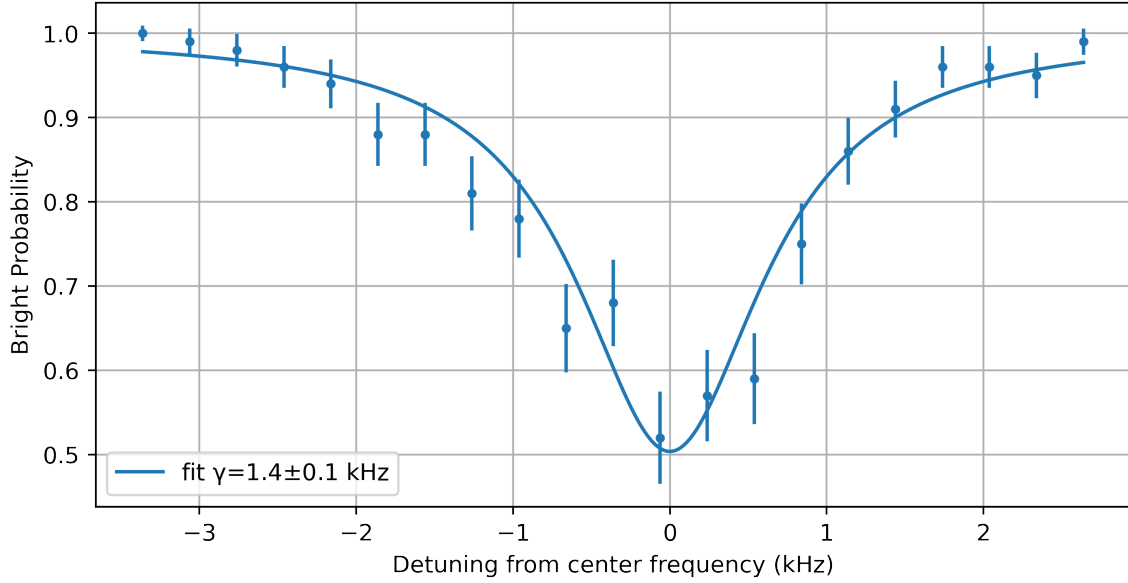


Figure 4.14. Frequency scan of 729 nm carrier transition with an interrogation time of 10 ms and lowered 729 nm power.

power broadening on the linewidth until further lowering of power and increasing of interrogation time did not decrease the width of the observed transition. We settled on an interrogation time of 10 ms to make sure we were not limited by power broadening. This provides an upper bound for the locked 729 nm linewidth because we can not separate the noise due to the magnetic field fluctuations from the laser linewidth. A greater reduction of the linewidth is possible [126], but is only necessary for frequency standard or optical qubit applications. By optimizing the modulation, PID parameters, and the alignment into the cavity, we lock the 729 nm laser to a linewidth upper bounded by 1.4 ± 0.1 kHz (Fig. 4.14). This setup maintains the lock for days at a time needing only to reset the diode laser cavity piezo voltage manually when the unlimited integrator voltage on the FALC pro reaches its limit or the laser goes multimode.

4.4.2 Polarization, Shifts, and Rabi Frequencies. There are ten allowed carrier electric quadrupole transitions for $S_{1/2} \leftrightarrow D_{5/2}$: two $|\Delta m| = 0$, four $|\Delta m| = 1$, and

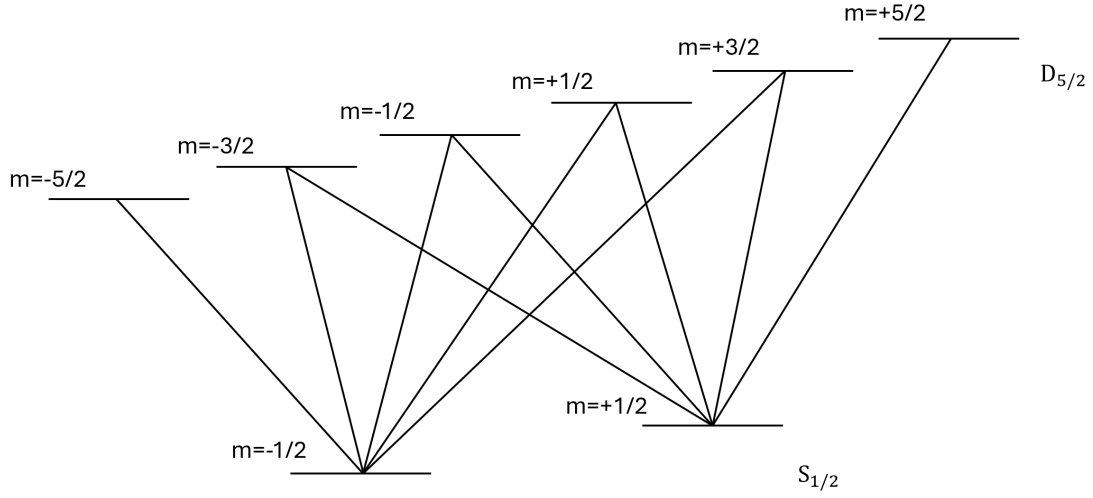


Figure 4.15. All ten allowed transitions for $S_{1/2} \leftrightarrow D_{5/2}$.

four $|\Delta m| = 2$ (Fig. 4.15) [131], where Δm denotes the difference in angular momentum, m , between two states.

The relative coupling strength of any transition, $g^{(\Delta m)}$, depends on the angle ϕ between the laser beam and the magnetic field and the angle γ between the polarization and magnetic field vector projected into the plane of incidence [131]:

$$\begin{aligned} g^{(0)}(\phi, \gamma) &= \frac{1}{2} |\cos \gamma \sin 2\phi| \\ g^{(\pm 1)}(\phi, \gamma) &= \frac{1}{\sqrt{6}} |\cos \gamma \cos 2\phi + i \sin \gamma \cos \phi| \\ g^{(\pm 2)}(\phi, \gamma) &= \frac{1}{\sqrt{6}} |\cos \gamma \sin 2\phi + i \sin \gamma \sin \phi|. \end{aligned} \quad (4.10)$$

See Fig. 4.16 for a diagram of the angles ϕ and γ and the heat maps for the three types of transitions. The Rabi frequency of any transition can be calculated as [132]

$$\Omega_{m,m'} = \left| \frac{e E_0}{2\hbar} \langle S_{1/2} | r^2 \mathbf{C} | D_{5/2} \rangle \sum_{q=-2}^2 \begin{pmatrix} 1/2 & 2 & 5/2 \\ -m & q & m' \end{pmatrix} g^{(\Delta m)} \right|, \quad (4.11)$$

where e is the elementary charge, E_0 is the electric field strength of the laser,

$\langle S_{1/2} | r^2 \mathbf{C} | D_{5/2} \rangle$ is the reduced matrix element [133], the summation terms are Wigner 3-j symbols, and $g^{(\Delta m)}$ comes from Eq. (4.10).

The frequency shift of the transitions due to the magnetic field can be calculated simply from the energy shift via the Zeeman effect:

$$\Delta E = g_j \mu_B B m_j, \quad (4.12)$$

where g_j is the Landé factor, μ_B is the Bohr magneton, and $|\vec{B}| = B$ is the strength of the external magnetic field. The frequency shift from the Zeeman effect is then simply $\Delta f = \Delta E / \hbar$. See Table 4.1 for the Zeeman shifts at our magnetic field strength.

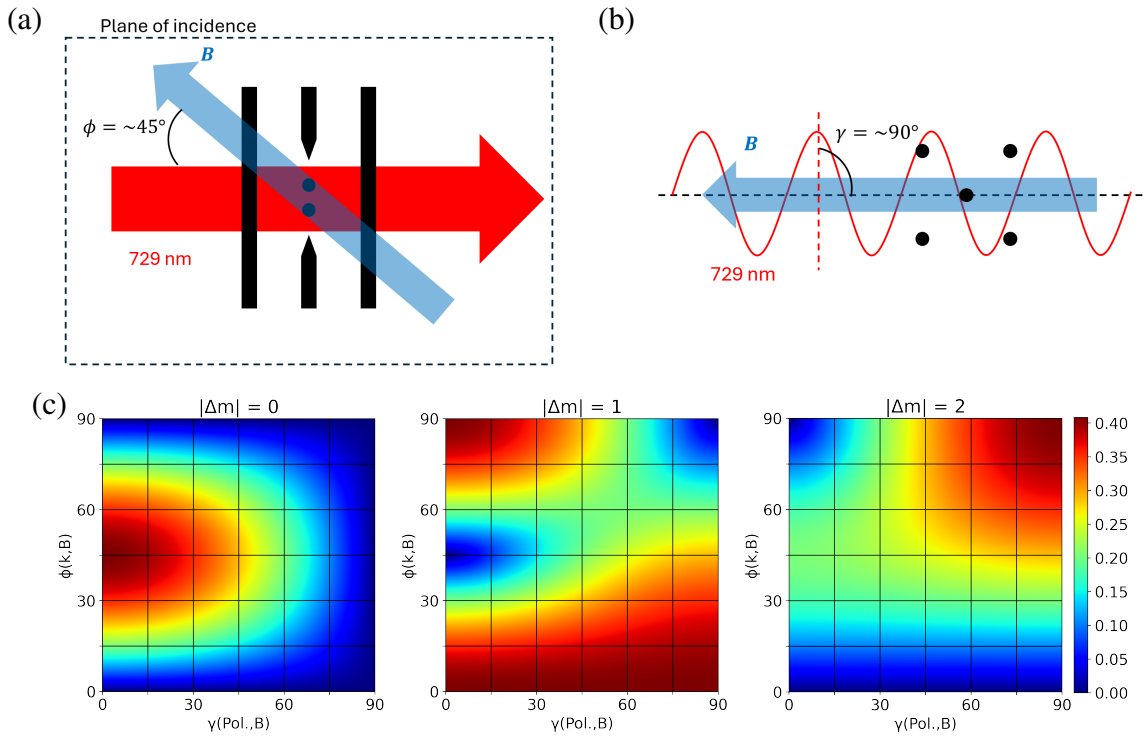


Figure 4.16. **(a)** Front view of trap showing ϕ : the angle between the 729 nm laser beam and direction of ion trap magnetic field, $\vec{B}$. **(b)** Top view of trap showing γ : the angle between the horizontal polarization of the 729 nm laser beam and the projection of $\vec{B}$ on the plane of incidence. **(c)** Relative coupling strength for the allowed $S_{1/2} \leftrightarrow D_{5/2}$ transitions parameterized as a function of angles ϕ and γ .

In our setup, the 729 nm laser beam enters the trap parallel to the table top with the intention to intersect $\vec{B}$ at $\sim 45^\circ$ (Fig. 4.16a). The angle γ is approximately 90° given the laser beam has horizontal polarization (Fig. 4.16b). Measuring the Rabi frequencies

Table 4.1. $S_{1/2} \leftrightarrow D_{5/2}$ magnetic field sensitivities, expected Zeeman shifts at our B-field of 1.567(2) G, measured Zeeman shifts, and relative couplings from Clebsch–Gordan (C-G) coefficients. Angular momentum for $|S_{1/2}, m\rangle$ is represented by m_S and $|D_{5/2}, m\rangle$ by m_D . Measured shifts and expected shifts for $|\Delta m| = 1, 2$ are within error of each other, while the two $|\Delta m| = 0$ transitions (the weakest driven transitions for our setup) are three and two standard deviations apart.

m_S	m_D	Sensitivity (MHz/G)	Expected Zeeman Shift (MHz)	Measured Zeeman Shift (MHz)	Relative Coupling ([C-G] ² /2)
-1/2	-5/2	-2.80	-4.387(3)	-4.375(2)	5/30
-1/2	-3/2	-1.12	-1.755(3)	-1.748(3)	4/30
-1/2	-1/2	0.56	0.877(3)	0.8845(2)	3/30
-1/2	+1/2	2.24	3.509(3)	3.515(5)	2/30
-1/2	+3/2	3.92	6.141(3)	6.142(9)	1/30
+1/2	-3/2	-3.92	-6.141(3)	-6.135(4)	1/30
+1/2	-1/2	-2.24	-3.509(3)	-3.504(3)	2/30
+1/2	+1/2	-0.56	-0.877(3)	-0.8727(4)	3/30
+1/2	+3/2	1.12	1.755(3)	1.760(2)	4/30
+1/2	+5/2	2.80	4.387(3)	4.387(2)	5/30

of all ten allowed carrier transitions gives fitted angle values of $\phi = 47 \pm 2^\circ$ and $\gamma = 85.05 \pm 0.03^\circ$. Looking at Fig. 4.16c, we most strongly drive the $|\Delta m| = 2$ transitions followed closely by $|\Delta m| = 1$ transitions and weakly drive the $|\Delta m| = 0$ transitions.

4.4.3 Off-Resonant Driving. The spectrum of frequencies that drive population between $S_{1/2} \leftrightarrow D_{5/2}$ includes not only the ten carrier transitions, but also many motional sidebands. This makes for a crowded spectrum across ~ 20 MHz since each carrier has red and blue sidebands for both radial secular frequencies at ~ 1.8 and ~ 1.6 MHz. Off-resonant driving is when population is driven via an unintended transition due to close proximity to the intended transition in frequency space. In addition, our laser has noise peaks [130] at 0.5 and 1.5 MHz, which can be seen from the trace of the lock signal in Fig. 4.17, that can unintentionally drive transitions. To measure this effect, a calcium ion was prepared in the ground state and a 729 nm pulse of length equal to the $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$ π -time was applied. The detuning of the

729 nm pulse frequency was varied ± 1.5 MHz from the transition resonant frequency to capture the noise peaks in the PDH lock. Additionally, the same experiment was repeated with pulse lengths of 10x, 100x, and 1000x the π -time. The results of these scans are plotted in Fig. 4.17 along with the trace of the frequency spectrum of the reflected cavity signal that the FALC pro mixes with the RF to produce the error signal. The 1x and 10x π -time scans show insignificant driving of population out of the ground state away from the carrier resonance. In the 100x π -time experiment, $\sim 10\%$ of the population is driven out of the ground state at the frequency corresponding to the first noise peak of the error signal. It is not until the 729 nm pulse is driven for 1000x the π -time that the noise peaks become well-resolved with over 40% of population being driven out of the ground state. For experiments, this means that 729 nm π -pulses would not off-resonantly drive population via transitions within 1.5 MHz above the percent level. This is important for methods and experiments that rely on the 729 nm to drive population with high fidelity (see Chapters V and VI).

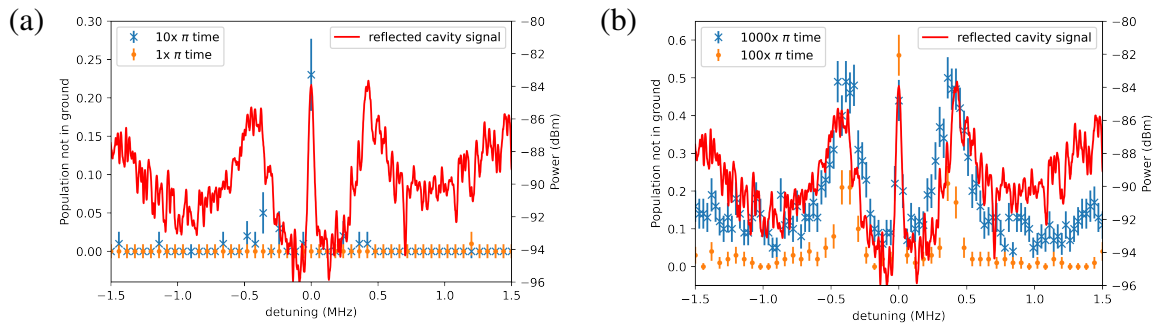


Figure 4.17. Population transfer from off-resonant driving overlaid on the trace of the frequency spectrum of the reflected cavity signal (red solid lines). **(a)** Pulse times are set to the π -time (yellow dots) and 10x the π -time (blue crosses). **(b)** Pulse times are set to 100x (yellow dots) and 1000x (blue crosses) the π -time.

CHAPTER V

MID-CIRCUIT SYMPATHETIC COOLING

The work contained in this chapter was performed by myself and Dr. Jameson O'Reilly with writing done by myself with editorial assistance from Dr. O'Reilly. The simulated SPAM error tables for *mg* crystal readout were produced by Gabe Gregory.

The optical, metastable, and ground (*omg*) architecture [27] combines the dual functionality of multi-species ion chains [25, 134] with the experimental simplicity of single species traps through the use of long-lived metastable states in several species including calcium. 'Dual functionality,' an important feature of scalable quantum devices, refers to the ability to perform dissipative operations like readout, sympathetic cooling [80, 135], and mid-circuit state readout [136] in parallel with long-term storage of information [24]. Dual functionality in multi-species ion chains is possible due to the different transitions between states in each species that require different sets of laser frequencies to address. These transitions between species are spectrally isolated with linewidths many orders of magnitude smaller than their differences, which prevents crosstalk, the unintentional addressing of individual ions. This comes with the drawback that roughly twice as many laser systems are needed for a two-species ion trap compared to a single species trap. This adds significant overhead not only in procuring these laser systems but also in engineering appropriate optical pathways and system maintenance. Additionally, a mismatch in charge-to-mass ratios between the different species results in decoupling of motional states, leading to inefficient sympathetic cooling and long times for entangling gates performed via a shared motional mode [137]. The use of a single species in the *omg* architecture means only one set of laser systems are required and there is no mass differential between ions resulting in efficient sympathetic cooling and shorter entangling gates.

An important aspect of any trapped ion system is heating. High temperature ions make for lower fidelity operations [50, 138] and needs to be compensated. Heating can exist for a variety of reasons. Long circuits involving many qubits and ions require mid-circuit re-cooling to mitigate background heating, referred to as idle or anomalous heating [139], which is caused by stray electric fields. Implementing long circuits in the quantum charge-coupled device (QCCD) architecture [14] relies on the ability to transfer, or shuttle, chains of ions between ‘zones’ on a surface trap, which results in excitations of the motional modes [26]. To prevent, detect, and correct errors during these long circuits, the dissipative ion can be used to non-destructively determine the logic qubit state via fluorescence detection for error correction[29, 30]. Fluorescence detection by driving the cycling transition also heats the ions from momentum kicks from absorption and emission of photons. Sympathetic cooling in dual species systems allows for cooling without compromising logic qubits [80, 135] because the transitions of the cooling ion are different from the logic ion. This is important for quantum information in ion trap systems with many ions to compensate the heating caused by essential processes. The dual species functionality of the *omg* architecture allows for sympathetic cooling, which we demonstrate by compensating heating effects with our cooling methods without limiting the ability to use logic qubits for quantum information applications.

In this chapter, we demonstrate state preparation and measurement (SPAM) and control of a two-ion crystal with one qubit encoded in the ground state and one qubit encoded in a metastable manifold, which we call an *mg* crystal. We also show that the metastable qubit can be sympathetically cooled by the ground qubit without a significantly negative impact on the coherence of the metastable qubit.

5.1 *mg* Crystal Preparation and Measurement

In multi-species ion-chains, the ion ordering is important to maintain a consistent secular mode structure but can be scrambled by collisions with background gas molecules on the timescale of minutes to hours. This makes swapping of ion positions essential, requiring additional DC potentials to break symmetries and force ions to switch positions [140]. In the *omg* architecture, ion ordering is defined through laser beam addressing and can be programmed at the start of each experimental shot.¹ In this work, we focus on the *mg* crystal: two $^{40}\text{Ca}^+$ ions where one is in the ground state $S_{1/2}$ manifold with $|\uparrow\rangle = |S_{1/2}, m = +1/2\rangle$ and $|\downarrow\rangle = |S_{1/2}, m = -1/2\rangle$ (*g*-qubit) and the other is in the metastable $D_{5/2}$ manifold with $|0\rangle = |D_{5/2}, m = +5/2\rangle$ and $|1\rangle = |D_{5/2}, m = +3/2\rangle$ (*m*-qubit).

5.1.1 Heralding *mg* Crystals. To herald an *mg* crystal, we first Doppler and EIT cool both ions and prepare them in $|\uparrow\rangle$ through optical pumping. Next, we apply a global 729 nm $\pi/2$ -pulse on the $|\uparrow\rangle \leftrightarrow |0\rangle$ or $|\uparrow\rangle \leftrightarrow |1\rangle$ transition depending on which state we want to prepare the *m*-qubit in. For a properly calibrated pulse time, each ion will be in an equal superposition of $|\uparrow\rangle$ and $|0(1)\rangle$. We collapse the wavefunction by performing a fluorescence check (FC) using the 397 and 866 nm lasers, see Section 4.1.2. A *g*-qubit fluoresces during the FC while any *m*-qubits remain dark, so we repeat the 729 nm $\pi/2$ -pulse and fluorescence detection whenever we detect either zero or two bright ions. The probability of heralding an *mg* crystal as a function of the number of heralding attempts n and probability of transferring one ion from $|\uparrow\rangle$ to $|0\rangle$ ($|1\rangle$) or vice versa with a pulse α is $p_{mg} = 1 - (\alpha)^n$. FC pulses are 1 ms, so the metastable state (natural lifetime of 1.168(9) s [102]) has a small probability ($< 0.1\%$) of decaying in this time frame. This

¹“There’s no trick to it, it’s just a simple trick.” — Brad Goodman, *The Simpsons* S5E7 “Bart’s Inner Child”

error does not build up with each successive FC since we reattempt heralding after each failed check.

Fig. 5.1 shows real images of two-ion crystals in our trap, their corresponding *mg* configurations, and a histogram of the photon counts from the photomultiplier tube (PMT) used to determine the configuration. The camera is used for alignment of the imaging system and is a high level troubleshooting tool. It is possible to integrate into experiments [141], but this requires developing a real-time image processing software tool and would result in a lower signal-to-noise ratio (SNR). The PMT has a high SNR for individual photon detections because it uses cascade multipliers to increase the voltage produced from each photon and is simple to integrate into experiments because it produces a transistor-transistor logic (TTL) voltage signal and does not require any imaging processing after detection. Additionally, the secular modes in *mg* crystals do not depend on their ordering and we address them with global beams, so none of the work presented here requires knowing the left/right qubit designation.

Our fluorescence detection with the PMT, which discriminates between zero, one, and two bright ions, is fundamentally limited to a fidelity of $> 99.999\%$ due to the overlap of the Poissonian distributions for the number of photons emitted from a cycling transition by a single bright ion and two bright ions. The Poissonian distribution is

$$P(k, \lambda) = \frac{\lambda^k e^{-\lambda}}{k!}, \quad (5.1)$$

where k is the number of instances (photons per FC) and λ is the expected number of instances. The Poissonian distribution is defined for integer values of k only. Given a stationary laser frequency and a FC period of 1 ms, the expected number of photons for zero bright ions is $\lambda_0 = 1$ count, one bright ion $\lambda_1 = 110$ counts, and two bright ions $\lambda_2 = 229$ counts. The overlap A of two distributions with λ_a and $\lambda_b > \lambda_a$ is $A = F(c, \lambda_b) + 1 - F(c, \lambda_a)$ where $F(k, \lambda)$ is the cumulative distribution function

Table 5.1. Probability table for detecting each type of two-ion configuration. Each column comes from 10,000 repetitions of an experiment.

Configuration	Unheralded	Heralded
Two m -qubits	23.0(4)%	0.00(0)%
mg Crystal	51.8(5)%	99.81(4)%
Two g -qubits	25.2(4)%	0.19(4)%

of the Poissonian distribution with expected value λ evaluated at instance k and $c = \lfloor \ln(\lambda_b - \lambda_a) / [\ln(\lambda_b) - \ln(\lambda_a)] \rfloor$ is the crossing point of the two distributions. The overlap of one and two bright ions distributions is < 0.00001 , which gives a fundamental error rate of 0.001%. The overlap of zero and one bright ions is exceedingly smaller. The separation of these distributions can also be adjusted by changing the length of time the fluorescence detection occurs. Decreasing time will decrease instances of leakage out of the $D_{5/2}$ manifold with the drawback that the overlap of the Poissonian distributions will increase. At 1 ms FC periods, the probability of leakage is 0.09% [102]. At 0.5 ms, the leakage probability is 0.04% while overlap becomes 0.002, which is an error rate of 0.2%, larger than the leakage rate at 1 ms.

Following a properly calibrated 729 nm $\pi/2$ -pulse, an FC will detect two g -qubits 25% of the time, two m -qubits 25% of the time, and an mg crystal 50% of the time. In practice, we do not detect an mg crystal exactly 50% of the time (see Table 5.1 for a typical detection matrix) due to the decay of the m -qubit out of $D_{5/2}$ manifold at a rate consistent with the natural lifetime [102] with the 1 ms fluorescence check.

5.1.2 mg Readout Scheme. To determine the results of any mg crystal experiment, a readout scheme is needed to detect each qubit state. Three FCs along with a series of 729 nm carrier π -pulses allows for the determination of each qubit state and flagging of errors including population decay from the metastable manifold, also known as “leakage.”

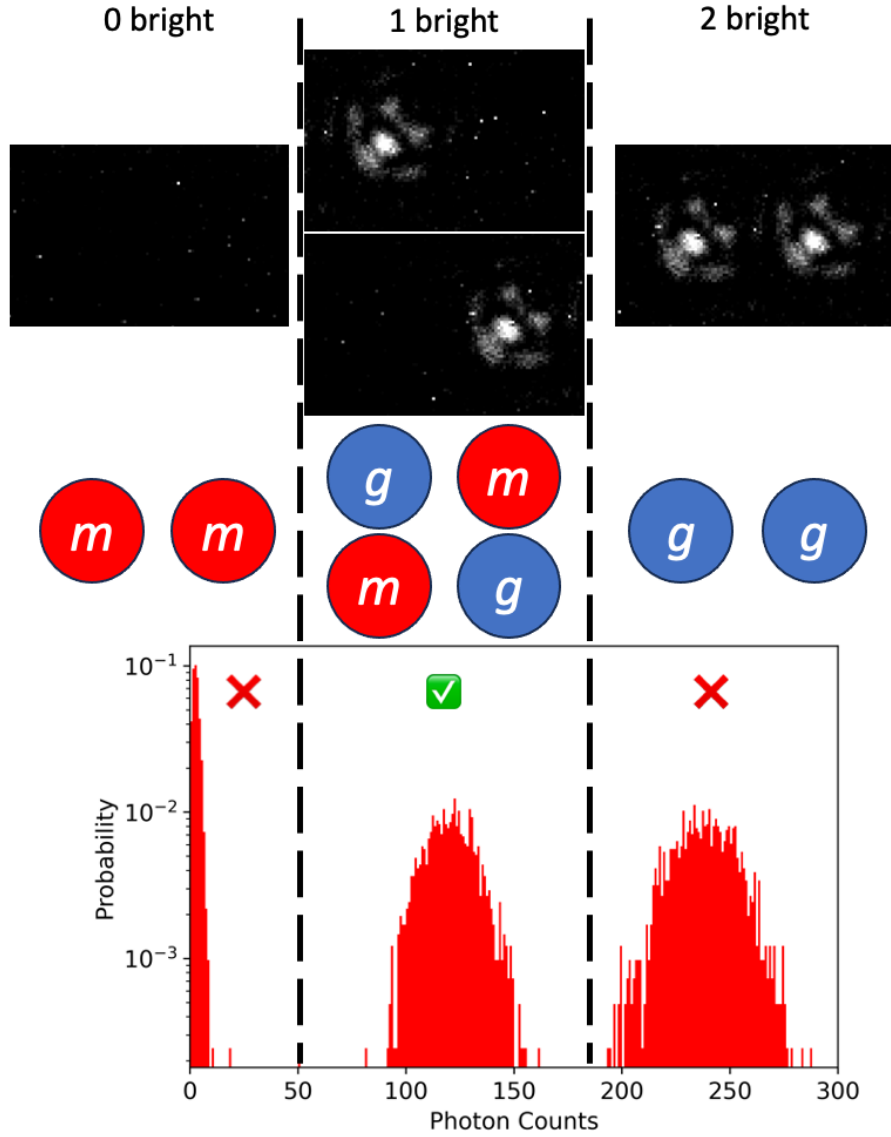


Figure 5.1. Real images of two-ion configurations in our trap corresponding to preparation in the $D_{5/2}$ manifold (dark ions) and $S_{1/2}$ manifold (bright ions). Below is a histogram of photon counts detected during each 1 ms fluorescence period from an example experiment involving two ions. Dashed lines are the set thresholds for discriminating between zero and one bright ions and one and two bright ions.

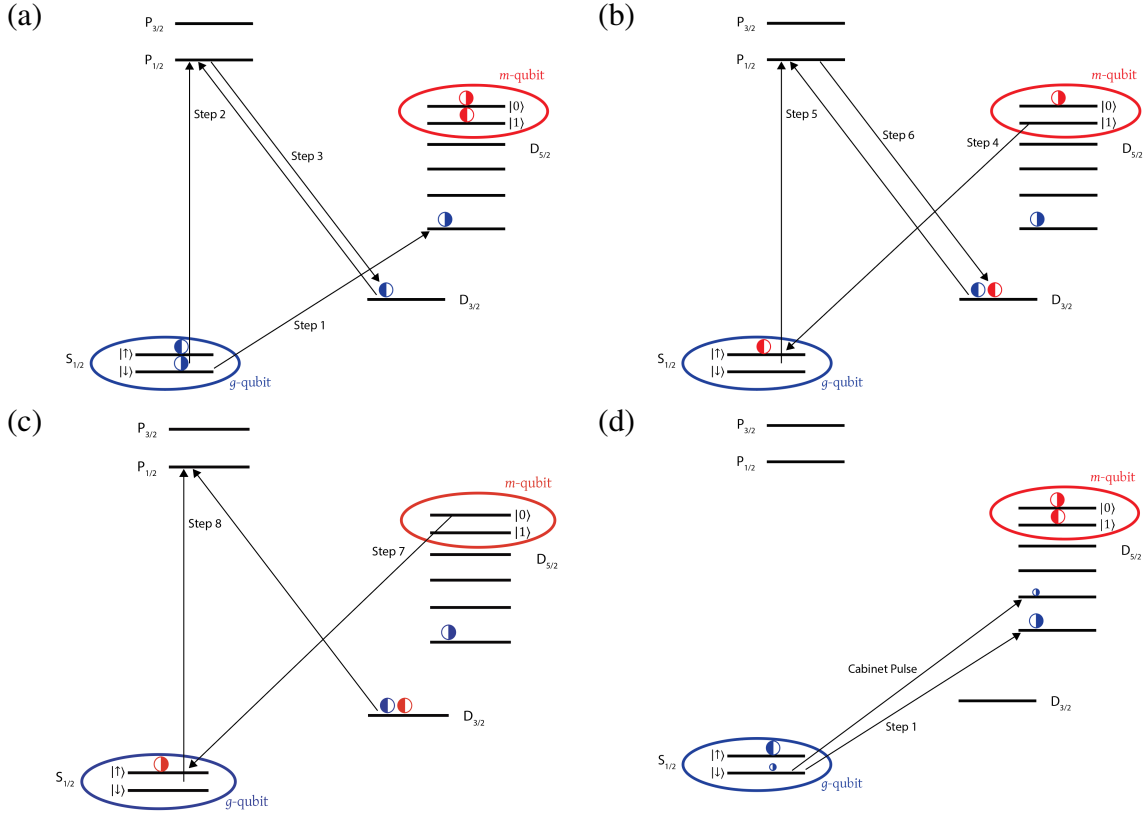


Figure 5.2. Diagrams of the mg crystal readout scheme. Half filled in circles represent populations in one of the two states for each qubit and where they are moved to in each step of the readout scheme. **(a)** Detection of the g -qubit: $|\downarrow\rangle$ population is shelved and remaining $|\uparrow\rangle$ population is detected and reshelled in $D_{3/2}$. **(b)** Detection of the m -qubit: $|1\rangle$ population is driven to $S_{1/2}$, detected, and shelved in $D_{3/2}$. **(c)** Final check of the m -qubit population: $|0\rangle$ population is driven to the ground state and detected. **(d)** Cabinet shelving of $|\downarrow\rangle$ population before the first FC. The smaller half-circle represents the probability that the first 729 nm π -pulse to $|D_{5/2}, m = -5/2\rangle$ did not fully shelve population from $|\downarrow\rangle$.

The basic detection process is as follows, starting with determining the g -qubit state, as diagrammed in Fig. 5.2a:

1. π -pulse on $|\downarrow\rangle \leftrightarrow |D_{5/2}, m = -5/2\rangle$ shelves any population in $|\downarrow\rangle$.
2. The first FC is performed.
3. $S_{1/2}$ population is optically pumped to $D_{3/2}$ by applying just the 397 nm beam.

If one bright ion is detected during the FC, the g -qubit was in $|\uparrow\rangle$. Otherwise, the 729 nm π -pulse shelved $|\downarrow\rangle$ population in $|D_{5/2}, m = -5/2\rangle$ where it will remain for the rest of the full readout sequence. Optically pumping any $S_{1/2}$ population into the long-lived metastable $D_{3/2}$ manifold ensures that it does not switch places with the m -qubit populations during subsequent 729 nm pulses. This is done by using the 397 nm laser to drive population to $P_{1/2}$. Without the 866 nm to drive population spontaneously decaying to $D_{3/2}$ back to $P_{1/2}$, population builds up in the $D_{3/2}$ manifold with over 99% probability in 10 μ s and will remain there until the next FC.

Next, the m -qubit state is determined, as diagrammed in Fig. 5.2b:

4. π -pulse on $|\uparrow\rangle \leftrightarrow |D_{5/2}, m = +3/2\rangle$ moves any population in $|1\rangle$ to the $S_{1/2}$ manifold.
5. The second FC is performed.
6. $S_{1/2}$ population is optically pumped to $D_{3/2}$ by applying just the 397 nm beam.

If one more bright ion is detected than was in the first FC (recall that population in $D_{3/2}$ will fluoresce), the m -qubit was in $|1\rangle$. Lastly, population remaining in $|0\rangle$ is detected, as diagrammed in Fig. 5.2c:

7. π -pulse on $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$ moves any population in $|1\rangle$ to the $S_{1/2}$ manifold.
8. The third FC is performed.

This final FC ensures that all metastable population was detected. The FC results and their corresponding state detections can be found in Table 5.2.

‘Leakage’ refers to the metastable population decaying because of the natural lifetime or spontaneously scattering a laser beam and is an especially damaging error

Table 5.2. Corresponding state for a given sequence of bright ion detections in mg crystals. ‘Leakage’ refers to the metastable population decaying because of spontaneous decay due to the natural lifetime or scattering of a laser beam. All other combinations of detections are categorized as a generic ‘error.’

Fl. Check #1	Fl. Check #2	Fl. Check #3	Result
0	0	1	$ \downarrow, 0\rangle$
0	1	1	$ \downarrow, 1\rangle$
1	1	2	$ \uparrow, 0\rangle$
1	2	2	$ \uparrow, 1\rangle$
0	0	0	leakage
1	1	1	leakage
2	2	2	leakage

in quantum computation [142]. If $|0\rangle$ or $|1\rangle$ population decays before it is purposefully driven to $S_{1/2}$, it will be detected during the incorrect FC and the error can be flagged, making it a less harmful ‘erasure’ error [143]. Decay to $|\downarrow\rangle$ before $|\downarrow\rangle$ population is shelved will result in zero bright ions for all three FCs if the g -qubit is in $|\downarrow\rangle$ and one bright ion for all three FCs if the g -qubit is in $|\uparrow\rangle$. If m population decays to $|S_{1/2}, m = -1/2\rangle$ after the $|\downarrow\rangle$ shelving pulse or it decays to $|S_{1/2}, m = +1/2\rangle$ before the first FC, then all three FCs will detect two bright ions. Permutations of the three FCs that are not listed in Table 5.2 are flagged as generic “errors.”

5.1.2.1 Cabinet Shelving and Bucket Deshelving. When driving population with a 729 nm π -pulse, we rely on the stability of the laser frequency (Section 4.4.1) and intensity at the ion to transfer population with high fidelity. However, the laser lock is not perfect and the intensity can fluctuate which results in a less than perfect transfer of population. Additionally, calibration of the 729 nm carrier π -time is imperfect, and finite ion temperatures cause damping due to the carrier Rabi frequency dependence on phonon number. To mitigate the resulting non-unity π -pulse fidelities, we incorporate additional π -pulses when attempting to shelve and deshelve qubit population in the mg crystal readout scheme.

Cabinet shelving is used in step 1 of the mg readout scheme but was suppressed in the previous section for clarity. A single 729 nm π -pulse on $|S_{1/2}, m = -1/2\rangle \leftrightarrow |D_{5/2}, m = -5/2\rangle$ sends $|\downarrow\rangle$ population to $D_{5/2}$ so that only $|\uparrow\rangle$ population will fluoresce. To increase the probability of shelving all $|\downarrow\rangle$ population, a second 729 nm carrier π -pulse on $|S_{1/2}, m = -1/2\rangle \leftrightarrow |D_{5/2}, m = -3/2\rangle$ can be used to remove any leftover population from the first pulse. This process can be repeated as long as there are available states to shelve to. For our encoding scheme in $^{40}\text{Ca}^+$, there are four states available: $m = \{-5/2, -3/2, -1/2, +1/2\}$, with $+3/2$ and $+5/2$ reserved for the m -qubit. Error on the π -pulse time calibration for $|\Delta m| = \{1, 2\}$ gives a population transfer of $> 95.3\%$ for a single pulse, so two pulses gives a population transfer of $> 99.7\%$.

As a reversal of cabinet shelving, bucket *deshelving* can be used to ensure that all targeted population gets deshelled from $D_{5/2}$. After an initial 729 nm carrier π -pulse to deshelve $|1\rangle$ population in step 3 or $|0\rangle$ population in step 5, the 397 nm can again be turned on for 10 μs to pump $S_{1/2}$ population to $D_{3/2}$, which allows for a second 729 nm carrier pulse to drive any population not yet driven out of $D_{5/2}$. Once again, two 729 nm carrier pulses each with $> 95.3\%$ probability of driving population makes the probability of properly deshelling m -qubit population to be $> 99.7\%$.

5.1.3 mg Crystal SPAM. To characterize our state preparation and measurement (SPAM) fidelity for mg crystals, each of the four combinations of the m and g -qubit states were prepared 10,000 times and we performed the mg readout scheme described above. To prepare the $|\downarrow\rangle$ state, an RF π -pulse on the g -qubit is performed after heralding an mg crystal. Each state is prepared and measured with a fidelity of over 98% after subtracting flagged leakage and error events (Table 5.3).

Table 5.3. State preparation and measurement fidelity of each of the four qubit state combinations in an *mg* crystal. Each fidelity is calculated from 10,000 *mg* crystal preparations. ‘Leakage’ refers to the metastable population decaying before the *m*-qubit population could be determined. All other combinations of detections are categorized as a generic ‘error.’ Fidelity is calculated as the percentage of the correctly prepared states to total prepared states after subtracting out states flagged as a leakage or error.

Prep \ Result	$ \uparrow, 0\rangle$	$ \uparrow, 1\rangle$	$ \downarrow, 0\rangle$	$ \downarrow, 1\rangle$	Leakage	Error	Fidelity (%)
$ \uparrow, 0\rangle$	9661	63	82	0	183	11	98.5(1)
$ \uparrow, 1\rangle$	17	9669	0	77	219	18	99.0(1)
$ \downarrow, 0\rangle$	107	0	9546	60	224	63	98.3(1)
$ \downarrow, 1\rangle$	1	116	15	9600	211	57	98.6(1)

Table 5.4. Simulation of error due to π -pulse time uncertainty on readout probability of each of the four qubit state combinations for an *mg* crystal. ‘Leakage’ refers to the metastable population decaying before the *m*-qubit population could be determined. All other combinations of detections are categorized as a generic ‘error.’

Prep \ Result	$ \uparrow, 0\rangle$	$ \uparrow, 1\rangle$	$ \downarrow, 0\rangle$	$ \downarrow, 1\rangle$	Leakage	Error	Fidelity (%)
$ \uparrow, 0\rangle$	99.8	0.0	0.0	0.0	0.1	0.1	100.0
$ \uparrow, 1\rangle$	0.1	99.8	0.0	0.0	0.0	0.1	99.9
$ \downarrow, 0\rangle$	0.0	0.0	100.0	0.0	0.0	0.0	100.0
$ \downarrow, 1\rangle$	0.0	0.0	0.0	100.0	0.0	0.0	100.0

5.1.3.1 *mg* Crystal SPAM Error Sources. Error in *mg* crystal SPAM was modeled with two sources: 729 nm π -pulse time uncertainty and unequal illumination of the ions with 729 nm light. These sources were considered separately in simulations.

Error from 729 nm π -pulse time uncertainty was experimentally determined by driving population with each of the three $S_{1/2} \leftrightarrow D_{5/2}$ transition frequencies used for *mg* crystal SPAM: $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$, $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +3/2\rangle$, and $|S_{1/2}, m = -1/2\rangle \leftrightarrow |D_{5/2}, m = -5/2\rangle$. The average error for these three transitions used to generate the simulated results in Table 5.4 is 2.3%. Infidelity due to 729 nm π -pulse time uncertainty is on the order of $\sim 0.1\%$.

Table 5.5. Simulation of readout error due to unequal illumination of the ions by 729 nm light for each of the four qubit state combinations in an *mg* crystal. Unevenness was simulated as a 2% difference in Rabi frequency between each ion. ‘Leakage’ refers to the metastable population decaying before the *m*-qubit population could be determined. All other combinations of detections are categorized as a generic ‘error.’

Prep \ Result	$ \uparrow, 0\rangle$	$ \uparrow, 1\rangle$	$ \downarrow, 0\rangle$	$ \downarrow, 1\rangle$	Leakage	Error	Fidelity (%)
$ \uparrow, 0\rangle$	99.9	0.0	0.0	0.0	0.1	0.0	100.0
$ \uparrow, 1\rangle$	0.1	99.8	0.0	0.0	0.0	0.1	99.9
$ \downarrow, 0\rangle$	0.0	0.0	99.6	0.1	0.1	0.2	99.9
$ \downarrow, 1\rangle$	0.0	0.0	0.1	99.5	0.1	0.3	99.9

Unequal illumination, or unevenness ϵ , of the ions with 729 nm light is defined as the percent difference of the Rabi frequency for each ion from the average Rabi frequency Ω_{ave} of the two ions. Unevenness was experimentally determined by driving population between $|\uparrow\rangle$ and $|0\rangle$ for times t long enough to see changes in maximum population transfer and fitting this to

$$P_{\uparrow}(t) = \left(\frac{1}{2} \left[\sin\left(\frac{\Omega_+ t}{2}\right) + \sin\left(\frac{\Omega_- t}{2}\right) \right] \right)^2 \quad (5.2)$$

where $\Omega_{\pm} = \Omega_{ave}(1 \pm \epsilon)$. For a typical unevenness of 2%, we find the results presented in Table 5.5, implying an error due to unequal illumination of 0.1%.

The remaining infidelity of $\sim 1.2\%$ is largely due to the *g*-state being nominally prepared in $|\downarrow\rangle$ and improperly readout in $|\uparrow\rangle$. This is from the infidelity of the RF π -pulse after heralding the *mg* crystal that moves population in $|\uparrow\rangle$ to $|\downarrow\rangle$ with an error of 1.2%. The *m*-qubit being improperly readout is likely due to off-resonant driving of the deshelling 729 nm π -pulses which drive $|0\rangle$ population to the ground manifold while attempting to deshelve $|1\rangle$.

5.1.4 *mg* RF Control. To show that *m* and *g*-qubits can be manipulated independently, Fig. 5.3 displays Rabi flopping of each qubit driven by RF pulses. After heralding an *mg* crystal, radiation is applied at the transition frequency for each qubit:

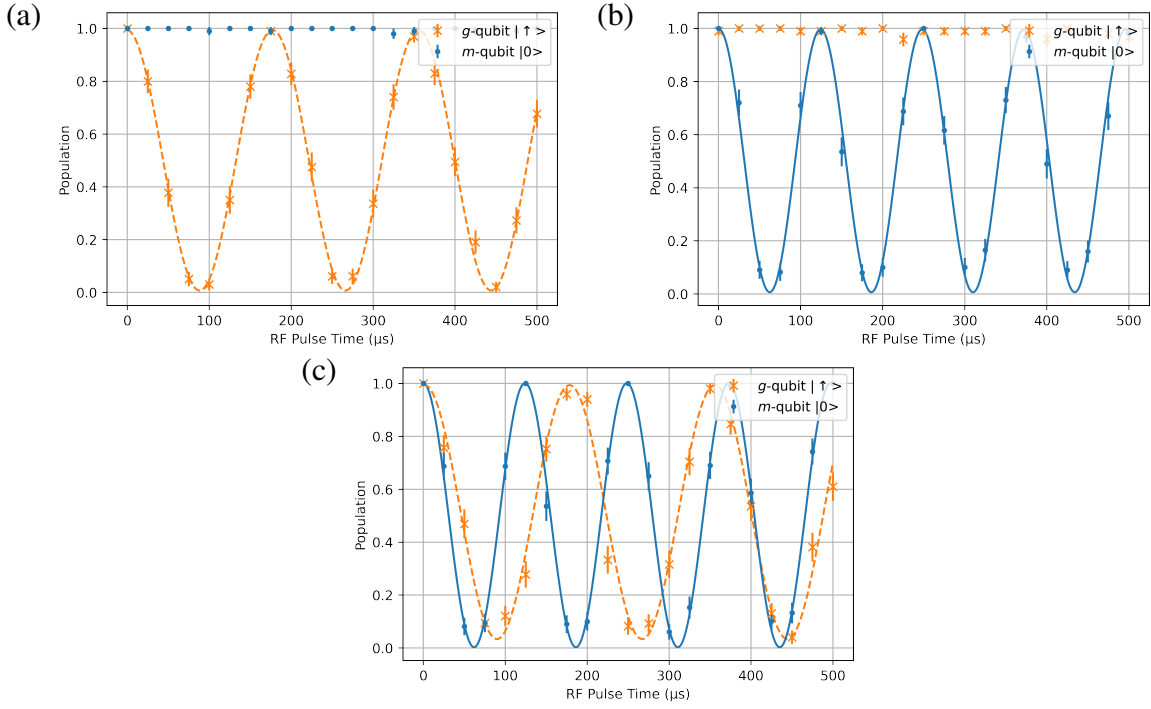


Figure 5.3. Individual RF Rabi flopping of (a) the g -qubit and (b) the m -qubit, as well as (c) simultaneous control of both the m and g -qubits. Lines are fits of the data to sinusoids.

~ 4.4 MHz for the g -qubit and ~ 2.6 MHz for the m -qubit. The population of each qubit is then determined using the mg readout scheme. As can be seen in Fig. 5.3a and b, driving one qubit can be done without effecting the population of the other qubit. Both qubit populations can also be driven in the same experiment, as seen in Fig. 5.3c, by applying each RF pulse in sequence.

5.2 Heating and Sympathetic Re-Cooling

There are several reasons why heating occurs in ions. Even when no operations are explicitly performed, trapped ions experience idle or anomalous heating [139]. Stray electric fields at the position of the ion couple to its charge resulting in fluctuating forces that can match the secular frequency and sidebands of the ion, which can result in measurable temperature increases. As trap size decreases and the length of time ions are left idle, this effect becomes larger [50]. Some ion trap constructions

allow for transportation of individual or chains of ions across 2D chips, referred to as shuttling [14, 144, 145], which is an essential feature of the quantum charged coupled device (QCCD) computer architecture [13, 14]. The mismatch in ion mass when using dual species, however, makes shuttling inefficient and requires additional stray field compensation voltages and quadrupole adjustments to shuttle as well as single ion species [26]. The *omg* architecture allows for simpler single species operation while allowing for dual species functionality. Shuttling chains of single ion species, however, still results in the motion of the ion becoming displaced, *i.e.*, becoming a coherent state, which is an increase in the occupation number in the Fock state basis (see Section 5.2.1 and Eq. (5.7)). This is effectively an increase in ion temperature and must be mitigated through cooling. Fluorescence detection is also used for mid-circuit measurements, which are needed for a variety of reasons including error-correction [29, 30] and characterizing entanglement phase transitions [146]. Fluorescence detection increases the temperature of the ion chain due to the momentum kicks from both absorbing and emitting photons.

In this section, we measure these heating effects to inform the re-cooling needs of the *mg* crystal. We explain how to readout ion temperature via the 976 nm Raman lasers on the *m*-qubit, measure the processes that heat the ion chain (anomalous heating, shuttling, and fluorescence detection), and demonstrate how sympathetic cooling on the *g*-qubit using Doppler and EIT cooling are able to compensate for these heating processes.

5.2.1 Measuring Fock State Distributions via Sideband and Carrier

Driving. When driving a carrier transition, the motional state of the ion is preserved ($\Delta n = 0$). To add or subtract a phonon, we can detune from the carrier transition by the secular frequency and drive a sideband. At non-zero temperatures, the n -dependence of the Rabi frequency causes damping as the electronic populations corresponding to each motional Fock state go out of phase with each other. To measure the temperature of an *mg*

crystal, we drove the m -qubit using the 976 nm Raman lasers (Section 4.3) on either the carrier or the first phonon-adding sideband (referred to as the blue sideband (BSB) from here on) of the center of mass motional mode of the ion crystal.

For a two level system with states $|0\rangle$ and $|1\rangle$ coherently driven via a resonant oscillatory field for a time t , the probability of being in $|0\rangle$ is [63, 147]

$$P_{|0\rangle}(\bar{n}, t) = \sum_{n=0}^{\infty} p_n e^{-\gamma\sqrt{n+1}t} \frac{\sin^2(\Omega_n t)}{2}, \quad (5.3)$$

where the probability of occupying state n for a thermal state is

$$p_n = \frac{1}{1 + \bar{n}} \left(\frac{\bar{n}}{1 + \bar{n}} \right)^n, \quad (5.4)$$

the n^{th} BSB Rabi frequency is

$$\Omega_n = \Omega \eta \sqrt{n+1}, \quad (5.5)$$

where n is the lower Fock state of the two being driven, γ is a phenomenological dissipation factor to account for decoherence, and the Lamb-Dicke parameter is

$$\eta = 2 \cos(\theta) \cos(\phi) \frac{2\pi}{\lambda} \sqrt{\frac{\hbar}{2m\omega}}, \quad (5.6)$$

where θ is the angle between the laser beam and the radial motion of the ions, ϕ is the angle between the laser beam and the axial motion of the ions, λ is the wavelength of the driving field, m is the mass of the trapped ion, and ω is the secular frequency of the motion under consideration. For our 976 nm laser setup, θ and ϕ are both 45° [33].

The motional state of the ion can also be a displaced or coherent state $|\alpha\rangle$ after shuttling [14, 144, 145] (see Section 5.2.3), where α is the displacement parameter. The probability of occupying state n is then [148, 149]

$$p_n = |\langle n | \alpha \rangle|^2 = e^{-|\alpha|^2} \frac{|\alpha|^{2n}}{n!} = e^{-\bar{n}} \frac{\bar{n}^n}{n!} \quad (5.7)$$

where $\bar{n} = |\alpha|^2$.

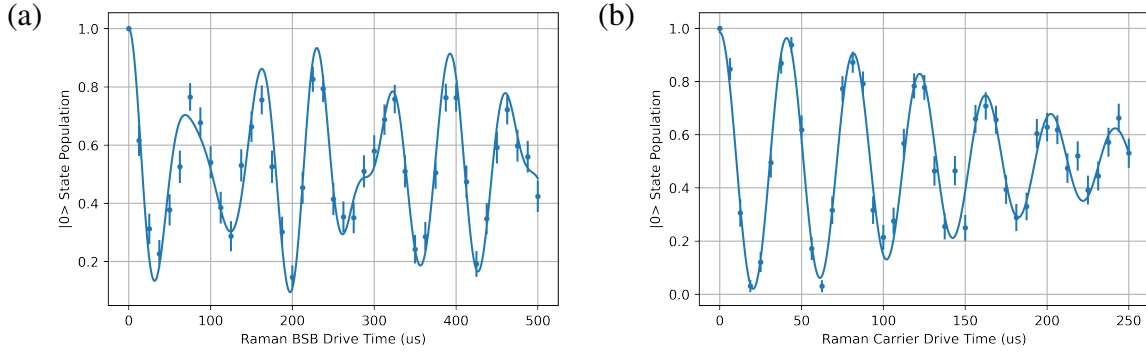


Figure 5.4. Example of data and fits for readout by driving the BSB and carrier transitions with 976 nm Raman beams. **(a)** The BSB data was collected after cooling an *mg* crystal in a thermal state with 1 ms of Doppler cooling and 400 μ s of EIT cooling. The fit gives $\bar{n} = 0.52(5)$. **(b)** The carrier data was collected after driving fluorescence on the *g*-qubit of an EIT limit cooled *mg* crystal for 800 μ s (Section 5.2.4). The fit gives $\bar{n} = 7.0(2)$.

The 976 nm BSB readout of the ion temperature is sensitive for $\bar{n} \lesssim 3$, above which the Rabi flops quickly decay to $P_{|0\rangle} = 0.5$. For better sensitivity at higher temperatures, the 976 nm carrier can be used to determine $\bar{n}$ as its Rabi frequency has a weaker dependence on n . Conversely, carrier driving is insensitive to small $\bar{n}$, which is why BSB driving is performed to determine temperatures below the Doppler cooling limit. As opposed to the BSB which only drives one motional mode, the carrier drives all four radial motional modes for two trapped ions. The probability of being in $|0\rangle$ is then

$$P_{|0\rangle} = \sum_{n_1=0}^{\infty} \sum_{n_2=0}^{\infty} \sum_{n_3=0}^{\infty} \sum_{n_4=0}^{\infty} p_{n_1} p_{n_2} p_{n_3} p_{n_4} \sin^2(\Omega' t) \quad (5.8)$$

where

$$\Omega' = \Omega \prod_i (1 - \eta_i^2 n_i) \quad (5.9)$$

is the carrier Rabi frequency for each motional mode and n_i and η_i correspond to each of the four radial mode occupation numbers and Lamb-Dicke parameters.

Examples of the data collected and fits for both BSB and carrier readout are given in Fig. 5.4.

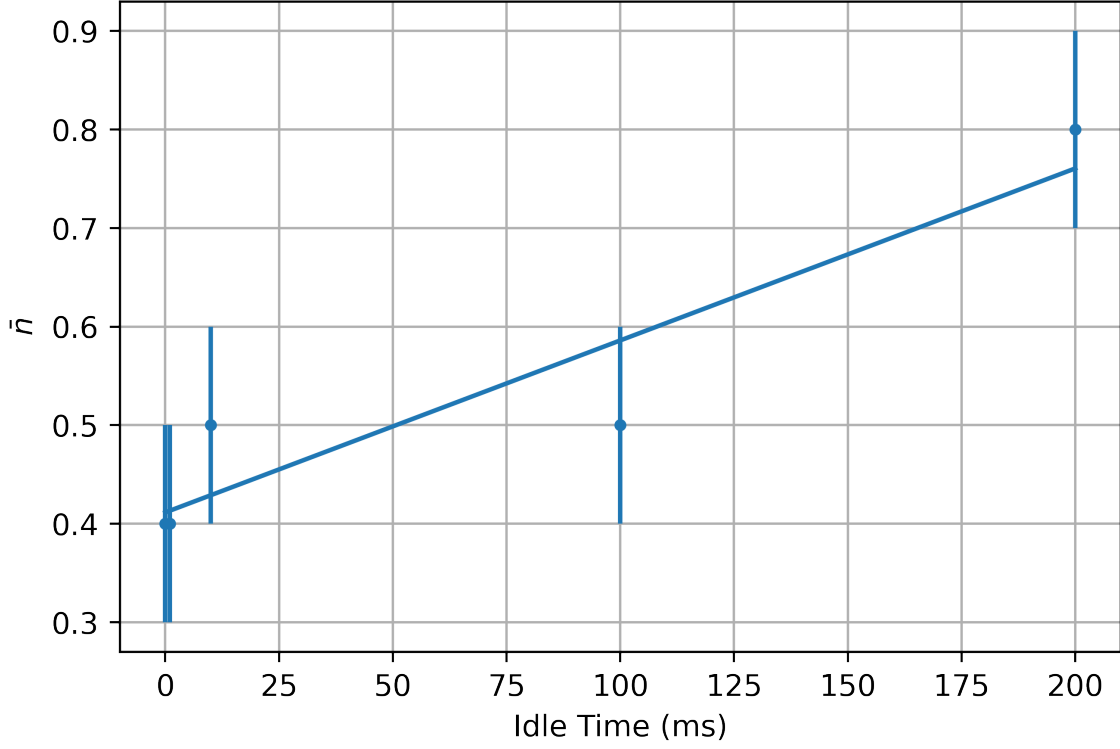


Figure 5.5. Average occupation state $\bar{n}$ for various idle times measured on the m -qubit via BSB readout (Section 5.2.1) fitted to a line with slope $1.7(4)$ quanta/s.

5.2.2 Idle Heating. Idle heating is measured in our trap by varying the time between preparing the mg crystal and driving the m -qubit BSB readout. Fig. 5.5 plots the results of five BSB readout experiments fitted to a line for a heating rate of $1.7(4)$ quanta/s, a reasonable rate given our trap parameters [139].

5.2.3 Displacement Heating. Our trap has no shuttling capabilities, but we can still use the RF drive and squeeze box (Section 2.5.2.1) to displace the motional modes of the ions. The displacement operator $\hat{D}$ acts on a ground state to create a coherent state such that

$$|\alpha\rangle = \hat{D}(\alpha)|0\rangle = \exp\left(-\frac{1}{2}|\alpha|^2\right) \sum_0^\infty \frac{\alpha^n}{\sqrt{n!}} |n\rangle. \quad (5.10)$$

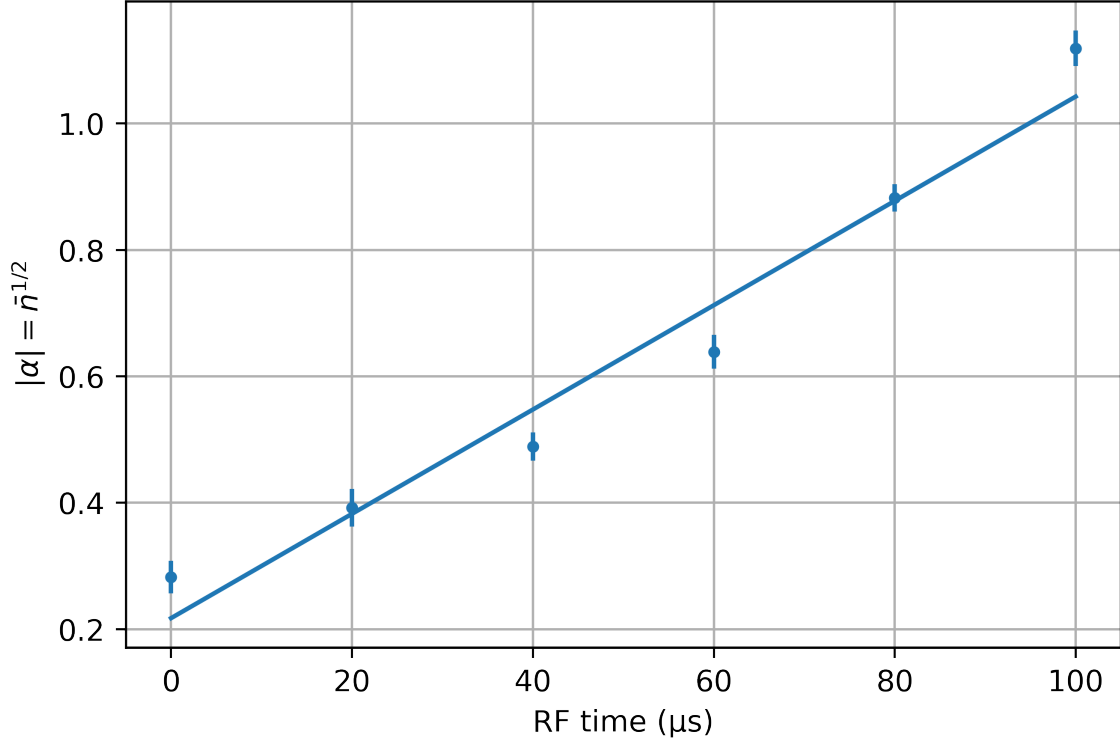


Figure 5.6. BSB readout of displacement $|\alpha|$ due to RF shuttling as a function of the RF pulse time for a fixed RF power. The data fits to line with a slope of $0.0083(8) \mu\text{s}^{-1}$.

Fig. 5.6 shows the results of displacing an mg crystal by increasing the RF drive pulse time for a fixed RF power. Readout is performed using the 976 nm Raman BSB and fit to a coherent state (Eq. (5.7)).

5.2.4 Fluorescence Heating. As a function of time, the heating of the motional state due to fluorescence follows [150]

$$\bar{n}(t) = \left(\bar{n}(0) - \frac{R}{D\hbar\omega} \right) e^{-\Gamma_0 D t} + \frac{R}{D\hbar\omega} \quad (5.11)$$

where ω is the radial secular frequency, Γ_0 is the effective scattering rate, R is the effective recoil energy, and $D = 8\Delta\hbar f_x k^2 / (m\gamma^2(1 + s') + 4m\Delta^2)$ is a Doppler-effect term that takes into account the laser detuning Δ . See [150] for explicit definitions and derivations and Table 5.6 for our values. Fig. 5.7 plots ion temperature measured via carrier flopping against fluorescence measurement time for our trap values. Data was

Table 5.6. Parameters used to calculate fluorescence heating rate for our trap setup. See [150] for explicit definitions and derivations. Symbols for each parameter used in [150] are listed alongside values.

Parameter	Symbol	Value	Units
mass ^{40}Ca	m	6.636×10^{-26}	kg
secular frequency	ω	$2\pi \times 1.8806$	MHz
isotropic emission geometry constant [151]	f_{xx}	1/3	-
absorption geometry constant [151]	f_x	1/2	-
wavelength	λ	397	nm
transition natural linewidth	γ	$2\pi \times 21.56$	MHz
Zeeman splitting	δ_B	$2\pi \times 1.455$	MHz
saturation parameter [152]	s	1.79	-
magnetic field strength	B	1.56	Gauss

taken in two data sets with $\Delta = 2\pi \times \{-7(2), -1(2)\}$ MHz where the uncertainty comes from fitting fluorescence peaks to calibrate the laser lock point.

The black data set (with a solid line from theory) for $\Delta = 2\pi \times -7(2)$ MHz approaches the calculated limit of $\bar{n} = 17_{-3}^{+6}$ in less than 2 ms whereas the red data set (with a dashed line from theory) for $\Delta = 2\pi \times -1(2)$ MHz does not reach a steady-state within the plotted region. The dashed line at long times is calculated to reach the limit $\bar{n} \sim 110$, however, the uncertainty in Δ extends from $2\pi \times -3$ MHz (limit of $\bar{n} \sim 37$) to $+2\pi \times 1$ MHz, which would result in runaway heating and the ion eventually leaving the trap. $\Delta = 2\pi \times -7(2)$ MHz reaches a lower temperature limit faster than $\Delta = 2\pi \times -1(2)$ MHz because Δ determines the average magnitude of phase-space kicks during photon scattering as well as the effective scattering rate Γ_0 [150]. For experiments, we set $\Delta = 2\pi \times -5$ to -10 MHz to quickly reach a lower heating limit. Detuning farther from resonance decreases the rate of emitted photons detected for state discrimination and results in larger overlaps of the Poissonian distributions for zero, one, and two bright ion detection.

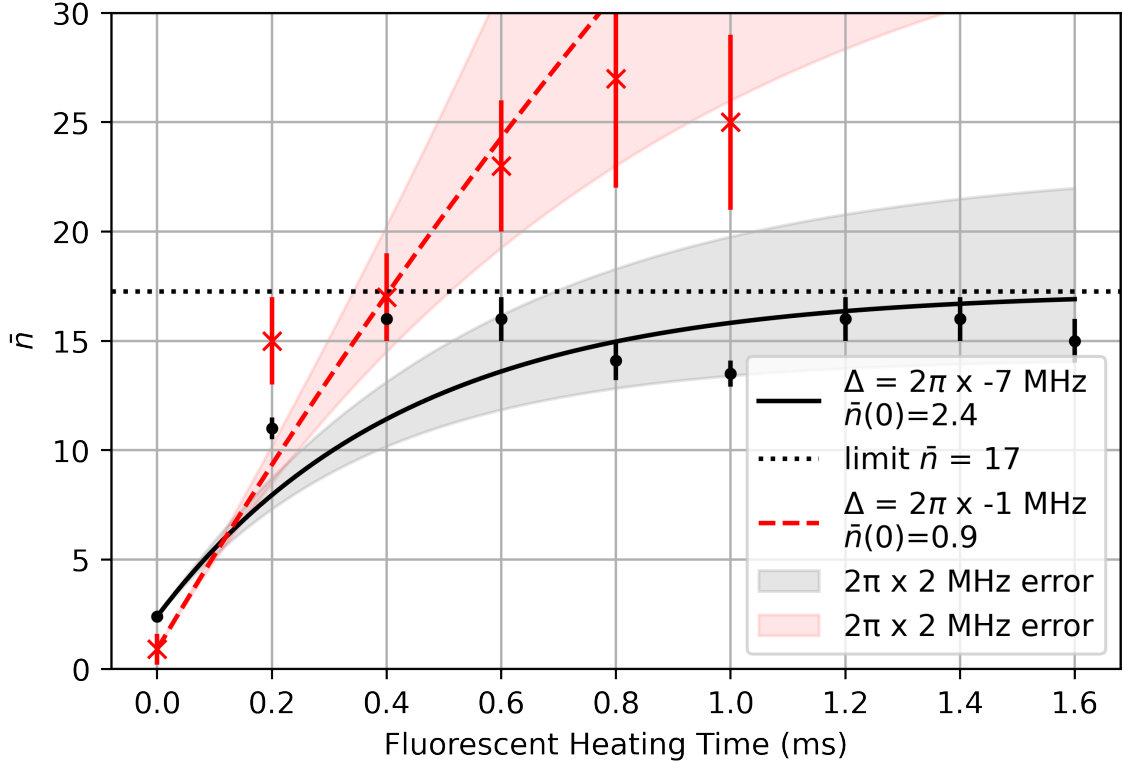


Figure 5.7. Measured average occupation number $\bar{n}$ from heating an mg crystal using the fluorescing lasers on the g -qubit and measuring the m -qubit via carrier readout (Section 5.2.1). Shaded regions represent uncertainty of laser detuning Δ for associated data set and theory curve.

5.2.5 Doppler Cooling. The rate equation describing the dynamics of

Doppler cooling is [153]

$$\frac{d}{dt}\bar{n} = -(A_- - A_+) + A_+, \quad (5.12)$$

where the rate coefficients are

$$A_{\pm} = \frac{\Omega^2}{\Gamma} \eta^2 [\cos^2(\theta) W(\Delta \mp \nu) + \alpha W(\Delta)], \quad (5.13)$$

where $W(\Delta) = 1/(4\Delta^2/\Gamma^2 + 1)$, Ω is the Rabi frequency, θ is the angle between the laser's direction of travel and the motional axis, ν is the secular frequency, Δ is the detuning of the cooling laser from resonance, Γ is the rate of spontaneous emission on the carrier, and α describes the average component of the recoil energy on the motional axis weighted

with the dipole pattern of spontaneous emission. Eq. (5.12) can be solved for the steady-state vibrational occupation number $\bar{n}$ as a function of time when $A_- > A_+$, *i.e.*, when the rate coefficient for red-sideband transitions is larger than the coefficient for blue sideband transitions. When $\Delta < 0$ ($\omega_L < \omega_0$), $A_- > A_+$ and we get a form similar to Eq. (5.11):

$$\bar{n}(t) = \bar{n}_0 e^{-(A_- - A_+)t} + \frac{A_+}{A_- + A_+} \left[1 - e^{-(A_- - A_+)t} \right], \quad (5.14)$$

where $\bar{n}_0$ is the initial vibrational occupation number and $\bar{n}_{lim,Doppler} = \frac{A_+}{A_- + A_+}$ is the steady-state value at long times.

When we set $\Delta = -20(2)$ MHz, the corresponding theoretical Doppler cooling limit is $\bar{n} = 5.8(3)$. Results from measuring $\bar{n}$ using carrier readout with variable Doppler cooling times are plotted and fitted in Fig. 5.8. The fit gives a Doppler cooling limit of $\bar{n}_{lim,Doppler} = 5.6(9)$, which agrees with our theoretical Doppler cooling limit.

5.2.6 Electromagnetically Induced Transparency Cooling. The EIT cooling limit can be calculated by solving a rate equation derived from the three-level master equation [103, 104]. The populations $P(n)$ of the vibrational number states n can be calculated as

$$\frac{d}{dt}P(n) = \eta^2 \{ A_- [(n+1)P(n+1) - nP(n)] + A_+ [nP(n-1) - (n+1)P(n)] \} \quad (5.15)$$

where

$$A_{\pm} = \frac{\Omega_g^2}{\gamma} \frac{\gamma^2 v^2}{\gamma^2 v^2 + 4 [\Omega_r^2/4 - v(v \mp \Delta)]^2}, \quad (5.16)$$

$\Omega - g(\Omega_r)$ is the Rabi frequency of the cooling (coupling) laser, γ is the linewidth of the internal transition between the ground and excited states, v is the secular frequency, and the detuning of both the cooling and coupling lasers has been set such that $\Delta = \Delta_g = \Delta_r$. Solving Eq. (5.15) for the steady state solution of the mean vibrational quantum number

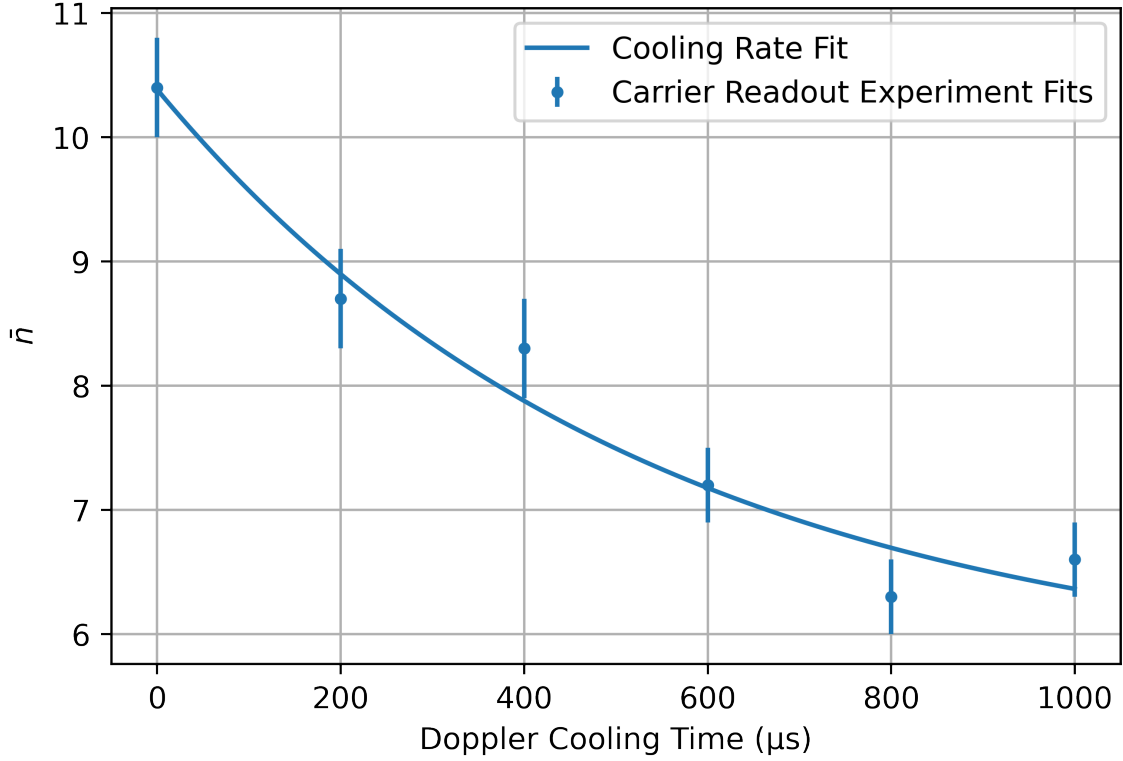


Figure 5.8. *mg* crystal Doppler cooled for variable time fitted to Eq. (5.14). Readout was performed using the 976 nm carrier transition. The fit gives a Doppler cooling limit $\bar{n}_{lim,Doppler} = 5.6(9)$.

$\bar{n} = \langle a^\dagger a \rangle = \sum n P(n)$ gives

$$\bar{n} = \frac{A_+}{A_- - A_+} = \frac{\gamma^2 v^2 + 4[\Omega_r^2/4 - v(v + \Delta)]^2}{4\Delta v(\Omega_r^2 - 4v^2)}. \quad (5.17)$$

For our setup, our σ beam has detuning $\Delta = 60(2)$ MHz and Rabi frequency $\Omega_\sigma = 31(3)$ MHz, which gives a theoretical EIT cooling steady state of $\bar{n} = 0.15^{+0.11}_{-0.08}$.

Temperature measurements were taken after preparing a coherent state corresponding to $\alpha \sim 1$ followed by a variable amount of EIT cooling time. We fit the results to an exponential decay

$$\bar{n}(t) = (\bar{n}_0 - \bar{n}_{lim,EIT}) e^{-Rt} + \bar{n}_{lim,EIT} \quad (5.18)$$

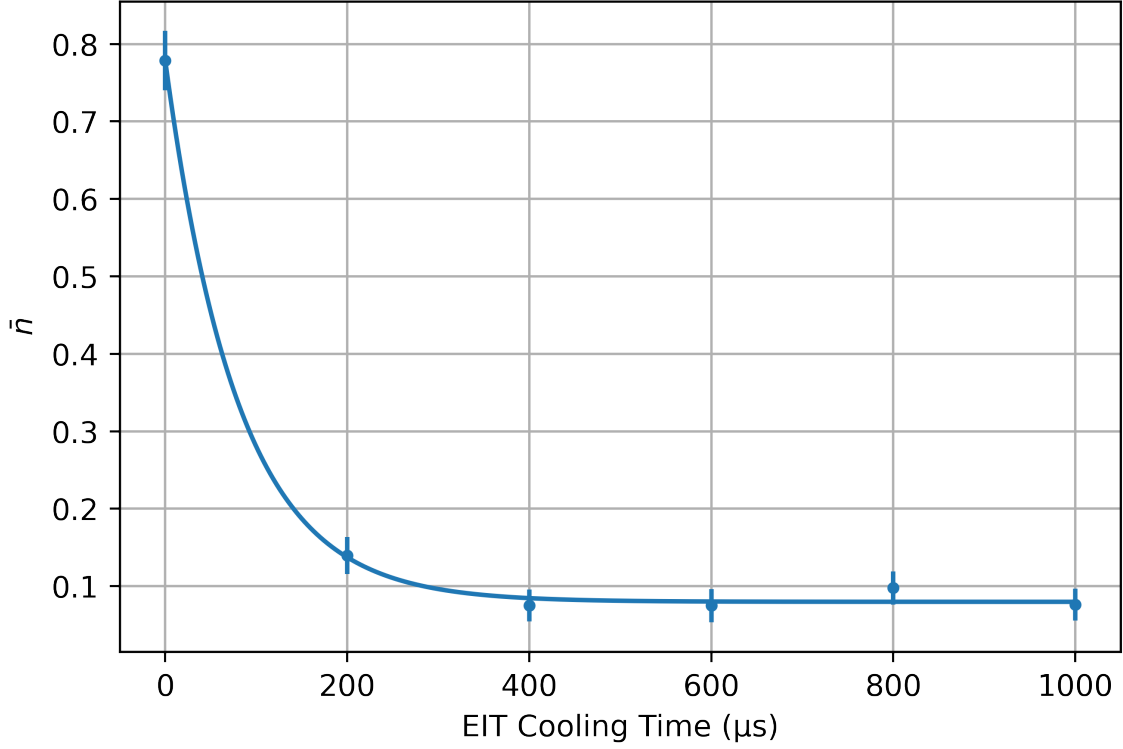


Figure 5.9. *mg* crystal EIT cooled for variable time fitted to exponential decay. The initial state was a coherent state corresponding to $\alpha \sim 1$. The fitted cooling limit is $\bar{n} = 0.080(6)$.

where $\bar{n}_0$ is the average occupational number before EIT cooling, R is the cooling rate, and $\bar{n}_{lim,EIT}$ is the EIT cooling limit. The data in Fig. 5.9 has a fitted temperature limit of $\bar{n} = 0.080(6)$, which agrees with our theoretical steady state.

5.3 *m*-Qubit Compatibility

To ensure high-fidelity *m-m* gates [34] are still possible in the full *omg* architecture, we tested the impact of various *g*-qubit operations on the lifetime of $|D_{5/2}, m = +5/2\rangle$, the coherence times of *m*-qubits, and a protected subspace of two *m*-qubits [154]. These operations include fluorescence checks (Section 4.1.2), Doppler cooling (Section 4.1.3.1), EIT cooling (Section 4.1.3.2), and quenching with the 854 nm lightshift as would be done during 729 or 976 nm sideband cooling. In the following sections, it will be shown that sideband cooling using the 729 or 976 nm lasers interferes

with coherent m -qubit operations as the quench limits the lifetime of $|D_{5/2}, m = +5/2\rangle$ and imparts noisy differential light shifts on the m -qubit.

5.3.1 Lifetime Measurements. The *omg* architecture relies on the long-lived natural lifetime of the metastable manifold to preserve the m -qubit [27]. Here we measure the effect our laser beams have on the lifetime to ensure the m -qubit does not decay too quickly. Various laser beams were directed at the ion for a time t and population was read out via fluorescence detection (Section 4.1.2). The results are plotted in Fig. 5.10 with fits for the population in $D_{5/2}$ $p = e^{-t/\tau}$ as a function of delay time and the lifetime τ . The 854 nm lightshift beam for the 729 nm sideband cooling quench is the only one that significantly shortens the lifetime of the m -qubit.

Although the 854 nm σ_+ laser beam is only meant to deshelve the $D_{5/2}$ manifold for $m \leq +1/2$ to complete the sideband cooling process, it has residual π and σ_- components that deshelve the qubit states as well. We also note that the beam we applied is -20 GHz detuned from resonance to minimize this scattering. All other results are within error of the lifetime of the $D_{5/2}$ manifold measured in the ‘no lasers’ experiment. The measured lifetimes exceed the natural lifetime (1.168(9) s [102]) due to the 393 nm beam not being fully extinguished, which causes scattering from $S_{1/2}$ to $D_{5/2}$. This prevents decay events from being detected and biases the lifetime measurement [31].

5.3.2 Coherence Times. For quantum computing applications, the coherence time of qubits must be longer than the operations performed to manipulate quantum states [155]. Each qubit state experiences AC Stark shifts from laser beams, which fluctuate due to intensity fluctuations in the beams. The m -qubit should not experience significant differential shifts, which would result in dephasing [156], from the fluorescence detection and cooling beams because they are 100s of THz detuned from any resonance and are much lower power than the 854 nm light shifting beam. Using detuning

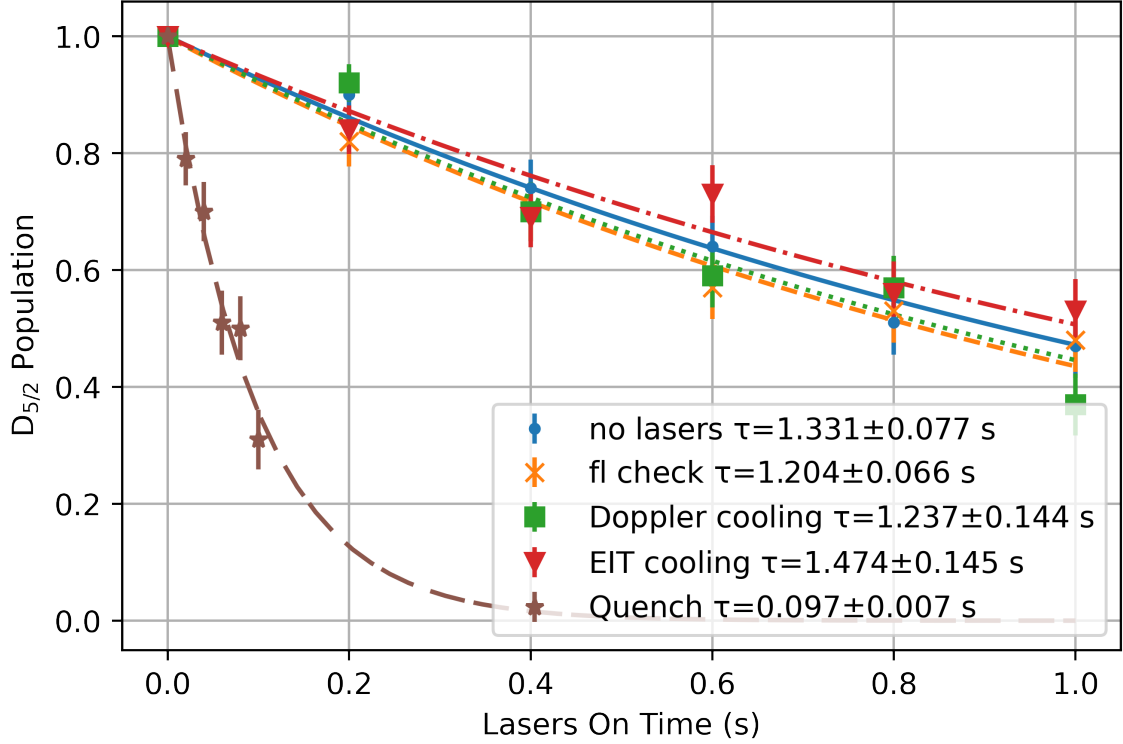


Figure 5.10. Lifetime measurements of the $|0\rangle$ state encoded in a single ion with a natural lifetime of $1.168(9)$ s [102] during various g -qubit operations.

and intensity of the beams to make an order-of-magnitude estimate of the Stark shifts on the $P_{3/2} \leftrightarrow D_{5/2}$ transition, the 866 and 397 nm beams produce Stark shifts less than a hundredth and a millionth the size of the Stark shift from the 854 nm light shift laser beam, respectively. This suggests operations with these laser beams impact on coherence times are orders of magnitude weaker than operations using the 854 nm beam.

m -qubit Ramsey experiments were performed during various operations on the g -qubit to check for excess dephasing. After various delay times t , the phase of a second $\pi/2$ -pulse was varied over one turn and fitted to a sinusoid with an amplitude that is referred to as the ‘Ramsey contrast.’ The Ramsey contrast versus delay time was then fitted to a Gaussian

$$f(t) = ae^{-\left(\frac{t}{\rho}\right)^2} \quad (5.19)$$

where t is the Ramsey delay time, a is the Ramsey contrast at $t = 0$, and ρ is the coherence time. Fig. 5.11a shows that the various g -qubit operations do not significantly effect the coherence time. Three out of the four ‘various lasers on’ experiments are within error of the ‘no lasers’ experiment, although this measurement is only sensitive to ~ 300 Hz-level shift fluctuations due to our baseline qubit dephasing. The pulse sequence for the Ramsey experiments is shown in Fig. 5.11b.

5.3.3 m - m Protected Subspace. Suppose an operation imparts a relative phase ϕ between the two states such that $|0\rangle + |1\rangle \rightarrow |0\rangle + e^{i\phi} |1\rangle$. Then the state $\Phi_+ = \frac{|00\rangle + |11\rangle}{\sqrt{2}}$ receives a phase such that $|00\rangle + |11\rangle \rightarrow |00\rangle + e^{2i\phi} |11\rangle$ whereas $\Psi_+ = \frac{|01\rangle + |10\rangle}{\sqrt{2}}$ receives a phase such that $|01\rangle + |10\rangle \rightarrow e^{i\phi} |01\rangle + e^{i\phi} |10\rangle = e^{i\phi} (|01\rangle + |10\rangle)$, which is a global phase that does not effect measurement outcomes. The entangled state Ψ_+ therefore represents a protected subspace from correlated phase noise, *i.e.*, it is decoherence free under this noise model. In fact, decoherence free subspaces exist within fully separable states as well [154]. For two ions exposed to the same laser beams, this creates correlated phase noise on the Stark shifts of each ion which can limit the coherence of our qubits. Our ions are close to each other ($\sim 7 \mu\text{m}$ apart) compared to our beam waists (25–45 μm radii for fluorescence detection and cooling laser beams), so they are likely to experience similar intensity fluctuations, *i.e.*, correlated noise that only imparts a global phase.

To determine if the laser operations we use break the protection of this subspace by introducing uncorrelated noise, we prepare two ions in $|0\rangle$ and use an RF $\pi/2$ -pulse to make an equal superposition of each ion in $|0\rangle$ and $|1\rangle$ such that

$$\Psi_p = \frac{|0\rangle + |1\rangle}{\sqrt{2}} \otimes \frac{|0\rangle + |1\rangle}{\sqrt{2}} = \frac{\Psi_+}{\sqrt{2}} + \frac{|0\rangle|0\rangle + |1\rangle|1\rangle}{2}. \quad (5.20)$$

We then perform a Ramsey experiment as described in [154] with delay time t . As t increases, the $|0\rangle|0\rangle$ and $|1\rangle|1\rangle$ coherences decay away as the relative phase from

correlated noise builds up while the global phase experienced by the state Ψ_+ maintains coherence.

After enough time has passed, the freely evolved mixed state of Ψ_p becomes

$$\rho_p = \frac{|\Psi_+\rangle\langle\Psi_+|}{2} + \frac{|00\rangle\langle 00| + |11\rangle\langle 11|}{4}, \quad (5.21)$$

with 50% probability of being in the Bell state Ψ_+ and 50% in one of the two separable states. Readout is performed by driving $|1\rangle$ population to $S_{1/2}$ via optical pumping with the 854 nm laser and performing detection via fluorescence check (see Section 4.1.2). This readout process is not the mg readout discussed earlier, but a readout specifically for m -qubit state detection [33, 34]. From the readout result, we can determine the parity, which is defined as $\Pi = P_{|00\rangle} + P_{|11\rangle} - P_{|01\rangle} - P_{|10\rangle}$. The parity measurement as a function of time goes like $\cos[\phi(t)]$. The contrast of the parity oscillates between -0.5 and 0.5 due to 50% of states ($|0\rangle|0\rangle$ and $|1\rangle|1\rangle$) not contributing to the signal since they are distributed evenly across the basis states of the two ions [154]. The first two data points of each plot in Fig. 5.12 are not fitted because the $|0\rangle|0\rangle$ and $|1\rangle|1\rangle$ states have not yet accumulated enough phase noise meaning the quantum state is not described by the mixed state of Eq. (5.21).

Fig. 5.12a plots parity measurements while both ions are exposed to various laser systems corresponding to operations on the absent g -qubit during the delay portion of the Ramsey experiment described above. See Fig. 5.12b for the pulse sequence. Decay of the $D_{5/2}$ manifold does not limit the measurements because a pre-readout fluorescence check determines if either of the ions has leaked out of $D_{5/2}$. For more information on leakage detection via fluorescence checks see [33, 34]. The fits in Fig. 5.12a are a sine function multiplied by an exponential decay with time constant ν . The protected subspace remains relatively undisturbed when no lasers are on, fitting to a contrast of $0.52(3)$ and $\nu = 1.2(3)$ s, much longer than our single-qubit coherence times. Fluorescence checking and

Doppler cooling show similar performance with fitted contrasts of 0.48(3) and 0.50(3) and $\nu = 2(1)$ and $3(2)$ s, respectively.

The EIT cooling trial fits to a contrast of 0.61(6) and $\nu = 0.5(1)$ and at longer times ($t \geq 0.3$ s) shows a bias towards positive values. The bias and smaller ν value can be attributed to leakage events being heated by the 397 and 866 nm lasers. For short times (1 ms cooling periods), the AOMs provide consistent control of the detuning and power of the laser beams. At longer times ($\gg 1$ ms) there are thermal effects on the AOMs which changes the alignment of the laser beam through the crystal and reduces the transmitted power [157]. EIT cooling relies on enhancing the red sideband absorption and suppressing the carrier and blue sideband absorption by matching the detuning of the intense coupling beam and a weaker cooling beam and matching the Stark shift induced by the coupling beam to the secular motion [103, 104]. With lower power beams due to thermalization at long times, the Stark shift induced by the coupling laser no longer matches the secular motion, carrier and blue sideband transitions are no longer suppressed, and, therefore, cooling becomes heating. The AOMs used for EIT cooling are the same for Doppler cooling and readout. Thermalization from repeated use in experiments is not a concern because the duty cycle of the AOMs is so low due primarily to the data processing time and other operations. In general, the Doppler and EIT cooling laser beams are not applied for long times compared to experiment times in practice because they can be calibrated to have large cooling rates [104] and approach steady-state limits quickly (see Sections 5.2.5 and 5.2.6). Refitting the EIT cooling lasers data to exclude times longer than 0.2 s, where the bias becomes obvious (dashed line), gives an amplitude of 0.53(4) and $\nu = 0.8(5)$.

The ‘Quench lasers’ label refers to the 854 nm lightshift laser that would be used to quench out of $D_{5/2}$ to $P_{3/2}$ after using the 729 or 976 nm lasers to sideband

cool. The 854 nm beam is the closest in frequency (detuned 20 GHz) to the resonant frequencies of transitions involving the qubit states, which is why it has the greatest effect on the protected subspace. The protected subspace immediately collapses with this laser showing that 729 or 976 nm sideband cooling would strongly impact our logic qubits during mid-circuit cooling even if the leakage shown in the lifetime measurement (Section 5.3.1) were overcome.

Given that Doppler and EIT cooling are unable to achieve ground state cooling and traditional forms of resolved sideband cooling are not compatible with our use of m -qubits, an additional cooling method must be employed to reach the motional ground state.

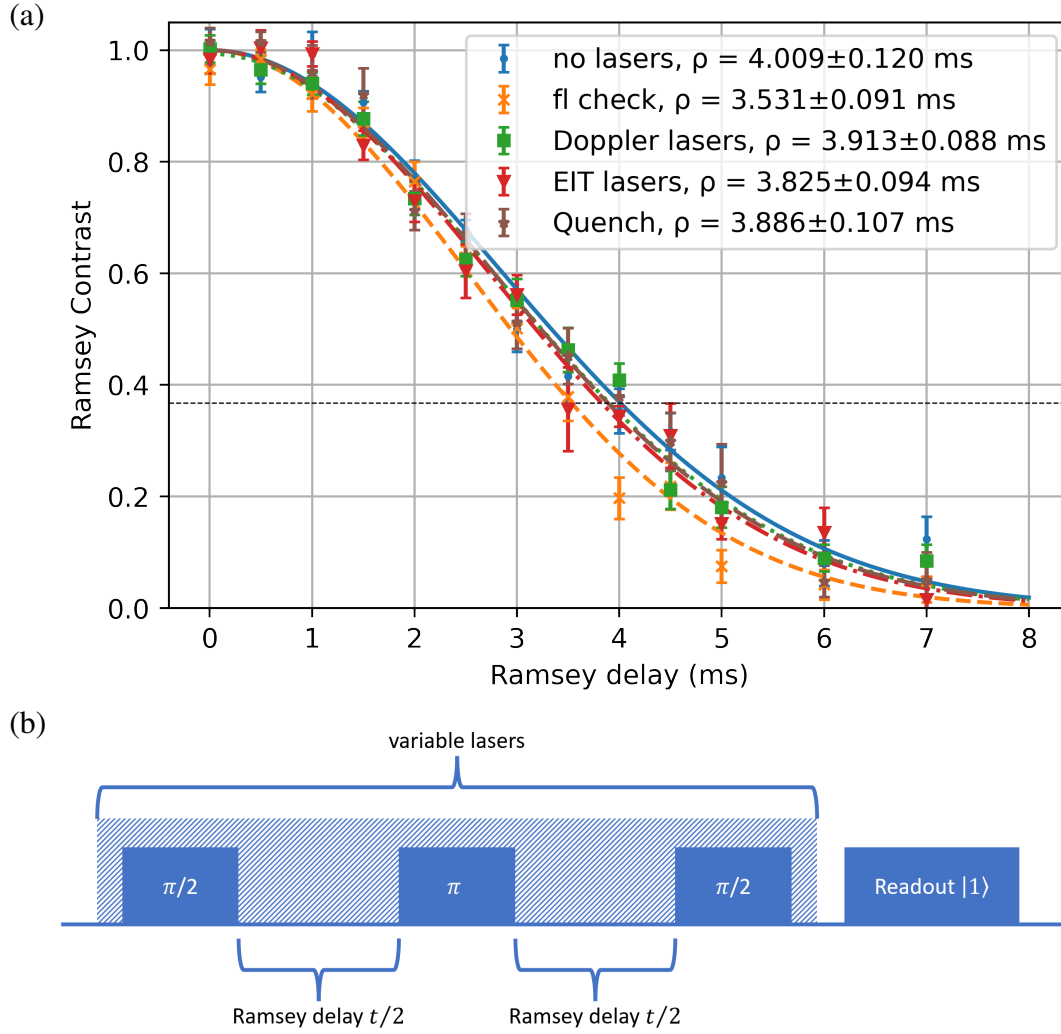


Figure 5.11. (a) Ramsey contrast plotted against delay time and fitted to a Gaussian distribution. The fits' intersection with the dashed line is where the Ramsey contrast equals one over Euler's number e , *i.e.*, when $t = \tau$. *(b)* Pulse sequence for Ramsey spin-echo experiments. π and $\pi/2$ -pulses are performed using RF. 'Variable lasers' are specified in the plot legend and occur at relevant conditions for each process.

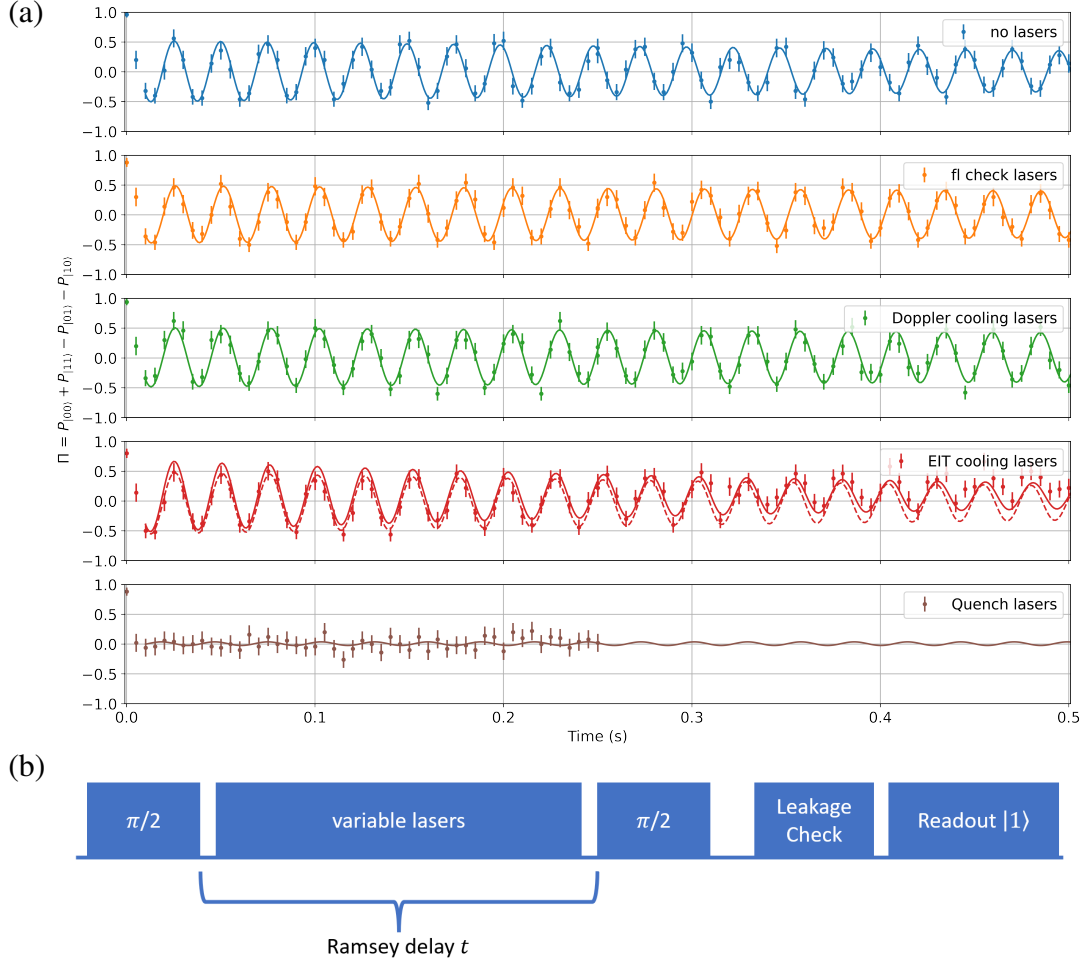


Figure 5.12. **(a)** Parity Π measurements of the protected subspace for two m qubits with various laser systems turned on during a Ramsey sequence. Lines are fits of the data to a sinusoid multiplied with an exponential decay. The dashed line for EIT cooling lasers excludes data for times $t > 0.2$ s. **(b)** Pulse sequence for the m - m protected subspace Ramsey experiments. $\pi/2$ -pulses are performed using RF. ‘Variable lasers’ are specified in plot legends and occur at relevant conditions for each process. To extend the experiment beyond the natural lifetime of the $D_{5/2}$ manifold of 1.168(9) s [102], a leakage detection [33, 34] is performed after the Ramsey sequence but before the readout procedure.

CHAPTER VI

PROJECTION COOLING

The work contained in this chapter was performed by myself and Dr. Jameson O'Reilly with writing done by myself with editorial assistance from Dr. O'Reilly. Connor Burns assisted with projection cooling simulations.

One application for the *omg* architecture in quantum computing is making non-destructive mid-circuit measurements [158], *e.g.*, using a *g*-qubit to determine the state of an *m*-qubit without disturbing it. Mid-circuit measurement techniques and uses are being explored [159] for applications such as error correction [29, 30] and efficient qubit use [160] and are possible in commercially available systems [14]. Both of these applications, and more, are important for the development beyond the Noisy Intermediate Scale Quantum (NISQ) [5] era.

Cooling the ions' motion close to the ground state is important to reduce temperature effects that limit the fidelity of qubit operations [50, 138] including mid-circuit measurements for determining qubit states. While a combination of Doppler and EIT cooling can reduce the occupation number of the trapped ions in the Fock basis, together they are limited to cooling above the motional ground state (see Sections 5.2.5 and 5.2.6). The typical solution for cooling close to the motional ground state ($\bar{n} \ll 1$) is via sideband cooling [51, 161], which we can perform using 729 or 976 nm lasers. Sideband cooling relies on using the phonon-subtracting sideband to drive a two level system where spontaneous emission occurs out of the excited state. Because the electronic ground state is driven to the excited state with the phonon-subtracting sideband, the energy of the photon absorbed by the atom is less than the energy of the spontaneously emitted photon resulting in a net loss of energy. In our system, we can either drive the 729 nm sideband from $S_{1/2}$ to $D_{5/2}$ or we can prepare in the $D_{5/2}$

manifold and use a 976 nm sideband. The 854 nm laser then drives the $D_{5/2} \leftrightarrow P_{3/2}$ transition (the “quench”) and population decays from $P_{3/2}$ back to $S_{1/2}$ by spontaneous emission of a photon. Further explanation of resolved sideband cooling in our system can be found in [32, 33]. Another route for sideband cooling would be using 397 nm Raman beams, however, this would require adding a set of beams that would not be used for anything else.

Sideband cooling via the 729 or 976 nm lasers disturbs the protected subspace of our m -qubits as described in Section 5.3.3 due to the 854 nm laser beam that provides the quench process out of $D_{5/2}$ imparting strong Stark shifts on the qubit subspace and, therefore, cannot be used for mid-circuit cooling. Sideband cooling using 397 nm Raman beams would not disturb the m -qubit subspace as $|\uparrow\rangle$ is used as the excited state before quenching through optical pumping to $P_{1/2}$ back to $|\downarrow\rangle$. But all sideband cooling shares another limitation. Population decaying from $P_{1/2}$ back to $S_{1/2}$ occurs through the scattering of a photon, which imparts momentum on the ion. Sideband cooling will be limited by the amount of energy associated with the momentum kick of this scattered photon [45].

Here, we explore an alternative to sideband cooling for our mg crystals: by instead detecting non-ground state motional population, we can herald the motional ground state and re-cool after failed attempts. This process will be referred to as ‘projection cooling’ [162, 163]. Neutral systems also utilize this process under different names [164]. In this chapter, we herald motional ground states in mg crystals as well as perform ancilla measurements, which rely on ground state cooling for high fidelity results.

6.1 Heralding Motional Ground States

Before projection cooling, we prepare two ions in the mg crystal configuration. After heralding, we Doppler and EIT cool followed by pumping the g -qubit

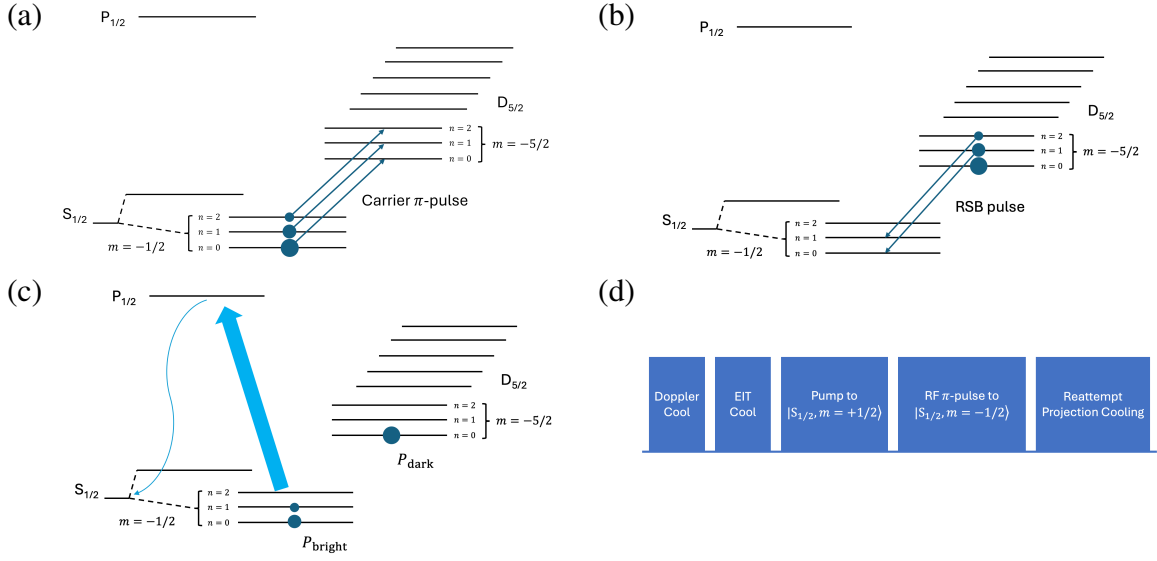


Figure 6.1. Diagram of projection cooling. Relative motional populations are represented by size of the circles corresponding to each Fock state. **(a)** The g -qubit is Doppler and EIT cooled, prepared in $|S_{1/2}, m = -1/2\rangle$, and then is shelved with a 729 nm carrier π -pulse to $|D_{5/2}, m = -5/2\rangle$. **(b)** A 729 nm RSB pulse on the first phonon-subtracting sideband (RSB) drives population in motional states $n > 0$ back to $|S_{1/2}, m = -1/2\rangle$. **(c)** Fluorescence detection is performed with the probability of detecting a bright ion $P_{\text{bright}} = 1 - P_{\text{dark}}$ where the probability of not detecting a bright ion P_{dark} is the probability of occupying the ground state after reaching the EIT cooling limit. **(d)** Pulse sequence for resetting the g -qubit after detecting a bright ion.

to $|S_{1/2}, m = +1/2\rangle$. An RF π -pulse then transfers the g -qubit population to $|S_{1/2}, m = -1/2\rangle$. The population is transferred to $|D_{5/2}, m = -5/2\rangle$ with a 729 nm carrier π -pulse. Next, we attempt to drive population back to $|S_{1/2}, m = -1/2\rangle$ via a 729 nm phonon-subtracting pulse. We will refer to this pulse as the red sideband (RSB) to be consistent with the phonon adding/subtracting nomenclature used in Chapter V even though it is technically blue-shifted from the carrier frequency. If the mg crystal's motional mode occupies $n = 0$, the RSB pulse will not be able to drive population back to $S_{1/2}$ as there is no phonon to subtract. To ensure all $n > 0$ population is transferred, bucket deshelling (Section 5.1.2.1) with $D_{3/2}$ is employed to allow multiple RSB pulses in succession. A fluorescence check is then performed to determine if the ion is in $S_{1/2}$. If the ion is bright, then it was not in $n = 0$ and Doppler and EIT cooling is performed

followed by a repeat of this projection cooling process. This is designated as a ‘rejected shot’ (Section 6.1.1). It is important to note that use of the phrase ‘rejected shot’ here is not post-selection of data, but rather a pre-selection of prepared states. Post-selection requires discarding experiments after they are performed and increases run-times by requiring a larger number of overall experiments. Pre-selection with projection cooling prevents experiments from continuing during the preparation stage, which reduces run times by not performing experiments that would have failed or been discarded after completion. If the ion is dark, we assume it occupies $n = 0$ and is ready for an experiment requiring motional ground state cooling of the addressed mode. The projection cooling pulse sequence on the g -qubit ion is depicted in Fig. 6.1. The m -qubit is not depicted as projection cooling is designed to prevent disturbing it by using $|D_{5/2}, m = -5/2\rangle$ to shelve the g -qubit and maintain as large a separation from the m -qubit population as possible.

6.1.1 Passing Fluorescence Checks. In a thermal state, the probability of occupying the Fock state $|n\rangle$ is given by

$$p_n = \frac{1}{(\bar{n} + 1)} \left(\frac{\bar{n}}{\bar{n} + 1} \right)^n, \quad (6.1)$$

so as the temperature of the ion increases, the probability of being in the motional ground state decreases and the probability of passing the projection cooling fluorescence check goes to zero. If the initial cooling of the ions is done only with Doppler cooling, the majority of heralding attempts will result in failure leading to long preparation times. In [162], heralding of the motional ground state was performed after cooling their ions to the Doppler limit of $\bar{n} \sim 18$, which resulted in heralding the motional ground state with a probability $p_0(\bar{n} = 18) = 5\%$. For our Doppler cooling limit $\bar{n} = 5.6$, the probability of passing the fluorescence detection is limited to 15.2%. To improve the probability of heralding the motional ground state, we include EIT cooling to get $\bar{n} \sim 0.08$, where the

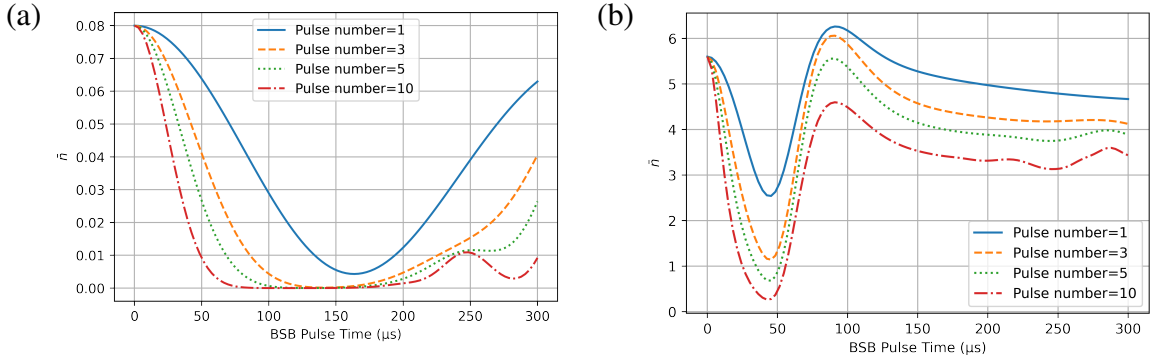


Figure 6.2. Theory model for final temperature after projection cooling with variable BSB pulse times and numbers at starting temperatures of **(a)** the EIT cooling limit of $\bar{n} = 0.08$ and **(b)** the Doppler cooling limit of $\bar{n} = 5.6$.

probability of success is limited to 92.6%, which greatly increases the speed at which the motional ground state is prepared.

Due to uncertainty in the RSB pulse time calibration and the n -dependence of the Rabi frequencies, a single RSB pulse cannot perfectly transfer non-zero motional population. To mitigate this, the work done in [162] used rapid adiabatic passage (RAP) to ensure complete population transfer of all n -states. RAP is insensitive to the coupling strength used to invert population of two internal states and from errors due to control parameters but requires more complicated frequency control and is more likely to off-resonantly couple to unwanted transitions. To simplify the population transfer process, we use multiple RSB pulses separated by dissipative resets (bucket deshelling with $D_{3/2}$, see Section 5.1.2.1) to approach the probability limit imposed by the initial $\bar{n}$. Calibration of the optimal RSB pulse time, explained in the following section, also improves non-zero n state population transfer.

6.1.2 Temperature Limits. To simulate projection cooling, we constructed a semi-classical model of population transfer between the $S_{1/2}$ and $D_{5/2}$ manifolds in the Fock state occupation basis. In this model we include probability of driving population with the 729 nm RSB pulses for various pulse times and number of pulses,

the n -dependence of the Rabi frequency for driving population, and the fidelity of bucket shelving to $D_{3/2}$. Fig. 6.2 plots the final temperature after passing the fluorescence check in projection cooling against RSB pulse time for several numbers of pulses. The initial temperatures are at the EIT and Doppler cooling limits. Sources of error not included in this simplified model include laser power and frequency fluctuations, off-resonant driving, and π -pulse time calibration uncertainty. However, the model captures several important features of this cooling technique.¹ The lower the initial temperature, the wider range of RSB pulse times that reach the minimum temperature. A greater number of pulses also reduces the minimum temperature with diminishing returns as there is a physical limit to the amount of population to be transferred. Interestingly, if the RSB pulse time is too long, heating can occur. As seen in Eq. (5.5), the n -dependence of the Rabi frequency means larger n states have shorter π -times. When the 729 nm RSB pulses are calibrated for the $n = 1$ state, larger states are driven back up to the $D_{5/2}$ manifold. For small $\bar{n}$ values, this effect is small since the probability of occupying $n > 1$ is small. But as $\bar{n}$ increases, the larger n states have higher occupation probability and the optimum 729 nm RSB pulse time decreases. If for larger initial $\bar{n}$ states the 729 nm RSB pulse time doubles the π -time of the most likely occupied n state, the lower occupied states are pre-selected out and the hotter states survive.

6.2 Minimum Temperature Measurements of mg Crystals

The phonon-subtracting sideband driven by the 976 nm Raman beams (referred to as the red-sideband (RSB) from here on) can be used to calibrate the optimal pulse length for the 729 nm RSB in projection cooling. This data is useful for calibration of the 729 nm RSB pulse time because it can be taken at a faster rate than 976 nm BSB readout from Section 5.2.1. After preparing the m -qubit in $|0\rangle$, the 976 nm RSB will only drive

¹“In this house, we obey the laws of thermodynamics!” — Homer Simpson, *The Simpsons* S6E21 “The PTA Disbands”

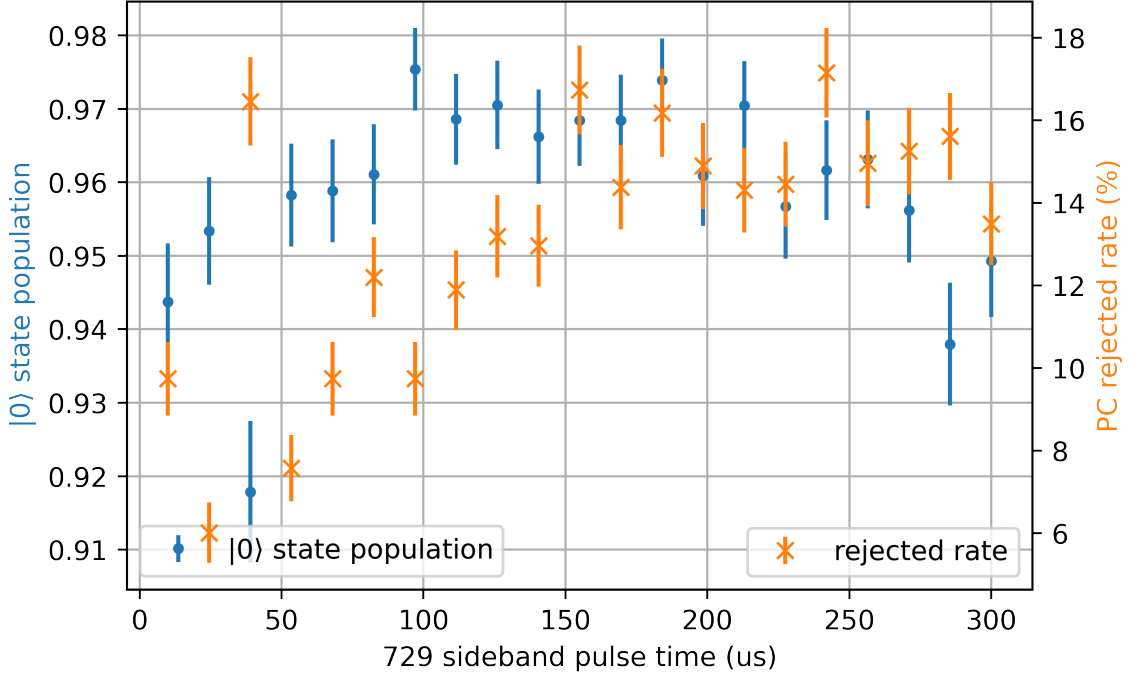


Figure 6.3. Experimental results for 976 nm RSB driving calibration of projection cooling 729 nm BSB pulse time using three pulses with initial $\bar{n}$ at the EIT cooling limit.

population into $|1\rangle$ if the motional state is in $n > 0$. The experiment in Fig. 6.3 used three 729 nm RSB pulses to projection cool. At a 729 nm RSB pulse time of zero, projection cooling will always pass the fluorescence check since no shelved g -qubit population will be driven to the ground state, so there is no change in temperature. Initially as the pulse time increases, more $n \neq 0$ population is driven to ground to pre-select out more hot states. Then the pulse time hits an optimum value for cooling where $|0\rangle$ population is minimally driven by the 976 nm RSB Raman beams and final $|0\rangle$ population is at a maximum. For lower initial $\bar{n}$, the optimum value approaches the π -time for the first RSB beam, which was calibrated to be 175 μs . The maximum $|0\rangle$ population in Fig. 6.3 occurs around 175 μs as well as in Fig. 6.2a which was calculated using a RSB π -time of 175 μs . As $\bar{n}$ approaches zero, the majority of thermal population not occupying $n = 0$ occupies $n = 1$. This makes the optimal time for the 729 nm RSB pulse equal the π -time. As $\bar{n}$

increases, higher Fock states become occupied and the RSB π -time becomes less effective at driving non-zero population. In Fig. 6.2b, the initial $\bar{n} = 5.6$ where the majority of the population occupies states $n > 1$ and the optimum RSB pulse time drops below 50 μs . For large $\bar{n}$, the maximum $|0\rangle$ population is limited by the final temperature after projection cooling. As $\bar{n}$ approaches zero, the maximum $|0\rangle$ population is limited by the mg crystal SPAM error for readout of the m -qubit (Section 5.1.3.1), which has not been subtracted from the plot in Fig. 6.3.

After optimizing the pulse timing using 976 nm RSB calibrations, we achieve a minimum temperature from 976 nm BSB readout of a projection cooled mg crystal of $\bar{n} = 0.02(1)$ with three cooling pulses, approximately a factor of four improvement in cooling compared to our EIT cooling limit (Section 5.2.6).

6.3 m -Qubit Compatibility

We performed experiments to determine the effects of the 729 nm laser beam used for projection cooling on our m -qubit in the same fashion as the experiments in Section 5.3 performed for fluorescence detection and Doppler, EIT, and sideband cooling. Fig. 6.4a and b show that the 729 nm laser beam on the RSB pulse frequency has no noticeable effect on a single m -qubit. In Fig. 6.4d, the parity for the protected subspace does not remain undisturbed for as long as the experiments for fluorescence check, Doppler cooling, and EIT cooling, but it still fits to a parity contrast of 0.55(5) and $\nu = 0.22(3)$ s. The 729 nm laser beam produces an estimated Stark shift on the $P_{3/2} \leftrightarrow D_{5/2}$ transition seven times as large as the 866 nm laser beam and 100,000 times as large as the 397 nm laser beam. It is still an order-of-magnitude smaller than the 854 nm light shift beam. In practice, the 729 nm RSB pulse times during projection cooling are on the order of 100 μs , a time-scale where the subspace of the m -qubit remains undisturbed.

6.4 Non-Destructive Mid-Circuit Measurements

Ground-state cooling via projection cooling enables us to perform quantum logic spectroscopy (QLS)-like [165] ancilla readout of the m -qubit using the g -qubit. The logic qubit's internal state is mapped to the ancilla qubit to allow for high efficiency readout of the logic qubit state non-destructively.

Here we present the procedure and results of using the optical (o) qubit to non-destructively measure the m -qubit population in an mg crystal. The o -qubit in $^{40}\text{Ca}^+$ refers to a two level system with one state in the $S_{1/2}$ manifold and one state in the $D_{5/2}$ manifold and driven with the narrow-linewidth 729 nm quadrupole laser (Section 4.4). The ion that is designated the g -qubit in heralding the mg crystal and projection cooling can be referred to as the o -qubit and ancilla qubit interchangeably at this point and the m -qubit is considered the logic qubit. For this procedure, the o -qubit will be $|S_{1/2}, m = -1/2\rangle \leftrightarrow |D_{5/2}, m = -5/2\rangle$.

Projection cooling heralds the mg crystal in the motional ground state with high probability, which means any non-zero motional occupation must be due to a phonon-adding process, specifically, a 729 nm sideband pulse. After successful projection cooling, the g -qubit is heralded to be in $|D_{5/2}, m = -5/2\rangle$. If projection cooling is not being used, a 729 nm π -pulse shelves the g -qubit in $|D_{5/2}, m = -5/2\rangle$. The following procedure is for non-destructively measuring the m -qubit:

1. π -pulse on $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$ drives $|0\rangle$ population to the ground manifold.
2. π -pulse on the first phonon-adding sideband of $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$ drives population back to $|0\rangle$, but now with all population in $n = 1$.

3. π -pulse on the first phonon subtracting sideband of $|S_{1/2}, m = -1/2\rangle \leftrightarrow |D_{5/2}, m = -5/2\rangle$ only drives shelved g -qubit population that is not in the motional ground state.
4. An FC is made.
5. Bucket shelve detected population to $D_{3/2}$.

If the m -qubit is $|0\rangle$, then the carrier and phonon-adding sideband pulses will have added a phonon to the shelved g -qubit, the phonon-subtracting sideband will have driven the g -qubit back to the $S_{1/2}$ manifold, and the FC will result in a bright ion detection. If the m -qubit is $|1\rangle$, none of the three π -pulses will have interacted with any population and the FC will result in no ion detection. At this point, the quantum state of the m -qubit has not been destroyed and we have measured it indirectly. To verify the non-destructive measurement is correct, we can now implement two FCs similar to the second and third FCs in the mg crystal readout scheme (Section 5.1.2):

6. π -pulse on $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +3/2\rangle$ moves any population in $|1\rangle$ to the $S_{1/2}$ manifold.
7. Bucket deshelve for a second π -pulse on $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +3/2\rangle$ to deshelve any remaining shelved population. See Section 5.1.2.1.
8. A second FC is made.
9. Bucket shelve all detected population to $D_{3/2}$.
10. π -pulse on $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$ moves any population in $|1\rangle$ to the $S_{1/2}$ manifold.

Table 6.1. Corresponding state for given sequence of ions detected in o -qubit ancilla readout scheme for non-destructive measurement of m -qubit verification. ‘Ancilla failure’ refers to the non-destructive measurement not corresponding to the m -qubit direct measurement. ‘Leakage’ refers to the metastable population decaying, which prevents determination of the state populations. All other combinations of detections are categorized as a generic ‘error.’

Fl. Check #1	Fl. Check #2	Fl. Check #3	Result
0	0	1	ancilla failure
0	1	1	$ 1\rangle$
1	1	2	$ 0\rangle$
1	2	2	ancilla failure
0	0	0	leakage
1	1	1	leakage
2	2	2	leakage

11. Bucket deshelve for a second π -pulse on $|S_{1/2}, m = +1/2\rangle \leftrightarrow |D_{5/2}, m = +5/2\rangle$ to deshelve any remaining shelved population. See Section 5.1.2.1.

12. A third FC is made.

The second and third FCs verify the m -qubit state determined by the non-destructive measurement and flags any instances of leakages. We can now update the mg crystal readout table (Table 5.2) to include failure of the ancilla detection. If the first FC detects no bright ions, then the m -qubit is non-destructively measured to be in $|1\rangle$ and the second FC must also detect just one bright ion. If no ion is detected for the second FC, this means the first three π -pulses failed and we label this as ‘ancilla failure.’ Similarly, if the first FC detects one bright ion, then the second FC must also detect just one bright ion since m -qubit population should remain in $|0\rangle$. See Table 6.1 for the o -qubit ancilla readout scheme FC result sequences and their corresponding meaning.

6.4.1 Optical Qubit Ancilla Measurement Results. We performed ancilla readout measurements at a range of ion temperatures by cooling the mg crystal different amounts using projection cooling and EIT cooling. The more EIT cooling performed before preparation and readout, the better we were able to non-destructively measure the

m -qubit. Using projection cooling after EIT cooling for 1 ms gave even higher fidelities. The results from non-destructive ancilla readout of the m -qubit via the optical transition are listed in Table 6.2. Fidelity of the ancilla readout is calculated as the probability of the g -qubit matching the prepared m -qubit after subtracting out leakages and errors. The probability of the m -qubit being readout as flipped from the prepared state is also listed.

From Section 5.1.3, mg SPAM error on the m -qubit being flipped is on the order of $\sim 0.1\text{--}1\%$. In the ancilla readout, the m -qubit being flipped remains within a standard deviation of this range of predicted SPAM error and is consistent across different temperatures, suggesting this procedure for mid-circuit measurement is non-destructive regardless of ancilla fidelity. Ancilla readout of $|1\rangle$ population is limited by imperfect motional ground state cooling. For ancilla readout of $|0\rangle$ population, the fidelity is limited by the sideband π -pulses. These pulses are slower by an order of magnitude compared to carrier pulses and experience more decoherence. Additionally, the strong n -dependence of the sideband Rabi frequencies degrades the fidelity of these pulses even further when the ion crystal is hot. Bucket deshelling as used in the mg readout scheme (Section 5.1.2.1) is convenient because of the short time scale at which additional π -pulses can be used, whereas in this ancilla readout scheme additional π -pulses would significantly increase instances of leakages.

In addition to showing how projection cooling improves our ability to make ancilla measurements on the m -qubit, we are also able to perform spectroscopy measurements on the RF transition of the m -qubit and measure driving of the m -qubit with RF. The fits plotted with the data in Fig. 6.5 give excellent overlap between the m and g -qubit readouts.

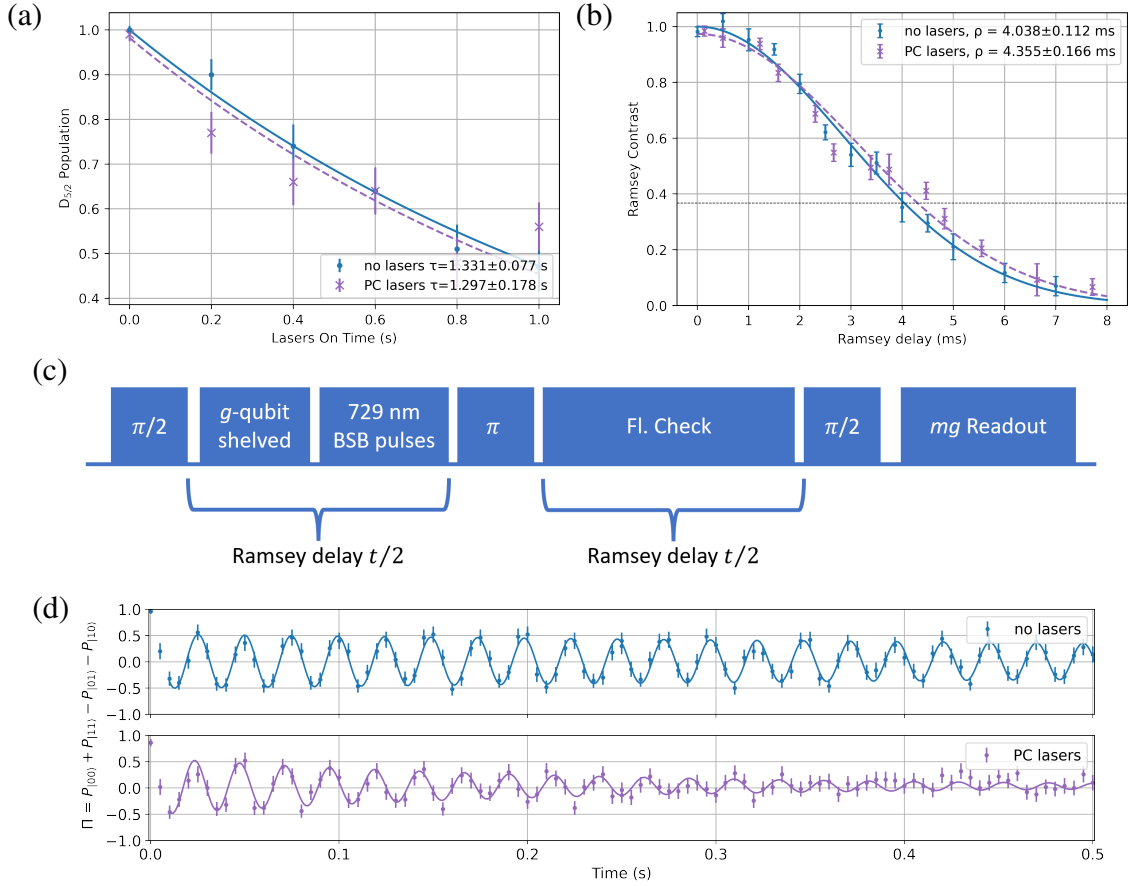


Figure 6.4. Experiments for testing projection cooling compatibility with m -qubits. **(a)** Lifetime measurements of a single ion prepared in $|0\rangle$ with a lifetime of $1.168(9)$ s [102] with the 729 nm laser on the RSB frequency for projection cooling turned on compared to no lasers turned on. **(b)** Ramsey contrast for the ‘no lasers’ and projection cooling experiments plotted against time and fitted to a Gaussian distribution (Eq. (5.19)). The fits’ intersection with the dashed line is where the Ramsey contrast equals one over Euler’s number e , *i.e.*, when $t = \rho$. **(c)** Pulse sequence for the projection cooling in the spin-echo experiment. The shelving pulse for the g -qubit is followed by some number of 729 nm RSB pulses in the first half of the spin-echo with a fluorescence check in the second half. The number of RSB pulses is a function of what the Ramsey delay time t is set to. **(d)** Parity Π measurements of the protected subspace for two m qubits with projection cooling turned on during a Ramsey sequence. The vertical axes of both the plots are from -1 to 1. Lines are fits of the data to a sinusoid multiplied with an exponential decay. Data and fit for the ‘no lasers’ experiment are the same as shown in Fig. 5.12, repeated here for convenience.

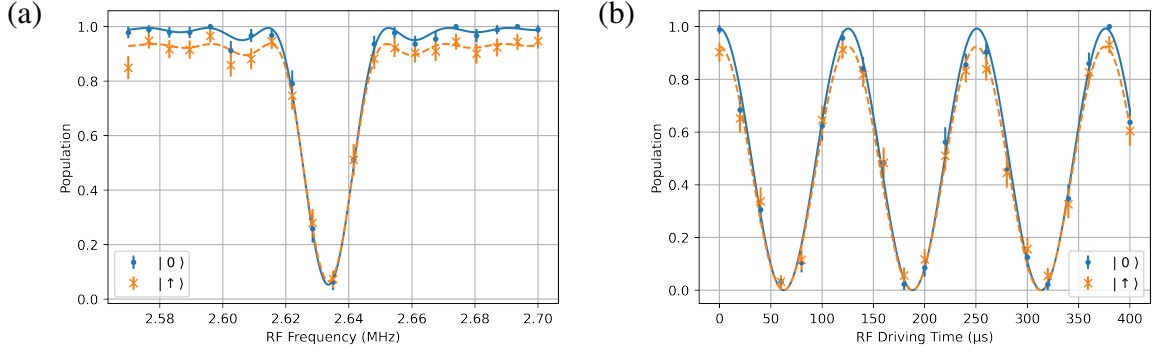


Figure 6.5. Comparison of the readout of the m qubit with readout via the o -qubit. **(a)** Scan of the RF driving frequency for the m -qubit fitted to a sinc-squared. The fit gives a center frequency of 2.6335(2) MHz for the m -qubit from readout of $|0\rangle$ population and 2.6336(3) MHz for the ancilla readout with $|\uparrow\rangle$ population. **(b)** Scan of RF driving pulse time on resonance of the m -qubit with RF fitted to a sine-squared. The fit gives a Rabi frequency of 250.9(6) kHz for driving $|0\rangle$ population and 250.5(7) kHz with ancilla readout of $|\uparrow\rangle$ population.

Table 6.2. Results of ancilla readout using the o -qubit to determine the m -qubit as described in Section 6.4.1. “Prep” refers to both the prepared state of the m -qubit as well as the cooling done. PC refers to projection cooling and is preceded by 1 ms of EIT cooling. Prepared states with “EIT” did not use projection cooling. Detection of $|\uparrow, 1\rangle$ and $|\downarrow, 0\rangle$ are considered “ancilla failure” as listed in Table 6.1. A BSB readout of an mg crystal was done at each of the cooling settings with the fitted $\bar{n}$ value listed in the rightmost column.

Result Prep	$ \uparrow, 0\rangle$ “Bright”	$ \uparrow, 1\rangle$	$ \downarrow, 0\rangle$	$ \downarrow, 1\rangle$ “Dark”	Leakage	Error	Ancilla Fidelity (%)	m Flip (%)	Measured $\bar{n}$
$ 0\rangle$, PC 5 pulses	861	6	67	6	57	3	92.2(8)	1.2(4)	0.04(2)
$ 1\rangle$, PC 5 pulses	3	22	0	936	27	12	97.4(5)	0.3(2)	
$ 0\rangle$, PC 0 pulses	836	1	56	3	102	2	93.4(8)	0.4(2)	0.13(3)
$ 1\rangle$, PC 0 pulses	6	93	1	871	26	3	90(1)	0.7(3)	
$ 0\rangle$, EIT 800 μ s	756	3	88	4	146	3	89(1)	0.8(3)	0.19(3)
$ 1\rangle$, EIT 800 μ s	3	107	1	862	24	4	89(1)	0.4(2)	
$ 0\rangle$, EIT 600 μ s	720	1	110	2	159	8	87(1)	0.4(2)	0.37(4)
$ 1\rangle$, EIT 600 μ s	1	197	0	769	24	9	80(1)	0.1(1)	
$ 0\rangle$, EIT 400 μ s	568	4	109	0	315	4	84(1)	0.6(2)	0.95(9)
$ 1\rangle$, EIT 400 μ s	6	298	2	635	48	11	68(1)	0.9(3)	
$ 0\rangle$, EIT 200 μ s	442	2	100	1	454	1	81(1)	0.6(2)	3.5(5)
$ 1\rangle$, EIT 200 μ s	2	453	4	483	55	3	52(2)	0.6(3)	
$ 0\rangle$, EIT 2 μ s	347	4	116	0	531	2	75(1)	0.9(3)	5(1)
$ 1\rangle$, EIT 2 μ s	0	469	0	431	98	2	48(2)	0(0)	

CHAPTER VII

SUMMARY AND OUTLOOK

This dissertation described work to design, construct, and characterize the Oregon Ions’ primary experimental apparatus, a rod trap, as well as the laser systems needed to trap, cool, prepare, and perform operations on $^{40}\text{Ca}^+$ ions in Chapters II and IV. Chapter III describes a project studying the extension of an alternative atomic experiment loading method found to be useful for alkaline metals in neutral traps, as well as strontium in ion traps, to calcium in ion traps. Finally, Chapters V and VI present work extending the utility of m -qubits within the *omg* architecture [27] through the heralding, control, and measurement of mg crystals as well as the demonstration of projection cooling for heralding motional ground states and its use in non-destructive ancilla measurements. These topics are ripe for continued exploration,¹ especially in the context of ongoing developments in the field of quantum sensing and information and the march through the NISQ era and beyond [5].

7.1 Projection Cooling Limits

In Section 6.1, the limits for projection cooling were investigated with a semi-classical numerical simulation followed by measurements. As already stated, the simulation did not include individual sources of error, such as laser frequency and power fluctuations, off-resonant driving, and π -pulse time calibration uncertainty. Work to include relevant error sources and uncover the true technical and non-technical minimum temperatures is ongoing.

The projection cooling scheme presented here specifically heralds the motional state of the co-trapped ions in $n = 0$ of the Fock number basis, but it could be adapted to

¹[Winged monkeys screaming and falling.] “Continue the research.” — Mr. Burns, *The Simpsons* S5E9 “The Last Temptation of Homer”

herald occupation of other states as well. Similar techniques have also been proposed as a way to probabilistically prepare rotational states of co-trapped molecular ions [166].

7.2 Mid-Circuit Measurement Improvements

Projection cooling was shown to improve mid-circuit measurements made in Section 6.4.1 when compared to cooling at the EIT limit. These measurements are primarily limited for two reasons. Preparation of the $|1\rangle$ state being readout by the ancilla qubit at 97.4(5)% fidelity is limited from imperfect ground state cooling, improvements of which are ongoing as mentioned above. Preparation of the $|0\rangle$ state being readout by the ancilla qubit at 92.2(8)% fidelity is limited by the sideband π -pulse fidelity, which are relatively slow compared to carrier pulses and experience more decoherence. Improving the deshelve fidelity is possible with bucket deshelve, but this would also increase instances of leakage.

Additional schemes for mid-circuit measurements with mg crystals are possible, including using the 976 nm Raman beams to create spin-dependent forces to drive closed loops in phase-space [34, 167] and non-destructively determine the m -qubit with the g -qubit. Other possible types of measurements that have been performed in dual species systems include parity readout using QLS [134] and Bell state stabilization [136] by using a g -qubit and two m -qubits. This will require expansion of trapping to three ions (minimal changes to fluorescence detection operations) and using protected motional modes to selectively designate the outer ions as m -qubits.

Further demonstrations and improvements of the effectiveness of the omg architecture are key to developing scalable, fault-tolerant systems with quantum error correction and moving beyond the NISQ regime.

REFERENCES CITED

- [1] Y. Manin, *Computable and uncomputable*, Sovetskoye Radio, Moscow **128**, 15 (1980).
- [2] N. P. de Leon, K. M. Itoh, D. Kim, K. K. Mehta, T. E. Northup, H. Paik, B. S. Palmer, N. Samarth, S. Sangtawesin, and D. W. Steuerman, *Materials challenges and opportunities for quantum computing hardware*, [Science](#) **372**, eabb2823 (2021).
- [3] G. Popkin, *Update: Quantum physics gets attention—and brighter funding prospects—in congress*, [Science](#) (2018).
- [4] J. Preskill, *Quantum Computing in the NISQ era and beyond*, [Quantum](#) **2**, 79 (2018).
- [5] J. W. Z. Lau, K. H. Lim, H. Shrotriya, and L. C. Kwek, *NISQ computing: where are we and where do we go?*, [AAPPS Bulletin](#) **32**, 27 (2022).
- [6] A. Bermudez, X. Xu, R. Nigmatullin, J. O’Gorman, V. Negnevitsky, P. Schindler, T. Monz, U. G. Poschinger, C. Hempel, J. Home, F. Schmidt-Kaler, M. Biercuk, R. Blatt, S. Benjamin, and M. Müller, *Assessing the progress of trapped-ion processors towards fault-tolerant quantum computation*, [Physical Review X](#) **7**, 041061 (2017).
- [7] P. Wang, C.-Y. Luan, M. Qiao, M. Um, J. Zhang, Y. Wang, X. Yuan, M. Gu, J. Zhang, and K. Kim, *Single ion qubit with estimated coherence time exceeding one hour*, [Nature Communications](#) **12**, 233 (2021).
- [8] K. R. Brown, J. Chiaverini, J. M. Sage, and H. Häffner, *Materials challenges for trapped-ion quantum computers*, [Nature Reviews Materials](#) **6**, 892–905 (2021).
- [9] M. Smith, A. Leu, K. Miyanishi, M. Gely, and D. Lucas, *Single-qubit gates with errors at the 10^{-7} level*, [arXiv preprint arXiv:2412.04421](#) (2024).
- [10] T. P. Harty, D. T. C. Allcock, C. J. Ballance, L. Guidoni, H. A. Janacek, N. M. Linke, D. N. Stacey, and D. M. Lucas, *High-fidelity preparation, gates, memory, and readout of a trapped-ion quantum bit*, [Physical Review Letters](#) **113**, 220501 (2014).

- [11] C. M. Löschnauer, J. M. Toba, A. C. Hughes, S. A. King, M. A. Weber, R. Srinivas, R. Matt, R. Nourshargh, D. T. C. Allcock, C. J. Ballance, C. Matthiesen, M. Malinowski, and T. P. Harty, *Scalable, high-fidelity all-electronic control of trapped-ion qubits*, [arXiv preprint arXiv:2407.07694](#) (2024).
- [12] R. Srinivas, S. C. Burd, H. M. Knaack, R. T. Sutherland, A. Kwiatkowski, S. Glancy, E. Knill, D. J. Wineland, D. Leibfried, A. C. Wilson, D. T. C. Allcock, and D. H. Slichter, *High-fidelity laser-free universal control of trapped ion qubits*, [Nature](#) **597**, 209–213 (2021).
- [13] D. Kielpinski, C. Monroe, and D. J. Wineland, *Architecture for a large-scale ion-trap quantum computer*, [Nature](#) **417**, 709–711 (2002).
- [14] J. M. Pino, J. M. Dreiling, C. Figgatt, J. P. Gaebler, S. A. Moses, M. S. Allman, C. H. Baldwin, M. Foss-Feig, D. Hayes, K. Mayer, C. Ryan-Anderson, and B. Neyenhuis, *Demonstration of the trapped-ion quantum CCD computer architecture*, [Nature](#) **592**, 209–213 (2021).
- [15] A. Bermudez, X. Xu, R. Nigmatullin, J. O’Gorman, V. Negnevitsky, P. Schindler, T. Monz, U. G. Poschinger, C. Hempel, J. Home, F. Schmidt-Kaler, M. Biercuk, R. Blatt, S. Benjamin, and M. Müller, *Assessing the progress of trapped-ion processors towards fault-tolerant quantum computation*, [Physical Review X](#) **7**, 041061 (2017).
- [16] C. Hempel, C. Maier, J. Romero, J. McClean, T. Monz, H. Shen, P. Jurcevic, B. P. Lanyon, P. Love, R. Babbush, A. Aspuru-Guzik, R. Blatt, and C. F. Roos, *Quantum chemistry calculations on a trapped-ion quantum simulator*, [Physical Review X](#) **8**, 031022 (2018).
- [17] S. Pal, M. Bhattacharya, S. Dash, S.-S. Lee, and C. Chakraborty, *Future potential of quantum computing and simulations in biological science*, [Molecular Biotechnology](#) **66**, 2201–2218 (2024).
- [18] L. B. Oftelie, M. Urbanek, M. Metcalf, J. Carter, A. F. Kemper, and W. A. de Jong, *Simulating quantum materials with digital quantum computers*, [Quantum Science and Technology](#) **6**, 043002 (2021).

- [19] C. W. Bauer, Z. Davoudi, A. B. Balantekin, T. Bhattacharya, M. Carena, W. A. de Jong, P. Draper, A. El-Khadra, N. Gemelke, M. Hanada, D. Kharzeev, H. Lamm, Y.-Y. Li, J. Liu, M. Lukin, Y. Meurice, C. Monroe, B. Nachman, G. Pagano, J. Preskill, E. Rinaldi, A. Roggero, D. I. Santiago, M. J. Savage, I. Siddiqi, G. Siopsis, D. Van Zanten, N. Wiebe, Y. Yamauchi, K. Yeter-Aydeniz, and S. Zorzetti, *Quantum simulation for high-energy physics*, [PRX Quantum](#) **4**, 027001 (2023).
- [20] C. L. Degen, F. Reinhard, and P. Cappellaro, *Quantum sensing*, [Reviews of Modern Physics](#) **89**, 035002 (2017).
- [21] Q. Zhuang, J. Preskill, and L. Jiang, *Distributed quantum sensing enhanced by continuous-variable error correction*, [New Journal of Physics](#) **22**, 022001 (2020).
- [22] J. Metzner, A. Quinn, S. Brudney, I. D. Moore, S. C. Burd, D. J. Wineland, and D. T. C. Allcock, *Two-mode squeezing and $SU(1,1)$ interferometry with trapped ions*, [Physical Review A](#) **110**, 022613 (2024).
- [23] K. A. Gilmore, M. Affolter, R. J. Lewis-Swan, D. Barberena, E. Jordan, A. M. Rey, and J. J. Bollinger, *Quantum-enhanced sensing of displacements and electric fields with two-dimensional trapped-ion crystals*, [Science](#) **373**, 673–678 (2021).
- [24] P. Drmota, D. Main, D. P. Nadlinger, B. C. Nichol, M. A. Weber, E. M. Ainley, A. Agrawal, R. Srinivas, G. Araneda, C. J. Ballance, and D. M. Lucas, *Robust quantum memory in a trapped-ion quantum network node*, [Physical Review Letters](#) **130**, 090803 (2023).
- [25] C. Bruzewicz, R. McConnell, J. Stuart, J. Sage, and J. Chiaverini, *Dual-species, multi-qubit logic primitives for Ca^+/Sr^+ trapped-ion crystals*, [npj Quantum Information](#) **5**, 102 (2019).
- [26] F. Lancellotti, S. Welte, M. Simoni, C. Mordini, T. Behrle, B. de Neeve, M. Marinelli, V. Negnevitsky, and J. Home, *Low-excitation transport and separation of high-mass-ratio mixed-species ion chains*, [Physical Review Research](#) **6**, L032059 (2024).
- [27] D. T. C. Allcock, W. C. Campbell, J. Chiaverini, I. L. Chuang, E. R. Hudson, I. D. Moore, A. Ransford, C. Roman, J. M. Sage, and D. J. Wineland, *omgblueprint for trapped ion quantum computing with metastable states*, [Applied Physics Letters](#) **119**, 214002 (2021).

- [28] H.-X. Yang, J.-Y. Ma, Y.-K. Wu, Y. Wang, M.-M. Cao, W.-X. Guo, Y.-Y. Huang, L. Feng, Z.-C. Zhou, and L.-M. Duan, *Realizing coherently convertible dual-type qubits with the same ion species*, [Nature Physics](#) **18**, 1058–1061 (2022).
- [29] J. Chiaverini, D. Leibfried, T. Schaetz, M. D. Barrett, R. B. Blakestad, J. Britton, W. M. Itano, J. D. Jost, E. Knill, C. Langer, R. Ozeri, and D. J. Wineland, *Realization of quantum error correction*, [Nature](#) **432**, 602–605 (2004).
- [30] B. M. Terhal, *Quantum error correction for quantum memories*, [Reviews of Modern Physics](#) **87**, 307–346 (2015).
- [31] I. D. Moore, *Easy on the Ions: Photon Scattering Errors from Far-Detuned Raman Beams in Trapped-Ion Qubits*, [Ph.D. thesis](#), University of Oregon (2023).
- [32] J. M. Metzner, *The Sound of Ions: Using Trapped Atomic Ion Motion for Quantum Computation and Sensing*, [Ph.D. thesis](#), University of Oregon (2024).
- [33] A. D. Quinn, *Novel Entangling Gates and Scalable Trap Designs for Trapped-Ion Quantum Computing*, [Ph.D. thesis](#), University of Oregon (2024).
- [34] A. Quinn, G. Gregory, I. Moore, S. Brudney, J. Metzner, E. Ritchie, J. O’Reilly, D. Wineland, and D. Allcock, *High-fidelity entanglement of metastable trapped-ion qubits with integrated erasure conversion*, [arXiv preprint arXiv:2411.12727](#) (2024).
- [35] W. Paul and H. Steinwedel, *Notizen: Ein neues massenspektrometer ohne magnetfeld*, [Zeitschrift für Naturforschung A](#) **8**, 448–450 (1953).
- [36] W. Paul, *Electromagnetic traps for charged and neutral particles*, [Reviews of Modern Physics](#) **62**, 531–540 (1990).
- [37] R. Blatt and D. J. Wineland, *Entangled states of trapped atomic ions*, [Nature](#) **453**, 1008–1015 (2008).
- [38] S. Seidelin, J. Chiaverini, R. Reichle, J. J. Bollinger, D. Leibfried, J. Britton, J. H. Wesenberg, R. B. Blakestad, R. J. Epstein, D. B. Hume, W. M. Itano, J. D. Jost, C. Langer, R. Ozeri, N. Shiga, and D. J. Wineland, *Microfabricated surface-electrode ion trap for scalable quantum information processing*, [Physical Review Letters](#) **96**, 253003 (2006).
- [39] N. Glikin, B. A. Stickler, R. Tollefsen, S. Mouradian, N. Yadav, E. Urban, K. Hornberger, and H. Häffner, *Probing rotational decoherence with a trapped-ion planar rotor*, [Physical Review Letters](#) **134**, 033601 (2025).

- [40] S. A. Moses, C. H. Baldwin, M. S. Allman, R. Ancona, L. Ascarrunz, C. Barnes, J. Bartolotta, B. Bjork, P. Blanchard, M. Bohn, J. G. Bohnet, N. C. Brown, N. Q. Burdick, W. C. Burton, S. L. Campbell, J. P. Campora, C. Carron, J. Chambers, J. W. Chan, Y. H. Chen, A. Chernoguzov, E. Chertkov, J. Colina, J. P. Curtis, R. Daniel, M. DeCross, D. Deen, C. Delaney, J. M. Dreiling, C. T. Ertsgaard, J. Esposito, B. Estey, M. Fabrikant, C. Figgatt, C. Foltz, M. Foss-Feig, D. Francois, J. P. Gaebler, T. M. Gatterman, C. N. Gilbreth, J. Giles, E. Glynn, A. Hall, A. M. Hankin, A. Hansen, D. Hayes, B. Higashi, I. M. Hoffman, B. Horning, J. J. Hout, R. Jacobs, J. Johansen, L. Jones, J. Karcz, T. Klein, P. Lauria, P. Lee, D. Liefer, S. T. Lu, D. Lucchetti, C. Lytle, A. Malm, M. Matheny, B. Mathewson, K. Mayer, D. B. Miller, M. Mills, B. Neyenhuis, L. Nugent, S. Olson, J. Parks, G. N. Price, Z. Price, M. Pugh, A. Ransford, A. P. Reed, C. Roman, M. Rowe, C. Ryan-Anderson, S. Sanders, J. Sedlacek, P. Shevchuk, P. Siegfried, T. Skripka, B. Spaun, R. T. Sprenkle, R. P. Stutz, M. Swallows, R. I. Tobey, A. Tran, T. Tran, E. Vogt, C. Volin, J. Walker, A. M. Zolot, and J. M. Pino, *A race-track trapped-ion quantum processor*, [Physical Review X **13**, 041052 \(2023\)](#).
- [41] R. J. Niffenegger, J. Stuart, C. Sorace-Agaskar, D. Kharas, S. Bramhavar, C. D. Bruzewicz, W. Loh, R. T. Maxson, R. McConnell, D. Reens, G. N. West, J. M. Sage, and J. Chiaverini, *Integrated multi-wavelength control of an ion qubit*, [Nature **586**, 538–542 \(2020\)](#).
- [42] X. Fernandez-Gonzalvo and M. Keller, *A fully fiber-integrated ion trap for portable quantum technologies*, [Scientific Reports **13**, 523 \(2023\)](#).
- [43] O. Wipfli, H. F. Passagem, C. Fischer, M. Grau, and J. P. Home, *Integration of a high finesse cryogenic build-up cavity with an ion trap*, [Review of Scientific Instruments **94**, 083204 \(2023\)](#).
- [44] G. Araneda, G. Cerchiari, D. B. Higginbottom, P. C. Holz, K. Lakhmanskiy, P. Obšil, Y. Colombe, and R. Blatt, *The Panopticon device: An integrated Paul-trap–hemispherical mirror system for quantum optics*, [Review of Scientific Instruments **91**, 113201 \(2020\)](#).
- [45] D. J. Wineland, C. Monroe, W. M. Itano, D. Leibfried, B. E. King, and D. M. Meekhof, *Experimental issues in coherent quantum-state manipulation of trapped atomic ions*, [Journal of Research of the National Institute of Standards and Technology **103**, 259 \(1998\)](#).
- [46] W. Jones, *Earnshaw’s theorem and the stability of matter*, [European Journal of Physics **1**, 85 \(1980\)](#).

- [47] D. J. Berkeland, J. D. Miller, J. C. Bergquist, W. M. Itano, and D. J. Wineland, *Minimization of ion micromotion in a Paul trap*, [Journal of Applied Physics](#) **83**, 5025–5033 (1998).
- [48] C. J. Foot, Atomic physics (Oxford University, 2004) Chap. 12, pp. 259–281.
- [49] J. Pedregosa, C. Champenois, M. Houssin, and M. Knoop, *Anharmonic contributions in real RF linear quadrupole traps*, [International Journal of Mass Spectrometry](#) **290**, 100–105 (2010).
- [50] M. Brownnutt, M. Kumph, P. Rabl, and R. Blatt, *Ion-trap measurements of electric-field noise near surfaces*, [Reviews of Modern Physics](#) **87**, 1419–1482 (2015).
- [51] C. Monroe, D. M. Meekhof, B. E. King, S. R. Jefferts, W. M. Itano, D. J. Wineland, and P. Gould, *Resolved-Sideband Raman Cooling of a Bound Atom to the 3D Zero-Point Energy*, [Physical Review Letters](#) **75**, 4011–4014 (1995).
- [52] D. J. Wineland, J. C. Bergquist, W. M. Itano, J. J. Bollinger, and C. H. Manney, *Atomic-ion coulomb clusters in an ion trap*, [Physical Review Letters](#) **59**, 2935–2938 (1987).
- [53] V. Krutyanskiy, M. Galli, V. Krcmarsky, S. Baier, D. A. Fioretto, Y. Pu, A. Mazloom, P. Sekatski, M. Canteri, M. Teller, J. Schupp, J. Bate, M. Meraner, N. Sangouard, B. P. Lanyon, and T. E. Northup, *Entanglement of trapped-ion qubits separated by 230 meters*, [Physical Review Letters](#) **130**, 050803 (2023).
- [54] C. Crocker, M. Lichtman, K. Sosnova, A. Carter, S. Scarano, and C. Monroe, *High purity single photons entangled with an atomic qubit*, [Optics Express](#) **27**, 28143–28149 (2019).
- [55] L. J. Stephenson, D. P. Nadlinger, B. C. Nichol, S. An, P. Drmota, T. G. Ballance, K. Thirumalai, J. F. Goodwin, D. M. Lucas, and C. J. Ballance, *High-rate, high-fidelity entanglement of qubits across an elementary quantum network*, [Physical Review Letters](#) **124**, 110501 (2020).
- [56] J. F. Fitzsimons, *Private quantum computation: an introduction to blind quantum computing and related protocols*, [npj Quantum Information](#) **3**, 1–11 (2017).
- [57] OregonIons, Squareatron-5000, <https://github.com/OregonIons/Squareatron-5000> (2020).

- [58] J. Siversns, L. Simkins, S. Weidt, and W. Hensinger, *On the application of radio frequency voltages to ion traps via helical resonators*, [Applied Physics B](#) **107**, 921–934 (2012).
- [59] P. S. Berman and V. S. Malinovsky, *Principles of Laser Spectroscopy and Quantum Optics* (Princeton University Press, 2011).
- [60] P. Meystre and M. Sargent, *Elements of Quantum Optics*, 4th ed. (Springer Science & Business Media, 2007).
- [61] S. Lloyd and S. L. Braunstein, *Quantum computation over continuous variables*, [Physical Review Letters](#) **82**, 1784–1787 (1999).
- [62] H.-K. Lau and D. F. V. James, *Proposal for a scalable universal bosonic simulator using individually trapped ions*, [Physical Review A](#) **85**, 062329 (2012).
- [63] D. M. Meekhof, C. Monroe, B. E. King, W. M. Itano, and D. J. Wineland, *Generation of nonclassical motional states of a trapped atom*, [Physical Review Letters](#) **76**, 1796–1799 (1996).
- [64] L. Ortiz-Gutiérrez, B. Gabrielly, L. F. Muñoz, K. T. Pereira, J. G. Filgueiras, and A. S. Villar, *Continuous variables quantum computation over the vibrational modes of a single trapped ion*, [Optics Communications](#) **397**, 166–174 (2017).
- [65] D. Gottesman, A. Kitaev, and J. Preskill, *Encoding a qubit in an oscillator*, [Physical Review A](#) **64**, 012310 (2001).
- [66] R. W. Heeres, P. Reinhold, N. Ofek, L. Frunzio, L. Jiang, M. H. Devoret, and R. J. Schoelkopf, *Implementing a universal gate set on a logical qubit encoded in an oscillator*, [Nature Communications](#) **8**, 1–7 (2017).
- [67] S. C. Burd, R. Srinivas, J. J. Bollinger, A. C. Wilson, D. J. Wineland, D. Leibfried, D. H. Slichter, and D. T. C. Allcock, *Quantum amplification of mechanical oscillator motion*, [Science](#) **364**, 1163–1165 (2019).
- [68] C. Flühmann, T. L. Nguyen, M. Marinelli, V. Negnevitsky, K. Mehta, and J. Home, *Encoding a qubit in a trapped-ion mechanical oscillator*, [Nature](#) **566**, 513–517 (2019).
- [69] R. T. Sutherland and R. Srinivas, *Universal hybrid quantum computing in trapped ions*, [Physical Review A](#) **104**, 032609 (2021).

- [70] D. J. Gorman, P. Schindler, S. Selvarajan, N. Daniilidis, and H. Häffner, *Two-mode coupling in a single-ion oscillator via parametric resonance*, [Physical Review A](#) **89**, 062332 (2014).
- [71] G. Kasprowicz, T. Harty, S. Bourdeauducq, R. Jördens, D. Allcock, D. Nadlinger, J. W. Britton, A. Sotirova, and D. Nowicka, *Urukul – Open-source Frequency Synthesizer Module for Quantum Physics*, [International Journal of Electronics and Telecommunications](#) **68**, 123–128 (2022).
- [72] G. Kasprowicz, P. Kulik, M. Gaska, T. Przywozki, K. Pozniak, J. Jarosinski, J. W. Britton, T. Harty, C. Balance, W. Zhang, D. Nadlinger, D. Slichter, D. Allcock, S. Bourdeauducq, R. Jördens, and K. Pozniak, in [OSA Quantum 2.0 Conference](#) (Optica Publishing Group, 2020) p. QTu8B.14.
- [73] R. J. Lewis-Swan, J. C. Castro Zuñiga, D. Barberena, and A. M. Rey, *Exploiting nonclassical motion of a trapped ion crystal for quantum-enhanced metrology of global and differential spin rotations*, [Physical Review Letters](#) **132**, 163601 (2024).
- [74] S. Gao, W. J. Hughes, D. M. Lucas, T. G. Ballance, and J. F. Goodwin, *An optically heated atomic source for compact ion trap vacuum systems*, [Review of Scientific Instruments](#) **92**, 033205 (2021).
- [75] D. R. Leibbrandt, R. J. Clark, J. Labaziewicz, P. Antohi, W. Bakr, K. R. Brown, and I. L. Chuang, *Laser ablation loading of a surface-electrode ion trap*, [Physical Review A](#) **76**, 055403 (2007).
- [76] R. Knight, *Storage of ions from laser-produced plasmas*, [Applied Physics Letters](#) **38**, 221–223 (1981).
- [77] G. Vrijsen, Y. Aikyo, R. F. Spivey, I. V. Inlek, and J. Kim, *Efficient isotope-selective pulsed laser ablation loading of $^{174}\text{Yb}^+$ ions in a surface electrode trap*, [Optics Express](#) **27**, 33907–33914 (2019).
- [78] C. C. Hsu, R. Larue, C. C. Kwong, and D. Wilkowski, *Laser-induced thermal source for cold atoms*, [Scientific Reports](#) **12**, 868 (2022).
- [79] A. Bautista-Salvador, G. Zarantonello, H. Hahn, A. Preciado-Grijalva, J. Morgner, M. Wahnschaffe, and C. Ospelkaus, *Multilayer ion trap technology for scalable quantum computing and quantum simulation*, [New Journal of Physics](#) **21**, 043011 (2019).

- [80] M. Guggemos, D. Heinrich, O. Herrera-Sancho, R. Blatt, and C. Roos, *Sympathetic cooling and detection of a hot trapped ion by a cold one*, [New Journal of Physics](#) **17**, 103001 (2015).
- [81] G. Pagano, P. W. Hess, H. B. Kaplan, W. L. Tan, P. Richerme, P. Becker, A. Kyprianidis, J. Zhang, E. Birckelbaw, M. R. Hernandez, Y. Wu, and C. Monroe, *Cryogenic trapped-ion system for large scale quantum simulation*, [Quantum Science and Technology](#) **4**, 014004 (2018).
- [82] O. Kock, W. He, D. Świerad, L. Smith, J. Hughes, K. Bongs, and Y. Singh, *Laser controlled atom source for optical clocks*, [Scientific Reports](#) **6**, 1–6 (2016).
- [83] T. E. Madey, R. E. Johnson, and T. M. Orlando, *Far-out surface science: radiation-induced surface processes in the solar system*, [Surface Science](#) **500**, 838–858 (2002).
- [84] A. Gozzini, F. Mango, J. Xu, G. Alzetta, F. Maccarrone, and R. Bernheim, *Light-induced ejection of alkali atoms in polysiloxane coated cells*, [Il Nuovo Cimento D](#) **15**, 709–722 (1993).
- [85] D. S. Barker, E. B. Norrgard, J. Scherschligt, J. A. Fedchak, and S. Eckel, *Light-induced atomic desorption of lithium*, [Physical Review A](#) **98**, 043412 (2018).
- [86] V. Coppolaro, N. Papi, A. Khanbekyan, C. Marinelli, E. Mariotti, L. Marmugi, L. Moi, L. Corradi, A. Dainelli, H. Arikawa, T. Ishikawa, Y. Sakemi, R. Calabrese, G. Mazzocca, L. Tomassetti, and L. Ricci, *Light desorption from an yttrium neutralizer for Rb and Fr magneto-optical trap loading*, [The Journal of Chemical Physics](#) **141**, 134201 (2014).
- [87] C. Klempt, T. van Zoest, T. Henninger, O. Topic, E. Rasel, W. Ertmer, and J. Arlt, *Ultraviolet light-induced atom desorption for large rubidium and potassium magneto-optical traps*, [Physical Review A](#) **73**, 013410 (2006).
- [88] L. Torralbo-Campo, G. D. Bruce, G. Smirne, and D. Cassettari, *Light-induced atomic desorption in a compact system for ultracold atoms*, [Scientific Reports](#) **5**, 1–10 (2015).
- [89] M. Yasuda, T. Tanabe, T. Kobayashi, D. Akamatsu, T. Sato, and A. Hatakeyama, *Laser-controlled cold ytterbium atom source for transportable optical clocks*, [Journal of the Physical Society of Japan](#) **86**, 125001 (2017).

- [90] K. Kromer, C. Lyu, J. Bieroń, M. Door, L. Enzmann, P. Filianin, G. Gaigalas, Z. Harman, J. Herkenhoff, W. Huang, C. H. Keitel, S. Eliseev, and K. Blaum, *Atomic mass determination of uranium-238*, [Physical Review C **109**, L021301 \(2024\)](#).
- [91] C. Schweiger, C. König, J. R. Crespo López-Urrutia, M. Door, H. Dorrer, C. E. Düllmann, S. Eliseev, P. Filianin, W. Huang, K. Kromer, P. Micke, M. Müller, D. Renisch, A. Rischka, R. X. Schüssler, and K. Blaum, *Production of highly charged ions of rare species by laser-induced desorption inside an electron beam ion trap*, [Review of Scientific Instruments **90** \(2019\)](#).
- [92] F. Mango and E. Maccioni, *Light-induced ejection of calcium atoms from polymer surfaces*, [The European Physical Journal D **50**, 253–256 \(2008\)](#).
- [93] D. M. Lucas, A. Ramos, J. P. Home, M. J. McDonnell, S. Nakayama, J.-P. Stacey, S. C. Webster, D. N. Stacey, and A. M. Steane, *Isotope-selective photoionization for calcium ion trapping*, [Physical Review A **69**, 012711 \(2004\)](#).
- [94] G. Imreh, *Implementing segmented ion trap designs for quantum computing*, [Ph.D. thesis](#), St. Peter's College Oxford (2008).
- [95] A. Bonch-Bruевич, T. Vartanyan, A. Gorlanov, Y. N. Maksimov, S. Przhibelski, and V. Khromov, *Photodesorption of sodium from the surface of sapphire*, [Journal of Experimental and Theoretical Physics **97**, 1077–1185 \(1990\)](#).
- [96] E. J. Salumbides, V. Maslinskas, I. M. Dildar, A. L. Wolf, E.-J. van Duijn, K. S. E. Eikema, and W. Ubachs, *High-precision frequency measurement of the 423-nm Ca i line*, [Physical Review A **83**, 012502 \(2011\)](#).
- [97] J. Rennie and S. Elliott, *Investigations of the mechanism of photo-dissolution of silver in amorphous germanium chalcogenide thin films*, [Journal of Non-Crystalline Solids **97-98**, 1239–1242 \(1987\)](#).
- [98] A. Prakash, A. Ayyadevara, E. Krishnakumar, and S. A. Rangwala, *Low divergence cold-wall oven for loading ion traps*, [Review of Scientific Instruments **95**, 033202 \(2024\)](#).
- [99] C. Becher, W. Gao, S. Kar, C. D. Marciniak, T. Monz, J. G. Bartholomew, P. Goldner, H. Loh, E. Marcellina, K. E. J. Goh, T. S. Koh, B. Weber, Z. Mu, J.-Y. Tsai, Q. Yan, T. Huber-Loyola, S. Höfling, S. Gyger, S. Steinhauer, and V. Zwiller, *2023 roadmap for materials for quantum technologies*, [Materials for Quantum Technology \(2022\)](#).

- [100] J. Mitroy, *Energy levels and oscillator strengths for neutral calcium*, [Journal of Physics B: Atomic, Molecular and Optical Physics](#) **26**, 3703 (1993).
- [101] M. Hettrich, T. Ruster, H. Kaufmann, C. F. Roos, C. T. Schmiegelow, F. Schmidt-Kaler, and U. G. Poschinger, *Measurement of dipole matrix elements with a single trapped ion*, [Physical Review Letters](#) **115**, 143003 (2015).
- [102] A. Kreuter, C. Becher, G. P. T. Lancaster, A. B. Mundt, C. Russo, H. Häffner, C. Roos, W. Hänsel, F. Schmidt-Kaler, R. Blatt, and M. S. Safronova, *Experimental and theoretical study of the $3d\ ^2D$ -level lifetimes of $^{40}\text{Ca}^+$* , [Physical Review A](#) **71**, 032504 (2005).
- [103] G. Morigi, J. Eschner, and C. H. Keitel, *Ground state laser cooling using electromagnetically induced transparency*, [Physical Review Letters](#) **85**, 4458 (2000).
- [104] R. Lechner, C. Maier, C. Hempel, P. Jurcevic, B. P. Lanyon, T. Monz, M. Brownnutt, R. Blatt, and C. F. Roos, *Electromagnetically-induced-transparency ground-state cooling of long ion strings*, [Physical Review A](#) **93**, 053401 (2016).
- [105] S. Stenholm, *The semiclassical theory of laser cooling*, [Reviews of Modern Physics](#) **58**, 699 (1986).
- [106] L.-L. Yan, L.-Y. Wang, S.-L. Su, F. Zhou, and M. Feng, *Verification of information thermodynamics in a trapped ion system*, [Entropy](#) **24** (2022).
- [107] B. Lounis and C. Cohen-Tannoudji, *Coherent population trapping and Fano profiles*, [Journal de Physique II](#) **2**, 579–592 (1992).
- [108] M. Ghadimi, E. M. Bridge, J. Scarabel, S. Connell, K. Shimizu, E. Streed, and M. Lobino, *Multichannel optomechanical switch and locking system for wavemeters*, [Applied Optics](#) **59**, 5136–5141 (2020).
- [109] J. Kim, K. Kim, D. Lee, Y. Shin, S. Kang, J.-R. Kim, Y. Choi, K. An, and M. Lee, *Locking multi-laser frequencies to a precision wavelength meter: Application to cold atoms*, [Sensors](#) **21** (2021).
- [110] J. Helmcke, S. A. Lee, and J. L. Hall, *Dye laser spectrometer for ultrahigh spectral resolution: design and performance*, [Applied Optics](#) **21**, 1686–1694 (1982).

- [111] J. Berrocal, E. Altozano, F. Domínguez, M. J. Gutiérrez, J. Cerrillo, F. J. Fernández, M. Block, C. Ospelkaus, and D. Rodríguez, *Formation of two-ion crystals by injection from a Paul-trap source into a high-magnetic-field Penning trap*, [Physical Review A](#) **105**, 052603 (2022).
- [112] J. Zhang, W. Yuan, K. Deng, A. Deng, Z. Xu, C. Qin, Z. Lu, and J. Luo, *A long-term frequency stabilized deep ultraviolet laser for Mg^+ ions trapping experiments*, [Review of Scientific Instruments](#) **84** (2013).
- [113] S. A. Saakyan, V. A. Sautenkov, E. V. Vilshanskaya, V. Vasiliev, B. B. Zelener, and B. V. Zelener, *Frequency control of tunable lasers using a frequency-calibrated-meter in an experiment on preparation of Rydberg atoms in a magneto-optical trap*, [Quantum Electronics](#) **45**, 828 (2015).
- [114] K. Saleh, J. Millo, A. Didier, Y. Kersalé, and C. Lacroûte, *Frequency stability of a wavelength meter and applications to laser frequency stabilization*, [Applied Optics](#) **54**, 9446–9449 (2015).
- [115] W. Riley and D. Howe, [Handbook of frequency stability analysis](#) (2008).
- [116] J. A. Barnes, in *Proceedings of the 19th Annual Precise Time and Time Interval Systems and Applications Meeting* (Redondo Beach, California, 1989) pp. 227–234.
- [117] A. E. Wallin, [Allantools](#) (2023), Revision 92a8b614, Online; accessed 6 April 2023.
- [118] Hamilton Technical Services, [Stable32](#) (2020), Ver. 1.62, Online; accessed 24 April 2023.
- [119] G. Grynberg, A. Aspect, and C. Fabre, *Introduction to quantum optics* (Cambridge University Press, 2010) Chap. 3 Principles of Lasers.
- [120] T. Shimasaki, E. R. Edwards, and D. DeMille, *Injection-locking of a fiber-pigtailed high-power laser to an external cavity diode laser via a fiber optic circulator*, [Review of Scientific Instruments](#) **90**, 056102 (2019).
- [121] R. Garstang, in *Pure and Applied Physics*, Vol. 13 (Elsevier, 1962) pp. 1–46.
- [122] I. S. Bowen and D. H. Menzel, *Forbidden lines in the flash spectrum*, [Publications of the Astronomical Society of the Pacific](#) **40**, 332–340 (1928).

- [123] E. D. Black, *An introduction to Pound–Drever–Hall laser frequency stabilization*, [American Journal of Physics](#) **69**, 79–87 (2001).
- [124] M. Nickerson, *A review of Pound-Drever-Hall laser frequency locking*, [JILA, University of Colorado and NIST](#) (2019).
- [125] J. Benhelm, *Precision Spectroscopy and Quantum Information Processing with Trapped Calcium Ions*, [Ph.D. thesis](#), University of Innsbruck (2008).
- [126] M. Chwalla, *Precision spectroscopy with $^{40}\text{Ca}^+$ ions in a Paul trap*, [Ph.D. thesis](#), University of Innsbruck (2008).
- [127] M. A. Sepiol, *Frequency stabilization of a 729 nm diode laser to an external high finesse reference cavity*, [Master’s thesis](#), ETH Zurich (2012).
- [128] K. Gao, *Optical frequency standard based on a single $^{40}\text{Ca}^+$* , [Chinese Science Bulletin](#) **58**, 853–863 (2013).
- [129] J. C. Bergquist, R. G. Hulet, W. M. Itano, and D. J. Wineland, *Observation of quantum jumps in a single atom*, [Physical Review Letters](#) **57**, 1699–1702 (1986).
- [130] L. Gerster, *Spectral filtering and laser diode injection for multi-qubit trapped ion gates*, [Master’s thesis](#), ETH Zurich (2015).
- [131] C. F. Roos, *Controlling the quantum state of trapped ions*, [Ph.D. thesis](#), University of Innsbruck (2000).
- [132] D. F. James, *Quantum dynamics of cold trapped ions with application to quantum computation*, [Applied Physics B](#) **66**, 181–190 (1998).
- [133] W. R. Johnson, *Atomic Structure Theory: Lecture on Atomic Physics* (Springer Berlin, Heidelberg, 2007).
- [134] K. Cui, J. Valencia, K. T. Boyce, E. R. Clements, D. R. Leibbrandt, and D. B. Hume, *Scalable quantum logic spectroscopy*, [Physical Review Letters](#) **129**, 193603 (2022).
- [135] J. B. Wübena, S. Amairi, O. Mandel, and P. O. Schmidt, *Sympathetic cooling of mixed-species two-ion crystals for precision spectroscopy*, [Physical Review A](#) **85**, 043412 (2012).
- [136] V. Negnevitsky, M. Marinelli, K. K. Mehta, H.-Y. Lo, C. Flühmann, and J. P. Home, *Repeated multi-qubit readout and feedback with a mixed-species trapped-ion register*, [Nature](#) **563**, 527–531 (2018).

- [137] K. Sosnova, A. Carter, and C. Monroe, *Character of motional modes for entanglement and sympathetic cooling of mixed-species trapped-ion chains*, [Physical Review A **103**, 012610 \(2021\)](#).
- [138] C. J. Ballance, *High-fidelity quantum logic in Ca^+* , [Ph.D. thesis](#), Oxford University (2014).
- [139] Q. A. Turchette, Kielpinski, B. E. King, D. Leibfried, D. M. Meekhof, C. J. Myatt, M. A. Rowe, C. A. Sackett, C. S. Wood, W. M. Itano, C. Monroe, and D. J. Wineland, *Heating of trapped ions from the quantum ground state*, [Physical Review A **61**, 063418 \(2000\)](#).
- [140] H. Kaufmann, T. Ruster, C. T. Schmiegelow, M. A. Luda, V. Kaushal, J. Schulz, D. von Lindenfels, F. Schmidt-Kaler, and U. G. Poschinger, *Fast ion swapping for quantum-information processing*, [Physical Review A **95**, 052319 \(2017\)](#).
- [141] D. Kiesenhofer, H. Hainzer, A. Zhdanov, P. C. Holz, M. Bock, T. Ollikainen, and C. F. Roos, *Controlling two-dimensional coulomb crystals of more than 100 ions in a monolithic radio-frequency trap*, [PRX Quantum **4**, 020317 \(2023\)](#).
- [142] S. Ma, G. Liu, P. Peng, B. Zhang, S. Jandura, J. Claes, A. P. Burgers, G. Pupillo, S. Puri, and J. D. Thompson, *High-fidelity gates and mid-circuit erasure conversion in an atomic qubit*, [Nature **622**, 279–284 \(2023\)](#).
- [143] Y. Wu, S. Kolkowitz, S. Puri, and J. D. Thompson, *Erasure conversion for fault-tolerant quantum computing in alkaline earth rydberg atom arrays*, [Nature Communications **13**, 4657 \(2022\)](#).
- [144] G. Shu, G. Vittorini, A. Buikema, C. S. Nichols, C. Volin, D. Stick, and K. R. Brown, *Heating rates and ion-motion control in a Y-junction surface-electrode trap*, [Physical Review A **89**, 062308 \(2014\)](#).
- [145] V. Kaushal, B. Lekitsch, A. Stahl, J. Hilder, D. Pijn, C. Schmiegelow, A. Bermudez, M. Müller, F. Schmidt-Kaler, and U. Poschinger, *Shuttling-based trapped-ion quantum information processing*, [AVS Quantum Science **2**, 014101 \(2020\)](#).
- [146] B. Skinner, J. Ruhman, and A. Nahum, *Measurement-induced phase transitions in the dynamics of entanglement*, [Physical Review X **9**, 031009 \(2019\)](#).
- [147] C. Noel, *High temperature studies of electric-field noise in a surface ion trap*, [Ph.D. thesis](#), University of California, Berkeley (2019).

- [148] D. Leibfried, D. M. Meekhof, C. Monroe, B. E. King, W. M. Itano, and D. J. Wineland, *Experimental preparation and measurement of quantum states of motion of a trapped atom*, [Journal of Modern Optics](#) **44**, 2485–2505 (1997).
- [149] R. Loudon, *The Quantum Theory of Light* (Claredon Press, 1983).
- [150] A. J. Rasmusson, I. Jung, F. G. Schroer, A. Kyprianidis, and P. Richerme, *Measurement-induced heating of a trapped ion*, [Journal of Physics B: Atomic, Molecular and Optical Physics](#) **57**, 225002 (2024).
- [151] A. J. Rasmusson, *Trapped Ion Cooling, Heating, and Thermometry*, [Ph.D. thesis](#), Indiana University (2024).
- [152] S. Ejtemaee, R. Thomas, and P. C. Haljan, *Optimization of Yb^+ fluorescence and hyperfine-qubit detection*, [Physical Review A](#) **82**, 063419 (2010).
- [153] J. Eschner, G. Morigi, F. Schmidt-Kaler, and R. Blatt, *Laser cooling of trapped ions*, [Journal of the Optical Society of America B](#) **20**, 1003–1015 (2003).
- [154] M. Chwalla, K. Kim, T. Monz, P. Schindler, M. Riebe, C. Roos, and R. Blatt, *Precision spectroscopy with two correlated atoms*, [Applied Physics B](#) **89**, 483–488 (2007).
- [155] J. I. Cirac and P. Zoller, *Quantum computations with cold trapped ions*, [Physical Review Letters](#) **74**, 4091–4094 (1995).
- [156] S. Schneider and G. J. Milburn, *Decoherence in ion traps due to laser intensity and phase fluctuations*, [Physical Review A](#) **57**, 3748–3752 (1998).
- [157] X. Zhang, Y. Chen, J. Fang, T. Wang, J. Li, and L. Luo, *Beam pointing stabilization of an acousto-optic modulator with thermal control*, [Optics Express](#) **27**, 11503–11509 (2019).
- [158] F. Wolf, Y. Wan, J. C. Heip, F. Gebert, C. Shi, and P. O. Schmidt, *Non-destructive state detection for quantum logic spectroscopy of molecular ions*, [Nature](#) **530**, 457–460 (2016).
- [159] M. DeCross, E. Chertkov, M. Kohagen, and M. Foss-Feig, *Qubit-reuse compilation with mid-circuit measurement and reset*, [Physical Review X](#) **13**, 041057 (2023).
- [160] C. Piveteau and D. Sutter, *Circuit knitting with classical communication*, [IEEE Transactions on Information Theory](#) **70**, 2734–2745 (2024).

- [161] S. A. Schulz, U. Poschinger, F. Ziesel, and F. Schmidt-Kaler, *Sideband cooling and coherent dynamics in a microchip multi-segmented ion trap*, [New Journal of Physics](#) **10**, 045007 (2008).
- [162] C. Lee, S. C. Webster, J. Mosca Toba, O. Corfield, G. Porter, and R. C. Thompson, *Measurement-based ground-state cooling of a trapped-ion oscillator*, [Physical Review A](#) **107**, 033107 (2023).
- [163] J. Eschner, B. Appasamy, and P. E. Toschek, *Stochastic cooling of a trapped ion by null detection of its fluorescence*, [Physical Review Letters](#) **74**, 2435–2438 (1995).
- [164] P. Scholl, A. L. Shaw, R. B.-S. Tsai, R. Finkelstein, J. Choi, and M. Endres, *Erasure conversion in a high-fidelity rydberg quantum simulator*, [Nature](#) **622**, 273–278 (2023).
- [165] P. O. Schmidt, T. Rosenband, C. Langer, W. M. Itano, J. C. Bergquist, and D. J. Wineland, *Spectroscopy using quantum logic*, [Science](#) **309**, 749–752 (2005).
- [166] I. Vogelius, L. B. Madsen, and M. Drewsen, *Probabilistic state preparation of a single molecular ion by projection measurement*, [Journal of Physics B: Atomic, Molecular and Optical Physics](#) **39**, S1259 (2006).
- [167] D. Leibfried, B. DeMarco, V. Meyer, D. Lucas, M. Barrett, J. Britton, W. M. Itano, B. Jelenković, C. Langer, T. Rosenband, and D. J. Wineland, *Experimental demonstration of a robust, high-fidelity geometric two ion-qubit phase gate*, [Nature](#) **422**, 412–415 (2003).