A Compact Cryogenic Package Approach to Ion Trap Quantum Computing

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Crystal Noel

Dissertation submitted in partial fulfillment of the requirements for the degree of Doctor of Philosophy in the Department of Electrical and Computer Engineering in the Graduate School of Duke University 2022
ABSTRACT

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Abstract

Ion traps are a leading candidate for scaling quantum computers. The component technologies can be difficult to integrate and manufacture. Experimental systems are also subject to mechanical drift creating a large maintenance overhead. A full system redesign with stability and scalability in mind is presented. The center of our approach is a compact cryogenic ion trap package (trap cryopackage). A surface trap is mounted to a modified ceramic pin grid array (CPGA) this is enclosed using a copper lid. The differentially pumped trap cryopackage has all necessary optical feedthroughs and an ion source (ablation target). The lid pressure is held at ultra-high vacuum (UHV) by cryogenic sorption pumping using carbon getter. We install this cryopackage into a commercial low-vibration closed-cycle cryostat which sits inside a custom monolithic enclosure. The system is tested and trapped ions are found to have common mode heating rate on the order of 10 quanta/s. The modular optical setup provides for a counterpropagating single qubit coherence time of 527 ms. We survey a population of FM two-qubit gates (gate times 120 µs - 450 µs) and find an average gate fidelity of 98%. We study the gate survey with quantum Monte Carlo simulation and find that our two-qubit gate fidelity is limited by low frequency (30 Hz - 3 kHz) coherent electrical noise on our motional modes.
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List of Abbreviations

AOM  acousto-optic modulator
CPGA  ceramic pin grid array
CW  continuous wave
DAC  digital-to-analog converter
DC  direct current
DDS  direct digital synthesizer
EIT  electromagnetically-induced-transparency
ELF  extremely low frequency
EMI  electromagnetic interference
EOM  electro-optic modulator
FM  frequency-modulated
GST  gate set tomography
ME  master equation
MEMs  Microelectromechanical systems
MS  Mølmer–Sørensen
NA  numerical aperture
NISQ  noisy intermediate-scale quantum
PBS  polarizing beam splitter
PCB  printed circuit board
PID  proportional-integral-derivative
PSD  power spectral density
RF  radio frequency
SPAM  state preparation and measurement

turbo  turbomolecular pump

UHV  ultra-high vacuum

UV  ultraviolet
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Lastly I would like to thank my Mom. She has been so supportive every step of my academic and personal life. I am lucky to have such a wonderful Mother. She has always been so invested in my learning and growth. She spent countless hours in my youth helping me when I was struggling or supporting my dreams when I was doing well. She inspired me through sharing memories of her Father Roger Bolz, who got an engineering degree by going to night school at Case Western Reserve. I so grateful to have parents that invested in my education and always encouraged my wild interests.
Merging of Difference and Unity

The mind of the great sage of India is intimately communicated between East and West: peoples’ faculties may be keen or dull, but in the Path there are no southern or northern ancestors. The spiritual source shines clearly in the light, branching streams flow in the darkness; grasping things is basically delusion, merging with principle is still not enlightenment. Each sense and every field interact and do not interact; when interacting they also merge, otherwise they remain in their own states. Forms are basically different in material and appearance; sounds are fundamentally different in pleasant or harsh quality; darkness is a word for merging upper and lower, light is an expression for distinguishing pure and defiled. The four gross elements return to their own natures like a baby taking to its mother: fire heats, wind moves, water wets, earth is solid; eye and form, ear and sound, nose and smell, tongue and taste—thus in all things the leaves spread from the root. The whole process must return to the source. Noble and base are only manners of speaking; right in light there is darkness but don’t confront it as darkness, right in darkness there is light but don’t see it as light. Light and dark are relative to one another like forward and backward steps. All things have their function; it is a matter of use in the appropriate situation. Phenomena exist like box and cover joining; principle accords like arrow points meeting. Hearing the words you should understand the source; don’t make up standards on your own. If you don’t understand the path as it meets your eyes, how can you know the way as you walk? Progress is not a matter of far or near, but if you are confused mountains and rivers block the way. I humbly say to those who study the mystery: don’t waste time.

The Merging of Difference and Unity or Sandōkai was written by Shítóu Xǐqiān in China in the 8th Century. Translated by Thomas Cleary printed here under fair use as an excerpt from the book Timeless Spring © 1980 Tuttle Publishing
Chapter 1

Introduction

The exponential growth of computing power, called Moore’s law [Mac11], has been a large driver of late 20th century and early 21st century boom in human productivity and economic growth [Mul16]. Moore’s law has been enabled by a consistent trend in the ability of the semiconductor industry to miniaturize the transistor, especially the dominant complementary metal–oxide–semiconductor (CMOS) platform. The size of transistors is reaching physical limits [MG15]. Additionally incentives in the semiconductor industry have shifted due to a splintering of the industry into many highly profitable divergent technologies [KHF18] which don’t address the miniaturization problem. There is some indication that other non-CMOS based electron transport technologies for classical computing will face similar limits due to heat dissipation [KHF18].

Quantum computers provide a fundamentally new computing paradigm which could open up new horizons for humanity. Quantum computers function on qubits which make available the resources of superposition, entanglement, and Hilbert space scaling to open up new algorithmic frontiers [NC10]. The most important example, which launched excitement in the field, is Shor’s algorithm [Sho94] which can solve the prime factorization problem in polynomial time, a feat which has exponential scaling on classical computers. Additionally as Richard Feynman first noted, a quantum system would be uniquely suited for the full simulation of other quantum systems [FHA18]. Quantum computers then could find unique use for chemistry and materials science [BMK10, KWPO+11, BBMC20, MGG20, BUM+21, ZJJ+22]. Ad-
ditionally Grover’s search algorithm can provide a quadratic speed-up over classical counterparts [Gro97].

There are many different physical platforms for storing and manipulating quantum information [SR21]. Ion traps are a leading qubit platform with experimentally verified high fidelity operation [WMI⁺98, MK13, BW08]. Atomic ions provide intrinsically stable qubits when hyperfine levels are used [LOJ⁺05, HAB⁺14, WUZ⁺17]. High-fidelity state preparation and detection [MSW⁺08, NVG⁺13, CCV⁺19], as well as logic gates driven by both lasers [HAB⁺14, MKC⁺15, BKGN⁺17, BHL⁺16, GTL⁺16, WCF⁺20] and microwave fields are readily available [WMI⁺98]. The performance of trapped ion qubits arises from the fact that these atoms are well isolated from the environment under ultra-high vacuum (UHV) conditions, and the control signals can be delivered to the qubits in the form of electromagnetic fields. Ions are a natural qubit so they are not subject to manufacturing variation. In order to create quantum computers which can run non-trivial quantum algorithms we require a large number of error corrected qubits [VMH13]. This is only possible with a modular integrated platform.

1.1 Rethinking the Ion Trap

The industrial world has been slow to embrace the ion trap platform because it is often viewed as unmanufacturable, and composed of technologies foreign to the chip industry. Atomic physics experiments typically consist of a complex assembly of tabletop optical components, typically assembled individually by hand. While providing ultimate flexibility in the layout of the experiments, such an approach is subject to drift with temperature changes and often requires frequent adjustments to

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1Section is adapted from [SIJ⁺21] where it first appears some parts are subject to © 2021 IEEE
keep the optical alignment optimized.

Our goal is to transform an ion trapping experiment into a reliable quantum computer. The setup must be broken down into functional blocks, each block being an integrated module, to dramatically improve the system stability. In this work, we describe a novel system design approach to creating the UHV environment for the trapped ions, and the stable optical system to deliver the necessary laser light to manipulate the ion qubits.

The centerpiece of our approach is the design, construction, and testing of a compact UHV assembly for traps operating at cryogenic temperatures [WSGS13, BPC+15, SJP+16, SDF+11, Hud12]. The UHV assembly consists of a copper lid attached to a 100-pin ceramic pin grid array (CPGA) package containing the trap. We refer to this assembly as the ‘trap cryopackage’. The trap cryopackage is cooled to cryogenic temperatures to achieve the pressure levels needed for stable ion trap operation [ASA+09, BKM16]. The trap cryopackage is designed to minimize the volume of the UHV environment. It only contains the components that are absolutely necessary for trapping: the ion trap, the atomic source, and activated carbon getter to achieve UHV at cryogenic temperatures.

We explore two design iterations of the cryopackage: the first being a fully sealed cryopackage which is assembled within a large packaging chamber; the second being a differentially pumped open package in which we take advantage of a controlled leak to achieve exceptional pressures. The two design iterations had unique experimental setups. The first generation was more exploratory in design. Ultimately this first generation never achieved trapping; however important characterizations of ablation loading were performed. Many lessons were learned from this setup which informed the design of the second generation platform. In this new platform we achieved trapping in a fully sealed UHV package delivered by ColdQuanta. We decided to
move to a differentially pumped open package assembled at Duke from which we report most experimental results.

Our setup also features a compact optical layout. Each optical function is designed onto a customized breadboard referred to as an ‘optical block’. The optical elements are precision mounted to a pre-designed reference point. The breadboard can be temperature stabilized to eliminate optical misalignment due to thermal drifts. This design philosophy can be extended to future atomic physics experiments and ion trap quantum computers.

1.2 Dissertation Overview

In Chapter 2 I review some basics of the ion trap as a a physical platform and explore how the ytterbium qubit is realized. In Chapter 3 I briefly overview how we manipulate qubit information with microwave and optical sources. In Chapter 4 I overview in detail the design of our compact cryogenic ion trap system including the fabrication details for the trap cryopackage. In Appendix A the recipe for our trap cryopackage is provided. In Chapter 5 I feature the characterization of the heating rates and vacuum performance of our cryopackage system including results from multiple installed traps. We converge on exceptional common mode heating rates in the 10 quanta/s regime.

Chapter 6 features the characterization of single and two qubit gate performance in our system. Here we demonstrate exceptional performance for single qubit optical coherence (527 ms) using counterpropagating individual optically addressed beams. In order to characterize the performance of the two-qubit frequency modulated (FM) Mölmer–Sørensen gates in our system, we conduct a survey of a diverse population of gates. We sample 37 robust [LLF+18] and filter function robust [KWF+22] FM gates
with gate times between 120 \(\mu s\) and 450 \(\mu s\). This gate survey provides insight into the unique nature of the motional dephasing process which limits our average gate fidelity to 98%. Our conclusions are backed up by quantum Monte Carlo master equation simulations performed of the entire gate survey on the Duke computational cluster. The ion experiences low frequency coherent noise which makes parity error predominant. The coupling mechanism of this low frequency noise has implications for the design of cryogenic ion traps and grounding of general ion trap control electronics.

In Appendix C I overview the design of our first generation prototype system which featured a fully sealed trap cryopackage. We explore the characterization of isotope selective ablation loading in the weak regime in this section and I talk about the system design faults which informed the next generation design.

1.3 Summary of My Contribution

Such a large and complex experiment features contributions from many people over time. Here I try summarize my unique contributions chapter by chapter.

Chapters 2 and 3 features a summary in my words of basic issues, at the physical level, which I think may be helpful for students. These protocols are not my original contribution.

In Chapter 4 the system sample chamber internal designs were done by Geert Vrijsen and Volkan Inlek. However over the years I have overseen the disassembly, modification, and reassembly all of these components myself. I designed the optomechanical setup (large underlying structure featuring breadboards) and the initial CW optical block module for demonstrating trapping. Volkan Inlek designed the ablation setup. I designed the optomechanical realization of the Microelectromechanical systems (MEMs) individual addressing beam using the optical block approach. The
MEMs chips and the fundamental optical design are the result of prior work in the group. Volkan Inlek made the original designs of the Raman upstream and CW upstream plates. Gloria Jia and Junki Kim conducted the vibration characterization. Volkan Inlek designed and Gloria Jia implemented and maintained the PDH transfer cavity locking system. Gloria Jia designed the global beam plate and the ion imaging plate.

While not discussed in this document extensively, I have worked to create the table radio frequency (RF) delivery and power supply infrastructure. I managed the upkeep of the experimental setup through the years. The original cryopackage lid design was done by Geert Vrijsen in conjunction with ColdQuanta engineers. The current Peregrine cryopackage lid and cold finger sample mount was designed by me. The current assembly recipe and installation procedure while influenced heavily by Volkan Inlek was codified and organized in its current form by me. Chamber preparation, modification, and trap installations were lead by me from late 2019. The traps were fabricated and packaged at Sandia.

In Chapter 5 the four rod trap was designed by Volkan Inlek and the fully sealed cryopackage was delivered by ColdQuanta. The data presented on heating rates and zig-zag chain pressure measurements were collected by me. The data presented on double charging as a function of RF voltage was collected by Volkan Inlek.

Chapter 6 features GST and qubit coherence data whose implementation and collection I led. The two qubit gate calibration procedures was heavily influenced by mentorship from Ye Wang and collaboration with Gloria Jia. The particular pulse code used to run the gates was debugged and implemented by me with the valuable discussions with Gloria Jia and Ye Wang. The two qubit gate survey data was collected by me. Ke Sun collected four of the f-robust frequency modulated gate datasets and also led the CPMG data collection shown in this chapter. The
simulation package (utilizing QuTiP) used to understand our data was written by Mingyu Kang; however, the particular code (utilizing this package) for the simulation was written by me informed by valuable discussions with Mingyu.

In Appendix C the prototype system design is presented. The prototype package was designed and tested by Kai Hudek and Byeong-Hyeon Ahn. I designed the optomechanical and optical for the experimental ion trap system. Initial trapping attempts were done in conjunction with Byeong-Hyeon Ahn. Later I worked with Geert Vrijsen and he advised me on the implementation of ablation characterization. We both worked in conjunction to develop the time-resolved spectroscopy techniques to estimate the trappable atoms/pulse; this work was later tested on a UHV system (which could trap ions) to verify the method’s efficacy [VAS+19].
Chapter 2

Modern Ion Trapping Basics

2.1 Basic Paul Trap

The Paul trap was invented in 1953 with a paper published by Wolfgang Paul [Pau90] and he partly won the Nobel prize for this insight in 1989. The technology has seen wide use in mass spectrometers, atomic spectroscopy and atomic clock technologies [Gho95].

Charged particles can be dynamically trapped in 2D via a quadruple configured oscillating potential. To understand the basics of ion trap operation we review a derivation presented in [Gho95].

A quadrupole field can be written:

\[ \phi = \phi_0 \frac{2}{r^2} \left( \lambda x^2 + \sigma y^2 + \gamma z^2 \right) \]  

(2.1)

We must satisfy the relation in Laplace’s equation \( \nabla^2 \phi = 0 \) and so we can operate on 2.1 with \( \nabla^2 \) thus:

\[ \frac{\phi_0}{2r^2} (2\lambda + 2\sigma + 2\gamma) = 0 \]  

(2.2)

If there is not an applied field then \( \phi_0 = 0 \) and Laplace’s equation is satisfied, but this solution is trivial. If we consider only the 2D plane and seek trapping there we can set \( \gamma = 0 \) (the z-coordinate parameter) and then set \( \lambda = -\sigma = 1 \). Making the
potential in the x-y plane simply:

\[ \phi(x, y) = \frac{\phi_0}{2r_0^2} (x^2 - y^2) \]  

(2.3)

An idealized case that makes the analytical solution easy is to create four ‘rods’ whose center facing boundary has a hyperbolic geometry held at some potential \( \pm \phi_0/2 \). The Paul trap is a dynamic trapping so it only becomes interesting when driven by time-dependent potentials. We then wire together each of the two opposing rods, and drive them with a time dependent potential \( \pm (U - V \cos(\Omega t)) \) as seen in Figure 2.1 where \( U \) is a direct current (DC) voltage and \( V \cos(\Omega t) \) is a RF drive. This can create an effective potential:

\[ \phi(x, y, t) = (U - V \cos(\Omega t)) \frac{x^2 - y^2}{2r_0^2} \]  

(2.4)

In reality such an effective potential is realized by driving one set of opposing rods with an RF source, while grounding the other set of rails. They can be segmented and the z-potential can be closed off with an approximately quadratic DC potential. This can also be achieved with end-caps held at DC. In either case the opposite potential to the driven RF is induced in the grounded rails at RF frequencies and in the RF rails the opposite DC potential is induced at DC realizing the potential of equation 2.4.

The electric field at every point can be calculated as \( \vec{E} = -\nabla \phi \) and thus the force experienced by a charged particle in the field is \( \vec{F} = q\vec{E} \). For a singly ionized species \( q = e \), the electron charge. To write down the equations of motion we can simply calculate \(-e[\partial \phi / \partial x, \partial \phi / \partial y, \partial \phi / \partial z] = m[\frac{d^2x}{dt^2}, \frac{d^2y}{dt^2}, \frac{d^2z}{dt^2}]\). Important properties of the Paul trap can be explored by ignoring the ‘trap axis’ and considering only the equations of motion.
Figure 2.1: A schematic representation of a simplified Paul trap geometry. In reality two sets of rods are grounded while RF is driven on the other two. This situation is more commonly used for solving the problem and they create an equivalent boundary value problem and potential.

![Figure 2.1](image)

for an ion in the x-y plane.

\[
\frac{d^2 x}{dt^2} + \frac{e}{mr_0^2}(U - V \cos(\Omega t))x = 0
\]

(2.5)

\[
\frac{d^2 y}{dt^2} - \frac{e}{mr_0^2}(U - V \cos(\Omega t))y = 0
\]

(2.6)

In order to get the equation into a standard form and aid analysis we define two new parameters \( a \) and \( q \) as well as substitute the argument of the cosine function with a dimensionless parameter \( \xi \):

\[
a = \frac{4eU}{mr_0^2\Omega^2}, \quad q = \frac{2eV}{mr_0^2\Omega^2}, \quad \Omega t = 2\xi
\]

(2.7)

Solving for and then substituting the \( U \) and \( V \) parameters and then working out the relationship between the differentials to obtain \( \frac{d^2 x}{dt^2} = \frac{\Omega^2}{4} \frac{d^2 x}{d\xi^2} \) and \( \frac{d^2 y}{dt^2} = \frac{\Omega^2}{4} \frac{d^2 y}{d\xi^2} \) and then
simplifying everything we arrive at the standard form of a Mathieu equation [Tho48]:

\[
\frac{d^2x}{d\xi^2} + (a - 2q \cos(2\xi))x = 0 \quad (2.8)
\]

\[
\frac{d^2y}{d\xi^2} - (a - 2q \cos(2\xi))y = 0 \quad (2.9)
\]

The dimensionless parameters \(a\) and \(q\) are sometimes called ‘trapping parameters’. In 2.7 \(r_0\) depends on the field (and therefore trap) geometry and is approximately the distance between the trapping location and the electrodes. The parameters \(a\) and \(q\) can be thought of as the normalized trapping strengths of the DC and RF potentials respectively. The Mathieu equation has two types of behaviors which depend upon the values of the trapping parameters, a periodic and unstable solution, and a periodic stable solution [Mar97]. Frozen in time and considered in the x-y plane this arrangement makes a saddle shaped hyperbolic potential in which the ion resides (in motion) in the center.

A periodic and stable solution will result in a bounded motion characterized by three ‘secular’ frequencies, sometimes called the ‘trap’ frequencies, or motional frequencies. The trap frequencies have a fundamental which is smaller than the driving RF frequency but all higher harmonics can be present when highly excited. Intuitively the trapping mechanism can be imagined by viewing the saddle shape in Figure 2.2, as the RF advances in phase with time, the potential flips up and down between the two saddle configurations. If the rate of flipping is fast enough that the ion cannot escape between the two orthogonal phases then what is called a pseudopotential null is formed. These two orthogonal states are shown in Figure 2.3. If we model the RF with a cosine function then at 0° phase a force is imparted which will push inwards from the top and bottom, while after a 180° of phase advancement, the opposite force will be applied, pushing inward horizontally. Swapping between
these two phases rapidly enough that the ion cannot escape will cause it to stay trapped in a closed orbit.

Figure 2.2: A visualization of the saddle shaped hyperbolic potential that the ion experiences in the x-y plane of the trap. In every RF period this saddle will rotate around the central point. If the mass, charge, RF frequency, and the RF/DC amplitudes are situated in the stability zone then the ion will remain localized in a closed oscillation.

An analogous situation that provides a very intuitive view would be to imagine a ball bearing sitting at the center of a spinning saddle shaped object, and the saddle must spin fast enough such that the ball bearing cannot escape its position horizontally, and the ball is kept in place vertically by gravity pushing downward and the normal force of the surface pushing back up. In this situation the action of gravity and the surface would be replicating the action of a DC potential well while the spinning saddle is the dynamic trapping force keeping it in place horizontally. However in this situation the ball would never be perfectly still but rather could undergo a constant small amplitude motion as it interacts with the saddle. If the ball were too light it may be knocked out; if the saddle was spinning too slowly, it wouldn’t maintain the ball and it will roll out.

If the ratio of $\frac{a}{q}$ is correct then the motion of the ion in the Paul trap will be
Figure 2.3: A schematic representation of a simplified Paul trap geometry. The color gradient represents the hyperbolic potential, the lines are electric field lines which flow from positive to negative charge.

small enough such that it will have a bounded orbit. For a given $U$ and $V$ the values of $a$ and $q$ are set by the charge to mass ratio $\frac{e}{m}$ and if we consider only singly ionized species then this becomes a mass filter with a filtering cutoff which can be tuned by adjusting the strengths of $U$ or $V$. In general the entire stability landscape for large values of $a$ and $q$ is complicated and can be seen in Figures 2.2 and 2.3 of [Gho95], but the relevant operational region for most applications, especially ion trap quantum computers, is the regime of small trap parameters near the origin. In Figure 2.4 the region near the origin is plotted in a reproduction of a common figure present both in [Pau90, Gho95] and many other references. We can visualize the trap parameter space as being divided into regions where motion in either axis can be stable or unstable. We can envision a line whose slope is established by the RF and DC potential strengths (shown in red in Figure 2.4), species of different masses will fall on this line and the ones with a mass in the stable region will remain trapped, while the ones which are lighter will be ejected. It is this action which made the Paul trap technology revolutionize mass spectrometers [Gho95, Mar97]. If we imagine
ramping the RF amplitude $V$ down in time this will make the slope of the red line increase. Eventually lighter mass ions will be ejected as their secular motion becomes large and they are ejected the trap. These ions can be detected in sequence during this RF amplitude ramping and thus we have a mass spectrometer.

![Graph showing the key zone of the Paul trap stability diagram.](image)

**Figure 2.4:** The key zone of the Paul trap stability diagram for most applications in mass spectrometry and ion-based quantum technologies. The black lines are borers of the trap stability regions, while the red line is an operation line created by the relative strengths of the DC ($U$) and RF ($V$) potentials.

In general we can numerically solve the Mathieu equations for the motion of an ion, however for the parameters that we use in ion trap quantum computers we can make a few simplifying assumptions and arrive at a potential well model called the pseudopotential approximation [CBB+05, Gho95]. We can think of the motion of a stably trapped ion as a sum of the secular motion, ($X$), at the trap frequency and ‘micromotion’, ($\delta$), at the RF frequency $\Omega$. The $x$-axis motion would then be their
sum $x = X + \delta$, and this is true for all three axes. A detailed derivation of the trap depth making the pseudopotential approximation for the z-axis of a 3D Paul trap can be found in [Gho95] section 2.5. One should keep in mind that this solution relies on making the approximation that the micromotion amplitude is much smaller than the secular motion ($\delta \ll X$) while the micromotion is at a much faster timescale ($\frac{d\delta}{dt} \gg \frac{dX}{dt}$). Furthermore we make the assumption that $a \ll q$ or that the normalized DC potential is much smaller than the RF potential. For our purposes it suffices to use the results of [Deh68, Gho95] as presented by [CBB+05] to give a basic idea for how this approximate solution behaves. When the pseudopotential approximation is applied the ion can be thought of as under the influence of a ‘ponderomotive’ potential.

$$U_P(\vec{r}^*) = \frac{e^2}{2m\Omega^2}(E^2(\vec{r}^*))$$ (2.10)

The $E(\vec{r}^*)$ depends upon the particular trap geometry and for a standard geometry as we have used above we can find:

$$U_P(\vec{r}^*) = \frac{e^2V^2}{4m\Omega^2r_0^2}(x^2 + y^2)$$ (2.11)

Under these simplifying assumptions which we seek to realize for all ion trapping systems used for quantum computation the motion of the ion takes on the form [CBB+05]:

$$x(t) \approx X_0 \cos(\omega t)[1 + \frac{q}{2}\sin(\Omega t)]$$ (2.12)

Using this pseudopotential model one can solve for the trap frequencies as a function of the system parameters which is shown below [MHS+04].

$$\omega_{P,x/y} = \frac{eV}{\sqrt{2m\Omega r_0^2}} = \sqrt{\frac{eVq}{4mr_0^2}}$$ (2.13)
One should remember that Equation 2.13 is this value for the infinitely long hyperbolic potential model. In real situations this relationship is more complicated and can be solved for numerically, and there will be a different relationship for the axial trap frequency which is related to the DC confining potential. Spatially dependent shape of the pseudopotential will be important for the trap geometries encountered which we will encounter as we try to scale the ion trap technology and make it capable of the complex operations needed for large-scale quantum computation.

2.2 Surface Paul Trap

In the previous section we used the basic model of a Paul trap to understand the properties of dynamic trapping in the four rod geometry. There are many ways to realize ion traps with this four-rod geometry in what we call ‘macroscopic’ traps. The most basic design is four rods, two grounded, two in which RF is driven with some needle shaped end-caps aligned to the trap axis. Another level of sophistication can be added by dividing the RF-grounded elements into segmented DC electrodes for applying arbitrary voltages along the length of the trap. For quantum computing applications it is desirable to have large trap frequencies, while keeping within the trap stability regime of Figure 2.4. If we examine Equation 2.13 we see that \( \omega \propto \frac{V}{r_0^2} \), thus if the ion-electrode distance is large then the RF voltage needed to maintain the same \( \omega \) grows faster. It is desirable to minimize \( r_0 \) so that the RF voltage requirement is not severe [DCL+15].

In practice it is difficult to use a rod geometry while having a small \( r_0 \). In response to this requirement blade geometries have become popular [SKHG+03]. The blade geometry is more amenable to precision construction at the scales needed [DCL+15]. These designs have been a workhorse where many quantum protocols and noisy
intermediate-scale quantum (NISQ) scale quantum simulations have been performed but they are not suitable for the more complicated protocols needed to demonstrate an error corrected logical qubit or to scale the system beyond that in either a modular architecture [MK13, KMK+11a] or a QCCD architecture [KMW02].

These protocols require the ability to shuttle, merge, and reorder the ions, and potentially the integration of microcavity optics will become necessary in the long-term [KMK11b, VRSS+17]. All of these protocols require many micro-scale DC control electrodes which cannot be accurately manufactured in the blade geometry. Also potentially important are unique trap geometries such as junctions where the ions can be shuttled in such a way that they turn corners and can be moved into distant trapping regions. Such protocols are also made feasible at large scale only with a standard uniformly manufactured device which is very predictable.

Only by leveraging microfabrication technologies can this level of precision and complexity be achieved. Microfabrication could be used to miniaturize a 3D quadrupole structure using multilayer designs. However there are many practical issues with this [CBB+05]. By mapping the Paul trap into a 2D plane we can achieve a near-hyperbolic potential as well. In order to see how this mapping works one must take note of the electric field lines in Figure 2.3 which flow from positive to negative charge and then compare them with the field lines drawn in Figure 2.5. Despite the 2D geometry we establish a region with equivalent field line geometry to the 3D case but with the ion hovering above the trap. While the total potential is very different near to the ion location, it will have a dominant quadrupole component and be approximately hyperbolic leading to a near-quadratic potential well in the pseudopotential approximation.
Figure 2.5: A schematic diagram of the surface Paul trap projection that most closely resembles our Sandia surface traps. Other mappings in a 2D plane as possible including a four electrode design in which every other electrode is RF/DC. This geometry is very useful as the inner grounded regions can be segmented DC electrodes very close to the ion for fine control. An arrow in green shows the \( r_0 \) equivalent for surface traps often called in this context the trap height. It is not always the case that the characteristic distance \( r_0 \) is equivalent to the ion height.

2.2.1 Surface Trap Design Features

Many diverse designs for surface traps are possible. Sometimes they are simply single layer metal-on-insulator traps with a rectangular geometry, but also more complicated multilayer routing can be used to achieve interesting trap geometries. In this work we have used both types of traps. However most of the scientific results were achieved in a trap fabricated on a CMOS compatible silicon process into a ‘bow-tie’ geometry. These types of traps can have 100 DC electrodes and they can be routed into such a narrow region by leveraging multilayer interconnect routing technologies of a silicon process. Whereas to get so many control electrodes in a single layer design necessarily leads to a large fan-out of the routing electrodes.

If we look at Figure 2.6 we will see an exaggerated schematic of a bow-tie shaped
Figure 2.6: A schematic representation of a surface Paul trap. This drawing is not to scale and particularly the size of the ions are exaggerated for viewing. In reality the electrode sizes could be 70 µm while a 7 ion chain would be 35 µm. The electrode sizes would depend strongly upon the trap design. The purpose of this drawing is to introduce the coordinate system commonly used for surface traps as well as labels of common features of the Sandia surface traps which we use. Compared to macroscopic traps which often use the z-axis for the trap axis, in surface traps we use the z-axis, as the vertical normal to the trap surface. This surface trap z-axis will also be the imaging axis collecting ion fluorescence. The Paul traps radial confinement axes then become the z and x axis. The y-axis is the trap axis where linear ion crystals will form along which the DC potential is established. The traps we have used in our studies have a bow-tie geometry with a central ‘isthmus’ which is as narrow as possible to maximize the available numerical aperture (NA) for individually addressed beams. There are a set of segmented inner DCs which in some traps have a slot in between for optical access through the bottom of the trap. The inner DC electrodes are also used commonly for the main confining potential while the large outer DC electrodes can be used for global micro-motion compensation voltages as well as adding rotations to the trapping axis.
surface Paul trap similar to the ones we use from Sandia. This is provided to fami-
iliarize the reader with the coordinate system we will use as well as some surface trap
terminology. In surface traps the coordinate system is customarily changed so that
the z-axis becomes the axis normal to the surface of the trap which will also be the
optic axis of the lens used for imaging ion fluorescence. The radial confinement axes
which replace x/y from the four rod geometry now become x/z where the x-axis is
the lateral trap dimension. The y-axis is then the trap axis along which ion chains
will form and the DC confinement potential will be established.

The advantage of the bow-tie shaped design is that numerous bond pads can
be exposed at the far edges which are wide, and they can all be routed towards
a more compact center. The more narrow the center trapping region can be, the
greater the available numerical aperture (NA) for optical access will be, and more
tightly focused beams can be brought in. It is especially important to minimize
scattering from the beam interacting with the trap when ultraviolet (UV) lasers are
used. When UV beams clip at the trap edge it can cause charging which effects
the trapping potential. More information about the need for tightly focused UV
beams can be found in Chapter 3. If one seeks to realize a logical qubit, mid-circuit
measurement with clean 369.5 nm beam is required and errant light will disturb the
qubits in other trapping regions. In Figure 2.7 we see a helpful plot inspired by
the HOA 2.0 manual [Mau16a] which has been included as a helpful reference for
optical designs and understanding. For a beam horizontal to the trap surface there
exists an optical beam-waist for maximum clearance at the edge of the trap which
will minimize scattering. The benefits of the HOA 2.0 trap are clear there is both a
larger value of maximal clearance and this beam can be more tightly focused.

In Figure 2.6 we see some salient features of surface traps we use such as sets of
‘inner DC’ electrodes which provide for fine control of the axial potential, ‘outer DC’
The clearance for two trap geometries once with the HOA 2.0 trap which has a 600 µm lateral distance between the trapping location and the edge of the trap. Another the EPICS 2 trap with a 1.635 mm distance between the trapping location and the trap edge. The clearance of a 355 nm beam focused at the ion evaluated at the trap edge normalized by the beam waist at that point. The results for 90° and 45° to the trap axis are shown.

Electrodes which are larger, sometimes not segmented, which help with micromotion compensation and can provide trap rotation. The RF is routed in one large electrode which surrounds the inner DC electrodes and serves as our two RF rods as we consider a 2D slice as in Figure 2.5. In this image the size of the ion chain is exaggerated greatly for visual schematic effect, a typical DC electrode size could be 70 µm while a 7 ion chain would span 35 µm in a typical potential. We also see a tightly focused beam brought into the plane of the surface trap, often the 355 nm Raman beams will be brought in mutually orthogonal to the trap axis (y-axis) and the imaging axis (z-axis). Our continuous wave (CW) cooling and repump beams will be brought in at 45° to the trap axis so that there will be projections along all three motional axes of the ions so that we can cool all three axes simultaneously with one beam.

Seen in Figure 2.8 we have on the right two examples of single layer gold on fused
silica traps which were used in the early stages of this project, and on the left we have two silicon-based Sandia traps used to collect the qubit data presented in this work. The HOA trap features Y-shaped junctions at either end for storing and reordering the ion chain in the center. It also features a long slot in the center for even greater optical access in the z-direction for individual addressing beams. This can also be used for loading the trap from below to prevent sputtering of material onto the trap. The Phoenix/Peregrine series is more recent and they are upgraded and somewhat simplified versions of the HOA 2.0 trap. The Phoenix features a slot like the HOA 2.0, the Peregrine (used in this work) excludes the slot feature and only has inner DC electrodes beneath the ions. These traps feature improved RF design with stronger grounds to lower RF losses, as well as did away with the Y-junctions that most users
were not utilizing.

The two traps shown on the right were produced originally for a cavity QED integration project and are a single layer of gold deposited upon fused silica with a dielectric mirror coating. It is more difficult to design tight geometry in the trapping region with a single layer trap because the many control electrodes need to be routed to the center and this leads to fan-out.

### 2.2.2 The Surface Trap Pseudopotential

![Figure 2.9](image.png)

Figure 2.9: Plots of the pseudopotential of an HOA 2.0 trap from our simulation software. (a) an example of a pseudopotential which has a trap axis rotation near 0° (b) The DC solution can be established to rotate the principal axes of the trap to 43°. This is useful for being able to Doppler cool both radial modes of motional effectively and demonstrates the flexibility of surface traps.

In order to produce real trapping fields for many diverse uses, finite element
simulation is needed. The trap simulations are provided by Sandia, and they use finite element simulations with 1 V applied to each electrode. In Mathematica the results of this can be scaled and added together to plot the pseudopotentials for a particular DC voltage array. We use this in order to generate good DC solutions for the application we desire, and to characterize the trap frequencies we should expect from a particular RF voltage amplitude. In Figure 2.9 we see a plotted simulation of the RF pseudopotential for an HOA 2.0 trap with a 250 V RF amplitude in the neighborhood of the null point. Figure 2.9(a) we see a standard radial null with no rotation, while in Figure 2.9(b) we see a solution where the outer DC electrodes are used to rotate the normal modes of ion motion and thus the trap axes to a 43° angle. The ability to rotate the trap axes allows one to tune the coupling of the Raman and cooling beams arbitrarily just by adjusting the trap voltages.

In Figure 2.10 we see the trapping pseudopotential plotted along the three principal axes. For the RF radial confinement portion we have Figure 2.10(a) and Figure 2.10(b) which are the lateral and vertical directions respectively. The vertical trapping dimension for surface traps is normally weaker. In Figure 2.10(c) we see one example of a DC potential. The DC solution can take many arbitrary forms as needed. Harmonic potentials will begin to have very unequal ion spacing for chains. A flatter potential can be established with near-equal spacing, double-well potentials can be used for hopping measurements, and with the correct DC control we can play DC voltage sequences which will shuttle the ions and merge and split chains, as well as potentially reorder them. We can see that the trap depths for a surface trap are on the order of 0.1 eV.
Figure 2.10: Plots of the pseudopotential wells as a function of position for an HOA 2.0 along the three principle axes. (a) The pseudopotential along the x-axis (or lateral) (b) the pseudopotential along the z-axis (or vertical) (c) the pseudopotential along the y-axis (or axial) the radial trapping is weakest along the vertical direction and the DC potential can take on many arbitrary forms depending upon the need due to the flexibility of surface traps.

2.2.3 Challenges Facing Surface Traps

While the repeatable trap fabrication, small feature size, and complex feature set of micro-fabricated surface traps are all necessary for advanced quantum computing protocols, there are downsides and challenges facing surface traps. They have a low trap depth compared to macroscopic traps which can handle higher RF voltages. This low trap depth makes them more susceptible to collisions, leading to lower single ion and chain lifetimes. Additionally there is the issue of anomalous heating from electric field noise on metal surfaces which has been studied [DOS+06] and found to increase the closer the ion sits to the trap surface [BKW18]. Because macroscopic traps can have much larger ion-metal distances they are much less subject to this concern. Furthermore there is the issue of trap charging due to UV trapping and Raman lasers necessary for many species. Early in the production of surface traps the community was uncertain about whether high fidelity quantum gates could be achieved on them despite these challenges. The work of Mount et al. [Mou15, MKC+15] proved that
high fidelity single qubit operations could be achieved. High fidelity state detection fidelity has also been achieved in many systems [CCV+19] as well as high fidelity two-qubit Mølmer–Sørensen (MS) gates [LBBK+16, WCF+20].

Still remaining as a challenge for the platform is the retention of long chains without reordering, and the reduction of common mode heating rates which can degrade achievable circuit depth. There are many indications that moving to cryogenic operation can greatly reduce common mode heating effects [LGA+08, NLK+14, CS14, Nie15]. Furthermore moving to cryogenic operation will make UHV more easily achievable and make the collisional energies lower, mitigating the negative effects of low trap depth. Moving to a cryogenic platform can mitigate the shortcomings of surface ion traps making them viable for deep quantum circuits with long chains and a much more attractive platform than macroscopic ion traps.

2.2.4 High Voltage RF Delivery

In order to achieve the high trap frequencies desired for quantum application using the motional modes, we require high voltage RF for ytterbium. This usually means the 30-55 MHz range at a 200-350 V amplitude. The most common way to achieve this is with the helical resonator [SSWH12]. Lumped element tank circuits which are simple RLC resonators can also be used such as [GNK+12].

2.3 Cooling and Trapping the Ytterbium Ion

The ion trapping field has worked with many different species, each with unique benefits and drawbacks. We have chosen to work with ytterbium ions. Ytterbium atoms and ions have been used for atomic clocks and frequency standards [CEB+91, HSP+13, JPS+12, SJP+16]. Ytterbium ions have a convenient closed cycle cooling
transition which can be driven by diode lasers. The $^{174}\text{Yb}^+$ isotope is commonly used for the testing of ion trapping and cooling, it is the most abundant isotope in natural ytterbium and can be most easily trapped and cooled. The $^{171}\text{Yb}^+$ isotope is used for the qubit due to the presence of a hyperfine ground state splitting of 12.6 GHz which is first-order magnetic field insensitive. This robustness against magnetic field noise makes it simple to attain long (second scale) qubit coherence times. The timescales of high fidelity state initialization, detection, and qubit manipulation are orders of magnitude less than the coherence time[OYM⁺07]. Additionally off-the-shelf mode-locked pulse lasers can be used for individually addressed qubit operations [HMM⁺10, IVH⁺14].

2.3.1 Neutral Ytterbium Ionization

Trapping ytterbium requires a neutral ytterbium flux. Traditionally most systems have used a thermally evaporated beam coming from an atomic ‘oven’. Usually this oven is a high aspect ratio metal tube with a ytterbium sample at the bottom. It is heated by passing a large current through the tube. And ytterbium will then evaporate and travel towards the trapping location. This method is difficult to integrate into cryogenic systems but this has been done [VWB⁺13a]. It requires the careful design of an oven which is heat-sunk at a higher temperature stage. This method is bulky and undesirable for use in this compact package. The thermal budget of our cryostat will be tight. In the process of working with the compact trap cryopackage we studied ablation loading extensively and learned how to calibrate an isotope selective ablation loading in the near-threshold regime. Details on our ablation process are given in Appendix C Section C.6.

With a flux of neutral ytterbium flowing through the trapping region we will shine
Figure 2.11: Plot of the ionization scheme for neutral ytterbium. We require 399 nm light to bring an outer electron from the $^1S_0$ to the $^1P_1$ state, where the electron can be stripped to the continuum with photons with wavelengths below 394.1 nm. We usually use intense light from our 355 nm pulse laser but have also used a CW 391 nm diode with 1-3 mW of power for the same task. The data on this plot is originally from [KHTY99, SMY05] as collected by [OMM+09].

ionization lasers according to the scheme of Figure 2.11. This involves shining a 399 nm laser stabilized to within 5 MHz and tuned to the isotope line that we wish to trap. The 399 nm light causes the outer electron to make the $^1S_0$ to $^1P_1$ transition. Once in this state another photon with a wavelength less than 394.1 nm is then needed to ionize an electron to the continuum. For this work we have used both 391 nm and 355 nm light for the second photon. When using a thermal oven it is often sufficient to rely upon the 369.5 nm cooling beam for this purpose, but for ablation loading we work to maximize photoionization efficiency. We have used a free running 391 nm laser diode delivering more than a 1 mW focused to the trapping location. Our 355 nm Raman laser is optimal for loading efficiency.
2.3.2 Doppler Cooling

Once the atom sheds its outer electron in the trapping region it immediately becomes trapped by the Paul trap’s potential and begins to scatter photons from the cooling laser. Doppler cooling relies on the force imparted by a resonant laser interacting with the ion. In this section a thorough explanation of Doppler cooling in ion traps is provided with a sense towards pedagogy. There is a limit to the effective temperature that can be reached via Doppler cooling. A few preliminaries are necessary for understanding the process and this discussion is based on the formalism present in [MvdS99] and some insights from [OMM+09]. Atomic transitions have a natural linewidth which can be broadened by resonant beam power, the Doppler shift of the light as seen by the atom, as well as pressure and other mechanisms [MvdS99], the frequency dependence of the scattering rate goes as

\[ \gamma_p = \gamma \rho_{ee} = \frac{s_0 \gamma}{1 + s_0 + \left( \frac{2A}{\gamma} \right)} \]

(2.14)

Where \( s_0 \equiv \frac{I}{I_s} \) and \( I_s = \frac{\pi \hbar c}{3\lambda^3 \tau} = \frac{\pi \hbar c \gamma}{3\lambda^3} \), and we define, as in Figure 2.13, the lifetime as \( \tau = \frac{1}{\gamma} \). In the denominator of \( I_s \) we could include the branching ratio as in [OMM+09] but this is a small correction. The value \( \rho_{ee} \) is the excited state density matrix element which for a pure state is the probability of the atom being excited. For large powers this can only ever saturate to \( \gamma/2 \) because the process is limited by spontaneous emission. The exact results of this come from solving the optical Bloch equations (OBE) in steady state. A good discussion of this exists in [MvdS99].

The fact that the scattering rate is detuning dependent is key for the mechanism of Doppler cooling. Doppler shifts seen due to atomic velocity can be thought of as effective detunings shifting the entire power broadened spectrum for that ion. If we look at Figure 2.12 we can see that for a laser source tuned red of the Doppler-free
spectral peak ion moving towards the beam, the velocity will shift the $\gamma_p$ spectrum to raise the scattering rate for ions moving towards the beam. For ions moving away from the beam, the opposite shift of the spectrum will lower the scattering rate further. For a laser source tuned blue of the Doppler-free spectral peak the opposite effect is observed. For each scattering event, the momentum of the photons will be transferred to the ion. For red detuned light there are more scattering events for atoms moving towards the beam and thus the net velocity dependent force acts in opposition to the ion’s momentum. This will narrow the velocity distribution of an ensemble of atoms [MvdS99]. For blue detuned light, more scattering events occur for atoms moving in the same direction as the light and thus it will increase the ion’s momentum, leading to a broadening of the velocity distribution for an ensemble [MvdS99].

Force is exerted on atoms and ions via their interactions with light and in Metcalf [MvdS99] a derivation is presented for this force for an atom at rest applying
Ehrenfest’s theorem in the form $F = \langle \mathcal{F} \rangle = \frac{d}{dt} \langle p \rangle$. One can go back to the OBE and then solve for the detuning dependent damping coefficient given below as solved by [MvdS99] for a two-level system (atom/ion) in motion in a traveling plane-wave. There is a force imparted in the form of 2.15 [MvdS99]:

$$F_{\text{ion}} = F_0 - \beta v$$

(2.15)

The exact forms of the equations for $F_{\text{ion}}$ and $F_0$ can be found in [MvdS99] but lots of intuition can be found via a short discussion of Equation 2.15. The form of this equation is similar to that of an object in free fall i.e. a constant force with a velocity dependent damping, which can be thought of like a drag, in opposition. Metcalf notes that $F_0$ is identical to the force experienced by an ion at rest absorbing and spontaneously emitting a photon and has the form $F_{sp} = h\kappa \gamma \rho_{ee}$. Comparing this with equation 2.14 we can understand that this force will also saturate with optical power to a constant and is also proportional to the photon momentum. The damping coefficient $\beta$ is shown below as solved by [MvdS99]:

$$\beta = -\hbar k^2 \frac{4s_0(\delta/\gamma)}{(1 + s_0 + (2\delta/\gamma)^2)^2}$$

(2.16)

For the case of a standing wave and an ion in motion this constant force $F_0$ would cancel and we would be left with a pure damping which would simply slow that ion [MvdS99]. Often in textbooks Doppler cooling is talked about in this context where an ‘optical molasses’ is formed via counterpropagating beams sent in all three orthogonal axes. In an ion trap however we have a restoring force keeping the atom localized in the trap which is much larger than the size of $F_0$. Furthermore with complete control of the radial trap axes, we can bring in a cooling beam at a 45°
angle with respect to the trap axes. Thus a beam brought in along the trap surface can have a projection in all three principal axes of motion. The trap frequencies are of the order 1 MHz, while the scattering rate saturates to $\gamma/2 \approx 10$ MHz and so, thought of as a classical oscillator it will scatter several photons moving towards the beam and several moving away with a red-detuned cooling beam. The ion will slow down when moving towards the beam at a rate $\gamma/2$ due to the damping force $-\beta v$; when moving away from the beam the scattering rate will decrease. Many cycles of this will continue to rob the ion of momentum.

A mechanical analogy that is quite similar is to imagine a mass suspended in space by springs in all three orthogonal axes that begins with a lot of kinetic energy, but is subject to a velocity dependent frictional force. The oscillations will damp out and an equilibrium will be reached.

The point at which Doppler cooling saturates is called the Doppler limit. The minimal achievable temperature is denoted by $T_D$. In order to understand how this limit arises, we must consider both the quantum nature of the interaction and the recoil energy of the absorption and emission process [MvdS99]. The atomic momentum can only change in discrete values of the photon momentum, $\hbar k$. The kinetic energy scale of this momentum change is defined as the recoil energy $E_r = \frac{\hbar^2 k^2}{2m} = \hbar \omega_r$. For spontaneous emission this results in a loss of atomic kinetic energy in a random direction into the light field. Absorption energy is taken from the light field and given to the atomic kinetic energy. For each scattering event then there is a net loss of $2\hbar \omega_r$ [MvdS99]. This causes the recoil induced heating of the atom, which will in steady state, come into balance with the velocity dependent damping force which cools. This steady minimum state energy depends upon the detuning but for the optimal case where $\delta = \gamma/2$ this energy is $k_B T_D = \hbar \gamma/2$ [MvdS99, OMM+09]. This value is derived for the case of counterpropagating beams where there is a factor of
two increase in the damping parameter $\beta$ and a factor of two increase in the rate of recoil heating $2\gamma_p$. When balancing these two effects these factors of two should cancel and thus for ions locally trapped, the result should be the same. For our species this value is easily calculated as $470.3 \, \mu K$. Putting this into perspective of our motional quanta, if our radial modes have a frequency $2.0 \, MHz$ and we utilize the energy in a quantum harmonic oscillator, $\hbar \omega (n + \frac{1}{2})$, then we can solve for the useful minimum motional quanta achieved in our ion via doppler cooling as:

$$n = \gamma \frac{\gamma}{2\omega} - \frac{1}{2}$$  \hspace{1cm} (2.17)

For our species $\gamma = 2\pi (19.6 \, MHz)$ resulting in $n_D = 4.4$ quanta.

In practical terms being able to achieve Doppler cooling in a species means you need a closed loop transition with a line-width in a good regime for both laser locking and a reasonable Doppler temperature. We use a 369.5 nm laser locked to a 1 MHz linewidth, which is focused to a $w_0 = 15 \, \mu m$ with a power in excess of $1 \, \mu W$ which will be comfortably above $I_{sat} = 51.04 \, W/cm^2$. More details on the lock and the CW modulation plate will be given in Section 4, however simply we lock this laser 210 MHz red of the Ytterbium $^2S_{1/2}|F = 1\rangle$ to $^2P_{1/2}|F = 0\rangle$ transition split the beam into two paths, one for detection/optical pumping, the other for Doppler cooling. The Doppler cooling path travels through a 14.7 MHz EOM to apply sideband at this frequency and then goes through a shutter acousto-optic modulator (AOM) operated at 200 MHz in order to give about 10 MHz detuned light for cooling. In Figure 2.13(a) we can see that this detuned cooling light drives transitions between these states. Most of the time they drop back down to the $^2S_{1/2}|F = 0\rangle$ state and the scattering events lead to cooling.

Different mechanisms will trap the population away from the cooling loop. One is
Figure 2.13: A diagram explaining the Doppler cooling scheme for $^{171}$Yb$^+$ adopted from [OMM+09, OYM+07] laser driven transitions are shown in cyan while 369.5 nm emission is shown in purple, 935 nm emission is shown in red. (a) For Doppler cooling we detune approximately $\Gamma/2 \approx 10$ MHz in order to maximize the damping and reach the lowest temperature. When decaying from the $^2P_{1/2}|F=0\rangle$ state there is 0.5% population leakage into the $^2D_{3/2}|F=1\rangle$ states which we can bring back into the cooling process by applying 935 nm light tuned to the $^2D_{3/2}|F=1\rangle$ to $^2D[3/2]_{1/2}|F=0\rangle$ transition. Due to the hyper-fine splitting of the $D$ levels we apply 3.0695 GHz side-bands to the 935 nm laser to drive the $^2D_{3/2}|F=2\rangle$ to $^2D[3/2]_{1/2}|F=1\rangle$ transitions as well. (b) The $^2D_{3/2}|F=2\rangle$ states become populated by off resonant driving to the $^2P_{1/2}|F=1\rangle$ states, whose decay can add population to the $^2S_{1/2}|F=0\rangle$ or the $^2D_{3/2}|F=2\rangle$ states. Removing population from the $^2D_{3/2}|F=2\rangle$ by driving it to the $^2D[3/2]_{1/2}|F=1\rangle$ state can also end driving population which gets trapped in the $^2S_{1/2}|F=0\rangle$ states. In order to bring these states back into the cooling cycle, we apply modulation with a 14.748 GHz electro-optic modulator (EOM) to our cooling beam.

the $^2P_{1/2}|F=0\rangle$ to $^2D_{3/2}|F=1\rangle$ transition which has a 0.5% branching ratio [OYM+07]. This state has a much longer lifetime than the cooling transition and so even though the branching is small over many cooling cycles all population will end up there and the ion will be dark. We tune a 935 nm laser to the $^2D_{3/2}|F=1\rangle$ to $^2D[3/2]_{1/2}|F=0\rangle$ transition with a beam with greater than 1 mW of power. It will be far above sat-
uration. This repump beam keeps the ion in the cooling loop. One complication arises for the hyperfine splitting in $^{171}$Yb$^+$ which is that there are off-resonant transitions from $^2S_{1/2}|F = 1\rangle$ to $^2P_{1/2}|F = 1\rangle$ which can then populate the $^2S_{1/2}|F = 0\rangle$ or the $^2D_{3/2}|F = 2\rangle$ states leading to lost population here. In order to avoid this use, require the 14.7 GHz EOM to drive population out of $^2S_{1/2}|F = 0\rangle$ in the scheme shown in Figure 2.13(b). Additionally we require 3.0695 GHz side-bands on our 935 nm laser which we also drive with an EOM. This extra 935 nm sideband will drive the $^2D_{3/2}|F = 2\rangle$ to $^2D[3/2]_{1/2}|F = 1\rangle$ transition which again can bring the ion back to the cooling loop directly or via the 14.7 GHz sideband offset scheme of Figure 2.13. With these beams in place we can cyclically drive our Doppler cooling transition and thus cool trapped ytterbium to the regime of 5 quanta.

When the magnetic field strength is low the $^2S_{1/2}|F = 1\rangle$ states become degenerate and a coherent dark state can be formed for certain polarizations of cooling light [BB02, OMM$^+$09]. This is alleviated by applying a magnetic field which will split the Zeeman states by 8-12 MHz.

The final complication that arises during Doppler cooling are long-lived darkening events. The most common event is the collisional excitation of the ion into the $^2F_{7/2}$ state [OMM$^+$09, BST92, LHLBH89]. The susceptibility to this also depends upon 369.5 nm cooling power and detuning. Often a 638 nm laser can be used to pump out of this state, however, in our system we use 355 nm light for this purpose. More exotic long-lived dark ion events have been observed in our cryogenic system including doubly charged ytterbium (which cannot be recovered) and the formation of Yb$H^+$ molecules which can be dissociated by strong 355 nm light in much slower timescales than the $^2F_{7/2}$ states [HJOS20, SY95]. Both the double charging events and the formation of Yb$H^+$ in our trap will be discussed in more detail in Chapter 5.
2.4 Establishment of the Ytterbium Qubit

With the ion trapped and Doppler cooled, it is now localized, then we require protocols which will establish the qubit manifold thus allowing us to do quantum computations. We require a method to accurately initialize a population to the $|0\rangle$ state as well as some method of discriminating the zero and one state. We also need to ensure that during algorithms we don’t have leakage into other states outside of the qubit manifold.

2.4.1 Initialization to the Dark State

In Figure 2.14(a) we have the scheme for pumping the population into the $|0\rangle$ state. In this case we require another beam path which has been split off upstream of our cooling EOM and AOM then sent through a different EOMs. These add a 2.105 GHz sideband to the light. Furthermore this AOM is driven at 210 MHz and the laser is tuned so that with this AOM shift it is near resonance. For pumping we turn on the 2.105 GHz sidebands and this drives the transitions $^2S_{1/2}|F = 1\rangle$ to $^2P_{1/2}|F = 0\rangle$ and the $^2S_{1/2}|F = 1\rangle$ to $^2P_{1/2}|F = 1\rangle$ simultaneously. These states have different allowed transitions including some transitions to the hyperfine $D$ splittings which we can pump out with the 3.0695 GHz sideband. Many transitions can occur during this process but the important thing is that the $^2S_{1/2}|F = 0\rangle$, our $|0\rangle$ state, has no possible transitions driven from it. While every other state is being driven eventually to $|0\rangle$. After a span of time all the population will end up initialized to $^2S_{1/2}|F = 0\rangle$ which we call $|0\rangle$. For our system this process takes less than 5 $\mu$s depending upon the available optical power.
2.4.2 State Detection

For state detection we use the resonantly driven beam which has passed through a 210 MHz AOM and turn off the EOM on this beam. In this case if the ion is observed as $|1\rangle$ then it will scatter a photon and appear as bright, while if it is observed as $|0\rangle$ then it will be dark. In the case that the $^2P_{1/2}|F = 0\rangle$ state decays to the $^2D_{3/2}|F = 1\rangle$ state then our 935 nm light will re-pump it back to $|1\rangle$ where it again has a chance to absorb and emit a 369.5 nm photon. We then seek to detect this light with a high NA lens. We seek to maximize detection efficiency of ion fluorescence so that the detection time can be as small as possible minimizing the population lost by resonant driving of the $^2S_{1/2}|F = 1\rangle$ to $^2S_{1/2}|F = 1\rangle$ state, a source of infidelity. We discriminate the states with a threshold for PMT counts during the detection period which is usually set to 1.
Figure 2.14: A diagram explaining the qubit scheme for $^{171}\text{Yb}^+$ adopted from [OMM+09, OYM+07] is shown. Laser driven transitions are shown in cyan while 369.5 nm emission is shown purple and 935 nm emission is shown in red. (a) In order to drive all the population to the $^2S_{1/2}\left\langle F = 0 \right\rangle$ we apply a 2.105 GHz sideband with an EOM to our near-resonance 369.5 nm beam. This drives both the $^2S_{1/2}\left\langle F = 1 \right\rangle$ to $^2P_{1/2}\left\langle F = 0 \right\rangle$ and the $^2S_{1/2}\left\langle F = 1 \right\rangle$ to $^2P_{1/2}\left\langle F = 1 \right\rangle$ transition simultaneously. Some population is lost into the $D$ states but the same scheme from Figures 2.13 is applied with the 935 nm laser driving transitions between both hyperfine $D$ splittings. The only state here without transitions being driven out is the $^2S_{1/2}\left\langle F = 0 \right\rangle$ level and thus within less than 5 $\mu$s we can initialize to our qubit $\left\langle 0 \right\rangle$ state. (b) Here the state detection scheme is similar to the Doppler cooling scheme except we use a near-resonant 369.5 nm beam to drive the $^2S_{1/2}\left\langle F = 1 \right\rangle$ to $^2P_{1/2}\left\langle F = 0 \right\rangle$ transition with the highest scattering rate possible.
Chapter 3

Manipulation of the Qubit

The execution of quantum algorithms require good control of the qubits. After initialization we require the ability to prepare arbitrary qubit states on the Bloch sphere, globally or individually. We also require the ability to perform two-qubit entangling gates. With these two ingredients we will have a universal gate-set for quantum computing [NC10]. Our two-qubit entangling gates makes use of the phononic degree of freedom which we can control coherently with counterpropagating Raman lasers. This good control of the phononic degree of freedom will also enable us to resolve motional sidebands so that we can prepare ground states of motion. We can achieve global qubit rotations with high fidelity using microwaves tuned to the 12.6 GHz hyper-fine transition. We will be able to individually address qubits by focused optical beams. Our ability to span the 12.6 GHz frequency gap is enabled by the wide-bandwidth frequency comb of mode-locked pulse laser technology.

A useful visualization we can use to imagine arbitrary manipulations of qubits is the Bloch sphere shown in Figure 3.1. Uniquely determined by the angles $\theta$ and $\phi$ we can write a state $|\psi\rangle = \cos(\theta/2)|0\rangle + e^{i\phi}\sin(\theta/2)|1\rangle$. Referred to throughout this work is the jargon $\pi$-time or $\pi$ rotation. The $\pi$-time is simply the interaction time required for the qubit to undergo a $\theta = \pi$ rotation on the Bloch sphere. If prepared into a $|0\rangle$ state we will have a $|1\rangle$ state after the interaction is completed.

The Pauli matrices, $\sigma_x, \sigma_y, \sigma_z = X, Y, Z$, form a complete basis for rotations on the surface of the Bloch sphere; that is any arbitrary manipulation of the state vector can be decomposed into Pauli matrices [NC10].
Figure 3.1: A diagram of the Bloch sphere. Positions on the surface of the Bloch sphere can be written as $|\psi\rangle = \cos(\theta/2)|0\rangle + e^{i\phi}\sin(\theta/2)|1\rangle$. The term ‘$\pi$-time’ or ‘$\pi$ rotation’ is used to denote the interaction time required for the qubit to undergo a $\theta = \pi$ rotation on the Bloch sphere. Prepared into a $|0\rangle$ state we will have a $|1\rangle$ state after the interaction is completed.

3.1 Rabi Oscillations with Microwave Sources

The simplest scheme for trapped ion qubit manipulation is to utilize a microwave source in free-space [BLMW04]. The microwave radiation must be directed at the ion trap with some collimation, and microwave horns with small opening angles are a simple means of achieving this [BLMW04]. Rabi rotations in response to oscillating fields are a standard technique whose physics maps onto many other quantum systems [NC10]. With reasonable microwave powers of less than a watt we can achieve Rabi frequencies above 100 kHz provided the horn can be located close enough to the trap [BLMW04]. In our system we utilize a 12.6 GHz frequency tuneable microwave source which we set to be roughly 120 MHz away from the qubit transition. Using a single sideband mixer we mix this microwave source with a direct digital synthesizer (DDS) signal which can be scanned by the control software.

We attach a microwave horn to a long metal rod and hang it over our chamber
pointing downwards into the unused top view-port. By simply pointing the horn into the cryostat we can easily obtain some signal by search. We scan our DDS in the neighborhood of 120 MHz with some arbitrary reasonable starting time on the order of 1 ms. Once we find the microwave frequency peak by this method, which due to luck may be quite narrow or quite strong, we can set this as our DDS frequency and instead scan the interaction time. We optimize the position of the microwave horn to obtain the fastest Rabi $\pi$-time we can. With minimal effort we can obtain $\pi$-times around 50 $\mu$s. No special effort is made to optimize this because we only use it for rare calibrations and diagnostics; it is not used for quantum circuits in our system.

### 3.2 Raman Rabi Oscillations with Frequency Combs

The optical frequency comb has been a workhorse in metrology and spectroscopy [HMM+10]. A mode-locked laser will produce a short pulse with a spectrum which can be approximated by a series of delta functions which are separated in by the repetition rate of the laser [STS91]. This is also the inverse of the time it takes for light to traverse the cavity ($c/2L$) [STS91]. Spanning the full 12.6 GHz hyperfine transition is technically difficult using modulation, or two phase locked lasers alone [HMM+10]. Frequency combs from mode-locked lasers can have bandwidths much broader than the qubit transition. We take advantage of this precisely spaced frequency landscape spanning far beyond our qubit frequency. Despite the fact that the light frequency is far-detuned from our atomic transitions, a ‘virtual’ transition can occur as photons from one comb tooth are absorbed and then remitted into the other comb tooth [HMM+10]. This can happen if the repetition rate or its harmonics is set to be in resonance with the qubit frequency. This can also occur by the interaction of two comb teeth which are offset by optical modulation. When this relationship
between the qubit frequency and the frequency comb (or beat note comb) exists, all of the frequency components will add constructively with their mate [HMM+10]. In order address the qubit resonance we just enforce the relationship in Equation 3.1 below.

\[ f_{\text{qubit}} = nf_{\text{rep}} \pm (f_1 - f_2) \]  

(3.1)

In our system we split the output of a single mode-locked Coherent Paladin Mini (3.5 W output) and then modulate both beam paths, one for applying control tones, and the other for stabilizing the frequency relationship between the combs. Locking the frequency relationship of Equation 3.1 is done by detecting a higher harmonic of the repetition rate and feeding changes in it to one of the control AOMs. This ‘feed forward’ technique is superior to directly controlling the laser cavity because of the limited bandwidth with which the cavity can be controlled [IVH+14].

In order to implement the feed forward lock we sample the laser at its output and send it to a fast photodiode. In order to track the repetition rate, we mix the photodiode signal with a local oscillator at 3.876 GHz and then send the output through a low-pass filter to remove the upper sideband. From this we obtain an error signal based on the 32nd harmonic of the repetition rate which we will see as a signal at 3.876 GHz – 32 × 119.58 MHz = 49.44 MHz. From here we can then rewrite Equation 3.1 as 12642.83 MHz = \( n \frac{3876 \text{MHz} - 49.44 \text{MHz}}{32} + (f_1 - f_2) \). From here we need to find some number, \( n \), of comb tooth separation which along with some pair of \( f_1 \) and \( f_2 \) will be within the range of our AOMs.

When setting this up it is convenient to choose one of these frequencies to be an integer and the other to be finely tuned. One example of a setup in our system is that we choose the individual AOMs to be 215.0 MHz. We use a double pass AOM
Figure 3.2: A diagrammatic representation of the frequencies in equation 3.1. This helps to illustrate the mechanism of the repetition rate feed forward method for Raman qubit manipulation by frequency comb.

with a center frequency of 210 MHz so this is a convenient number which then gives $f_1 = 430$ MHz. If we chose $n = 104$ we will see that by solving for $f_2$ we should expect our qubit transition to occur by driving at 222.645 MHz on this AOM. This will only occur if we feed forward changes in the repetition rate (as detected by this error signal at 49.7 MHz) to the other AOM around 215 MHz. In reality this number will be slightly offset due to Stark shifts. Figure 3.2 is helpful for visualizing these variables and the frequency comb.

If we apply similar logic as above, but to a single beam with two AOM tones we can set the polarization appropriately (circular) and drive Raman gates on ‘carrier’ transition. We use this term carrier when talking about Raman operations to denote transitions which only change the hyperfine qubit (or spin) degrees of freedom, as in the microwave gates.
The motion of ions is governed by the harmonic oscillator Hamiltonian and chains of ions can be described as coupled harmonic oscillators. Arranging the frequencies as described before, a repetition rate locked tone on one side (we choose the individually addressed side) and another tone which can be scanned in relation to this, and setting these beams as counterpropagating we can create a net momentum vector $\Delta \vec{k}$. This large net momentum vector is enabled by the two-photon Raman process. The coupling parameter, also called the Lamb-Dicke parameter $\nu$, can be large for two-photon processes because it becomes proportional to the difference of the photon’s momentum vectors [Hay12]. By arranging the beams either orthogonally or in a counterpropagating geometry we can create large interaction $\Delta \vec{k}$. If the ion is in the Lamb-Dicke regime, meaning the $\sqrt{n} \eta < 1$ where $n$ is the mean phonon occupation of the ion’s motional mode when we will be able to resolve and manipulate our ions motion.

When the second AOM (not locked to the repetition rate) is scanned from the above example we will then see a carrier transition around $f_c = 222.654$ MHz. Also in the case of the counterpropagating beams we will see two radial modes at $f_{x,z} = f_c \pm \nu_{x,z}$. If the momentum vector has a projection on the trap axis then we will also see the axial motional mode at $f_y = f_c \pm \nu_y$. If there are $N$ ions then we will see $N$ motional modes for each degree of freedom. These normal modes of motion arise similarly to any coupled oscillator and a good discussion of them is found in [Hay12]. If we consider the two radial modes of motion there is the common mode of motion and the tilt mode. The tilt mode is related to the radial common mode $\nu_{x,z}$ and axial common mode $\nu_y$ by $\nu_{\text{tilt}} = \sqrt{\nu_{x,z} - \nu_y}$ [Hay12].

By red detuning (from the carrier) our addressing laser by $\nu_{x,y,z}$ we can drive the states $|0n\rangle \rightarrow |1n - 1\rangle$ for the chosen mode, we call this driving the red motional sideband of that mode. By blue detuning our addressing laser by $\nu_{x,y,z}$ we can drive
the states $|0n\rangle \rightarrow |1n + 1\rangle$, we call this driving the blue motional sideband. The Rabi
$\pi$-time for these manipulations is longer than the carrier $\pi$-time, if the $\bar{n} \approx 0$ then
the factor by which it is longer is by $1/\eta$.

Using this ability to manipulate quantum states of motion we can cool below the
Doppler limit using a simple algorithm. We prepare a Doppler cooled $|0n\rangle$ state, and
apply a motional $\pi$ timed pulse at the red sideband frequency. This brings the state
to $|1n - 1\rangle$, from here we use optical pumping to return it to the $|0\rangle$ spin state, but
now it is $|0n - 1\rangle$. This is can be repeated until the ground state is reached. This
procedure is called sideband cooling \[WMI^+98\].

This procedure is complicated by the fact that the motional $\pi$-time will continue
to get longer, approaching the $1/\eta$ factor as we cool. We must then lengthen the
interaction time as $n$ decreases. Furthermore this procedure must be done for each
motional mode individually in series. If we restrict our cooling to only the radial
modes then this must be done $2L$ times, where $L$ is the chain length (in ions). For
our case we adjust the time by a simple linear algorithm ramping up from a short time
close to the carrier $\pi$-time to a time $\frac{1}{\eta} \times \pi$-time in even steps. While cooling the other
modes the previously cooled modes can heat by the action of optical recoil \[WMI^+98\]
or by the natural heating rate of the trap. A choice is made as to which mode is
cooled last for gate operations. In our case the common mode heating rate will be
sufficiently low and we perform our two-qubit gates on the tilt mode so we cool the tilt
most last. After sideband cooling we can prepare base phonon occupation numbers
for two ions of $(\bar{n}_{\text{common}}, \bar{n}_{\text{tilt}}) = (0.35 \text{ quanta, } 0.05 \text{ quanta}).$
3.3 Two-Qubit Gates with Raman lasers

In order to have a complete set of quantum gates we require a two-qubit entangling gate which can prepare Bell states such as $\frac{|00\rangle + |11\rangle}{\sqrt{2}}$ [NC10]. We use the Mølmer–Sørensen gate which has the unique property of generating entangled states with no explicit dependence on $n$ [SM99]. This means that our entanglement gate can be used on thermal distributions of number states and does not depend on the particular $\bar{n}$ at which we start the gate, and will also be more tolerant to heating than other schemes [SM99]. The gate is performed by applying two driving tones near the motional frequency used for the gate which are equally detuned. This detuning is symmetric about the carrier. In Figure 3.3 we see the classic diagrammatic representation of the Mølmer–Sørensen gate. In this figure we see the spin energy levels of a two ion system along with the motional states, $n-1$, $n$, $n+1$, of the odd population states. In order to realize the Mølmer–Sørensen gate we utilize an AOM where we drive two tones simultaneously, the red and blue tone, each related to the carrier $f_c$, the motional frequency $\nu$, and a single detuning $\epsilon$ by the relations $f_r = f_c - \nu - \epsilon$ and $f_b = f_c + \nu + \epsilon$ respectively.

If we inspect Figure 3.3 the only energy conserving transition which can be driven by a two-photon interaction is $|00\rangle|n\rangle \leftrightarrow |11\rangle|n\rangle$. Once the interaction is turned on the populations of all states will evolve in time according to the scheme in Figure 3.4 which we plot in terms of the gate angle $\Theta$. This gate angle is the angle of the state with respect to the $XX$ axis. At $\Theta = \pi/4$ the entangled state $\frac{|00\rangle + |11\rangle}{\sqrt{2}}$ and we call this the gate time, at $\Theta = \pi/2$ we see that the state disentangles and becomes $|11\rangle$. If the interaction is allowed to continue then this process will repeat, it is only at odd numbers of gate times that we see entanglement.

More detail is given on the requirements and noise sources for Mølmer–Sørensen
Figure 3.3: In this diagram we see a visualization of the Mølmer–Sørensen gate. The red and blue arrows represent laser drives which are created by adding an RF tone to the AOM with offset from the carrier, \( f_r = f_c - \nu - \epsilon \) and \( f_b = f_c + \nu + \epsilon \) respectively where \( \nu \) is the motional frequency. The energy conserving transition which can be driven in this scheme becomes \(|00\rangle|n\rangle \leftarrow |11\rangle|n\rangle \) [SM99].

Figure 3.4: The evolution of the Mølmer–Sørensen gate in time, in terms of angle with respect to the XX-axis. The entangled state \( \frac{|00\rangle + |11\rangle}{\sqrt{2}} \) is generated at \( \Theta = \pi/4 \), the time at which this happens we call the gate time. The ions are then disentangled at \( \Theta = \pi/2 \) and the state \(|11\rangle\) is produced.
gates in Chapter 6. Drift in system parameters such as laser power, red/blue tone balance, motional frequencies, and excessive ion heating will all contribute to imperfections in the gate. We measure the imperfections mainly by the inaccuracy with which we realize the gate angle $\Theta$ and the residual population in the states $|01\rangle$ and $|10\rangle$ when we reach the gate time. The gate can be considered dynamically as a phase space loop. The motional states are displaced in phase space by the interaction and must undergo a loop in which they return to the origin. This loop in phase space will enclose an area which corresponds to the gate angle. The residual displacement at the end of the interaction is proportional to the remaining odd population $p_{01} + p_{10}$, inaccuracies in this enclosed area lead to inaccurate rotations with respect to the $XX$-axis. Much recent work has been done on gate schemes which are robust to system drift [LLF$^+$18, KLZ$^+$21, MEH$^+$20]. These schemes usually include the modulation of the gate control tones (amplitude, frequency, or phase) in ways that manipulate the path through phase space ensuring the displacement error is minimized. These gates can be much more tolerant under the influence of noise and miscalibration than the raw gate.

\[ p_{00}, p_{01}, p_{10}, p_{11} \] to denote the populations of the $|00\rangle$, $|01\rangle$, $|10\rangle$, $|11\rangle$ states respectively. This is observed by repeating the experiment (usually 100 times) and measuring the ion fluorescence (as bright or dark). In accordance with the postulates of quantum mechanics, when measured in the computational basis, these experimentally derived populations should be approximate the probability of the ions being in each state if observed [NC10]. These probabilities relate to the state amplitudes ($c_{00}$, $c_{01}$, $c_{10}$, $c_{11}$) by $p = |c|^2$ if the normalization condition is met [NC10].
Chapter 4

A Compact Cryogenic Ion Trap Platform

4.0.1 Motivating the Trap Cryopackage Design

The installation of surface ion traps into Ultra-high vacuum (UHV) chambers is a time consuming and delicate process [Mou15]. Using the UHV platform restricts choice of materials for all chamber internals. The strict requirements of the assembly procedure help to achieve the lowest pressures possible. The surface ion trap itself must be handled carefully from fabrication through installation in order to prevent dust contamination. Facilities which seek to utilize surface ion traps are required to have access to a clean room facility, a large oven for chamber baking, and invest several weeks of labour and processing if they seek to replace the ion trap or service the chamber interior. Additionally the large size of the UHV chambers is ill suited for removal and servicing without the surrounding optical alignment being disrupted.

Cryogenic temperatures offer access to UHV pressures with relative ease due to the fact that all surfaces participate in the pumping via adsorption on surfaces at low temperatures [Eki06]. Below about 10 K most gasses have frozen out leaving only helium and hydrogen remaining. Additionally any collisions that do occur originate from a low temperature reservoir and so the gas particles have small kinetic energies. Due to the low collision energies fewer events which are detrimental to complex quantum algorithms, such as chain reordering and catastrophic chain loss occur. Even the best pressures that can be achieved at UHV still will suffer from these occasional catastrophic events coming from 300K background atoms colliding with the chain.
A new platform is needed for ion trapping experiments which is more amenable to mass production and distribution of standardized parts. However manufacturability and convenience should never come at the cost of system performance. Thus we sought a design with better performance than state-of-the-art UHV chambers which would also solve many of the practical hindrances of the ion trap platform.

We conceived of a new design approach in which the UHV space was reduced to the minimal features needed to trap an ion. At the center of our approach is a compact cryogenic ion trap package, the ‘trap cryopackage’. The only features which need be included in the UHV space are the ion trap itself and an ion source. The trap cryopackage must have adequate optical access for high NA imaging and tightly focused Raman lasers. Feedthroughs for DC and RF voltages are required and are naturally furnished by the ceramic pin grid array (CPGA) package to which the trap is attached. In order to ensure excellent vacuum levels we include sorbtion getters, in this case activated carbon. This carbon getter will trap the remaining hydrogen and helium present in the cryopackage when kept at low temperatures.

We created an ambitious first generation design which featured a fully vacuum sealed package. It included many major changes to the ion trap system all at once. A delicate, isotope selective ablation loading was required to work reliably. The helical resonator of many ion trap systems was replaced with a RLC lumped element circuit driven by low output impedance amplifier. The design details of this system and some key work on ablation characterization are included in Appendix C. In this system we never were able to demonstrate ion trapping. With all of these major system changes it was difficult to know if the cryopackage, the RF circuit, or the ablation were to blame. We had good confidence that our ablation loading technique was correct. Confidence in the ablation technique was bolstered when the same technique was used to trap in a room temperature UHV chamber [VAS+19]. These ablation techniques
were later applied to our second generation system with no problems [SIJ+21]. The most likely candidates for the inability to trap were achievable RF amplitude, the possibility of cryopackage leaks, but chiefly a poor thermal design which left the metal lid—containing the carbon getter—the hottest part of the setup.

Our RF circuit driven by low output impedance amplifier was a clever approach, but it is always limited by the ability of a particular amplifier to drive current at a particular frequency. At 30 MHz this approach would have adequate voltage amplitude but in the 40-50 MHz regime we couldn’t achieve the voltages required for good ytterbium manipulation and were left searching for a suitable amplifier.

Our sample chamber design depended upon the cooling of the trap printed circuit board (PCB), and thus the cryopackage lid had an indirect cooling path through the CPGA. This meant that while the cold finger below would sit at 5 K with RF at full amplitude, we would measure lid temperatures of 20 K or more. We would later find out that these temperatures would cause gasses to desorb from the carbon getter making stable trapping difficult in excess of 15 K. Lastly our fully vacuum sealed package suffered from leaks which were large enough to trap humidity inside the package during transport and storage. While also too small to allow for the conductance necessary evacuate the package before cyrocooling. We eventually converged on an approach where the cryopackage would feature a meandering path which would create a differentially pumped volume upon cryocooling. This meandering path would have enough conductance that the cryopackage could be well-evacuated using a turbomolecular pump (turbo) before cyropumping commenced.

For our second generation design it was decided to obtain a cryostat whose sample chamber could accommodate the inclusion of a helical resonator which could be used to drive the RF. Furthermore during the design phase of the previous experiment, it was noticed that high NA imaging was limited by three layers of viewports (pack-
age, 40 K window, room temp vacuum) which added a lot of optical path length before the lens can be placed. In order to overcome this limitation we sought to include an ion imaging lens which could compete with the 0.6 NA photon gear lens of other experiments [CCV+19]. We then included our imaging lens inside the cryostat with alignment stage. This new experiment then would benefit greatly from a more spacious sample chamber.

4.1 The Cryostat\(^1\)

We utilize a closed-cycle Gifford-McMahon (GM) cryocooler which has been designed into a mechanically stable module by Montana Instruments (Montana Instruments Cryostation s200). Montana Instruments integrates a Sumitomo SRDK-101E into their low vibration system complete with sample chamber and feedthroughs. This cryostat features low levels of vibration despite the mechanical motion of the Gifford-McMahon cryocooler. The cryocooler has two stages; the stage 1 cooler sits nominally at 30 K with 3-5 W cooling power at 45 K and the stage 2 cryocooler has 0.1-0.2 W of cooling power at 4.2 K. We speak more about the thermal performance under various heat loads in section 4.4.2. The Montana system as delivered had a base temperature of 5K on the cold finger and so we usually refer to this as the nominal base temperature. Another unique design feature of this cryostat compared to traditional designs is that the sample space is anchored on the tabletop surface rather than hanging from the top [PHK+18, BvMP+16, VWB+13b], which makes swapping of the trap cryopackages and upgrade of the cryostat sample chamber straightforward. It has the added advantage that optical alignment can be tested before the cryostat is closed up for cooldown.

\(^1\)Section is adapted from [SIJ+21] where it first appears some parts are subject to © 2021 IEEE
A labelled diagram of the Montana Cryostation system. The vacuum housing and lid with o-ring seals provide the vacuum environment needed for cryogenic operation. The radiation shield is mounted to the 90 K stage, and includes all windows necessary for optical access. A gold coated copper sample mount is attached to the 5 K cold-finger platform. The trap is mounted within this sample mount such that the trap axis (defined as the x-axis) points along the vertical axis. Appearing first in [SIJ+21] © 2021 IEEE

A diagram of the cryostat and its internals is shown in Fig. 4.1. The internal sample chamber includes a large flat 90 K stage with a grid of tapped holes for mounting various components. In the center there is an exposed 5 K base plate, upon which we mount a gold coated copper sample mount that houses and cools our trap cryopackage. A radiation shield with optical viewports through all necessary axes is thermally anchored to the 90 K stage. The vacuum housing has two o-ring
seals, one making contact with the base of the cryostat, and the other on the top side sealing the lid. Several temperature sensors are installed to monitor the temperature of the system, including the 5 K base plate and the sample mount near the trap.

This cryostat will feature plenty of room for a more complicated interior. New features such as a high NA imaging lens and a cold helical resonator can be included.

4.2 Trap Cryopackage Design and Assembly

The vacuum inside the cryostat is maintained with o-ring seals, as opposed to copper gaskets which are typically used in UHV-chambers. This sealing mechanism limits the pressure to the low $10^{-8}$ Torr range. To achieve UHV vacuum levels in the trapping region, we create a secondary chamber assembly to hold the trap, which supports a UHV environment within the cryostat via differential pumping. The design of the trap cryopackage is shown in detail in Fig. 4.2. The assembly consists of a lid sealed to a ceramic pin-grid array (CPGA) package, on which the ion trap is mounted. The lid is machined out of copper, the side windows are made out of N-BK7 glass (5 mm diameter with 2 mm thickness), and the imaging window is made out of sapphire (15 mm diameter with 2 mm thickness). All windows are AR coated and attached to the lid with cryogenic-compatible epoxy (Epotek T7110).

Sapphire was chosen due to the coefficient of thermal expansion (CTE) being well matched with metals such as copper and titanium. An additional benefit of sapphire is the low permeation of hydrogen and helium. The downside is a birefringence-induced astigmatism that leads to polarization-dependent image focus. To avoid this issue, glass material such as fused silica should be used, but care must be taken to manage the CTE mismatch during the cooldown cycle. Limited anecdotal evidence

\*Section is adapted from [SIJ^+21] where it first appears some parts are subject to © 2021 IEEE
suggests that Epotek T7110 may work better for fused silica windows than Epotek T7109 due to the lower shear stiffness. Furthermore care must be taken to avoid epoxying the edges of the windows to the copper, rather only the bottom should be attached. The side windows were chosen consistently as N-BK7 to avoid birefringence problems on the Raman beams. These beams can tolerate slightly higher absorption losses and stress-induced cracking has not been seen.

The top imaging viewport features a machined trench which serves as a meandering narrow pumping port for evacuating the volume inside the cryopackage at room temperature before the system cools down. The 100-pin alumina CPGA package is modified with a ‘ringframe’ brazed to it, which is 1 mm wide and protrudes 2.5 mm from the CPGA surface. This ringframe mates with a matching groove in the lid, designed for a tight fit with indium wires in between. The ringframe serves both as an alignment mechanism for the lid and as mechanical reinforcement for their mating.

The indium wires increase the thermal conduction between the lid and the CPGA package, and serve as a breakable mechanical bond between the two pieces. A true hermetic seal is not required for our differentially pumped open-lid design, and the adhesion provided by the indium is sufficient to keep the lid and the CPGA together during handling and installation.

The top of the lid features eight tapped holes for affixing the lid firmly to the cold finger sample mount (Fig. 4.2(a). This top surface is the main thermal and mechanical anchor for the assembly on the sample mount. A few pieces of activated carbon getter material, encapsulated by a fine copper mesh, is inserted into a cavity of the lid interior. Upon cooldown, the lid is cooled most efficiently, and the cryopumping by the getter achieves a UHV environment within the sealed enclosure.

Some important internal features of the lid can be seen in Fig. 4.2(b). The first is an oblong conical ground shield feature which is machined beneath the imaging
**Figure 4.2:** Compact ion trap cryopackage design (a) On the exterior, it features eight tapped holes on the top and a spiral meandering pathway cut into the copper lid for pumping. The imaging window covers and makes the fourth wall of this meandering pathway, which is used for gas evacuation. AR coated N-BK7 side windows are epoxy-sealed to this copper lid. (b) A cross section shows the interior features of this assembly. A sapphire imaging window is epoxy-sealed to the top. The internal features include a cavity for the storage of carbon getter packaged in a copper mesh, and a holder for storage of the Yb ablation target. The ion trap stack is packaged with an interposer for routing wire bonds, and a spacer to raise the top surface of the trap to clear the ring-frame height in the lid. Appearing first in [SIJ+21] © 2021 IEEE

window to allow the full numerical aperture of the imaging optics to be available while maximally shielding the potential impact of charge buildup at the exposed dielectric window from influencing the ions. There is also an internal holder to house a ytterbium ablation target. This ablation holder has optical access from a window directly across from the target, such that the path of atomic flux is orthogonal
to the CW laser beam path for photoionization lasers to enable isotope-selective loading [VAS+19].

The trap (Sandia HOA 2.0 [Mau16b]) is mounted on an interposer, placed on a ceramic spacer to bring the trap surface in line with the windows of the lid. The interposer also contains trench capacitors (TCs), to filter the RF signal on the DC electrodes. The HOA 2.0 trap we used in the experiment had an additional layer of gold (0.5 µm) deposited on the surface of the electrodes.

The trap cryopackage is assembled in a clean room to minimize dust contamination and is installed into the PCB by pressing it into the 100-pin socket. This PCB and trap cryopackage assembly is installed into the 5 K sample mount in the cryostat chamber. We include in detail the recipe for making the trap cryopackage lid in Appendix A and we also include mechanical drawings of the Peregrine lid and cold finger sample mount in Appendix B.

4.3 Cryostat Internal Design

Figure 4.3 shows a cutaway of the interior elements included in the sample chamber of the cryostat from two useful views. Figure 4.4 shows a real image of the cryostat interior with labels. Some elements are more easily viewed in the real image, while others more easily viewed in schematic cutaway. The copper colored central tower piece is called the cold finger sample mount. We call the gold pieces below it the 5K cold finger which is thermally and electrically tied to the stage 2 cryocooler. The central metallic colored plate we call the experimental platform or simply the platform, this is nominally a 40K plate with no other heatload, but we have measured it as 90 K when all of our equipment is installed. This is thermally and electrically tied to the stage 1 cryocooler.
Inside the sample chamber, we install three custom experimental elements on the 90 K stage. First is the RF helical resonator (more detail can be seen in Figure 4.5), this is designed to be smaller than most UHV chamber helical resonators and to mount neatly with a square edge to the sample chamber platform. This can be seen installed in a cutaway (less the scorpion tail) in figure 4.3(a) as a copper colored object with a green helical interior. Here the green helix (representing the antenna) is held mechanically stable by Teflon coil holders. A real image of the device can be seen in Figure 4.4. This resonator has a coaxial cable input and a simple thick copper wire output. It is mounted close to the trap to avoid having a long cable between the resonator and the trap. A ‘scorpion tail’ style copper pipe ground shield is extended beyond the resonator to come within 1.5 inches of the trap PCB. Fully connecting the ground shield at the trap PCB would dump too much heat into the trap PCB. In practice we connect a stainless steel wire between the helical resonator body (ground) to the ground plane of the PCB. This wire has a 10 mil (0.250 mm) diameter and is a few inches long, it is soldered to the trap PCB and pressed onto the resonator by screw and washer. This practice is questionable as it only provides a good ground at low frequencies where such a ground is not desirable. This matter is discussed more in the final chapters of this document.
Figure 4.3: Two cutaway views of the Montana cryostat internals (a) in this view we see all of the system lasers. We can see the MEMs individual beams incident from the left going through the Raman projection lens atop a z-stage, the global beam comes in from the right. The CW lasers are incident from the top left, from the top right we have the ablation laser. Machined into the the cold finger sample mount is a cavity for storing a temperature sensor mounted to a pill. (b) In this view we can see the top of the CW laser delivery plate. We see the imaging lens in this view in a cutaway. We also point out mounts for natural magnets as well the locations of o-ring seals for the top and bottom of the sample chamber.
Also mounted on the 90 K stage are two lenses specially designed for operation at cryogenic temperatures. The first and larger of which is a 0.6 NA ion imaging lens (Photon Gear) with a 150 µm field of view. A cross section of this (less the proprietary multilens stack) can be see in Figure 4.3(b). This is mounted inside the cryostat to achieve a short working distance to the ions. In order to align this lens we designed a specialized alignment stage. It has two micrometer-adjusted X and Y positioners which can only be adjusted at room temperature. The stage also includes a custom kinematic-mount, with three cryogenic precision actuators (Attocube Inc.) that enables the adjustment of the tip, tilt, and z-translation of the lens. These are critical knobs to minimize imaging aberrations. The imaging lens is designed to be an infinite conjugate lens which allows us to use any off-the-shelf lens outside the cryostat to form an image with desired magnification. The infinite conjugate design also reduces aberrations that would arise due to traveling through the 90 K and room temperature view-ports.

In practice it is quite difficult to achieve diffraction limited imaging quality due to the lack of in-situ x/y adjustment. The strategy used is to design a jig which allows a camera to be mounted with a simple long-working-distance zoom lens, whose position can be guaranteed by dowel pin alignment. This can image through this lens and see the trap slot, and the lens edges at the same time. We can adjust the x/y positioners to ensure that the trap slot is aligned with the center of the lens in x. Ideally this imaging system can also resolve the trap alignment markers, so that we can ensure the trap center can be vertically aligned to lens center. During cooldown the micrometers can drift, and the size of the mechanical stack will contract, so in practice it can take a few cooldown cycles to reduce off-axis abberations to an acceptable level. An improvement in the alignment procedure or added in-situ degrees of freedom are needed to achieve system efficiencies which can compete with [CCV+19].
The second lens is a Raman projection lens (Photon Gear) for the final tight focusing of individual addressing beams to the ion chain. We can achieve a 1.7 \( \mu \text{m} \) focused beam waist at the ion with this lens. This lens is placed manually at room temperature, with an \textit{in-situ} translation stage for linear focus control only which is sufficient. This lens and stage are shown in cutaway form in Figure 4.3(a). Brass-tipped set screws are used to clamp down to the top of the lens (best seen in Figure 4.4(b), the lens stage is simply resting atop of an aluminum block designed to place it at the correct height. We use machining screw tolerances to set the lens x and y position which is accurate enough with the beam steering degree of freedom to achieve the correct focus.

4.3.1 Trap Mounting and Cooling

The most important functional aspect of the experiment is the mounting and cooling structure for the trap. The interface with which the cold finger sample mount meets the cryopackage lid is the main mechanical reference surface for the experiment. The sample mount itself is dowel pin located with respect to the cold finger cooling surface below, with two dowel pins on opposite sides. The sample mount itself has a round receptacle and a slot for mating with these dowel pins which will locate the sample mount and hold its rotation repeatably static. Built into the sample mount are four 2 mm dowel holes which register the lid location in this plane. The Peregrine lids then have corresponding dowel holes (with a tighter tolerance) and dowel pins are inserted into the cryopackage. The dowel pin feature was added for the Peregrine upgrade. Previously the location was fixed by 8 tightly countersunk screws. The ion location is fixed relative to the cold finger then by the lid design. If an update is made to trap packaging, or if different traps are used it is much easier to redesign and manufacture a new lid which will correctly locate the ion with respect to the experiment. While
the traps location in packaging is important with respect to the lid, and lid registers this to the sample mount within a $\pm 50 \mu m$ tolerance, the harder to control locations of the trap PCB and its supporting mechanical stack are unimportant and subject to thermal contraction and thus we don’t have to depend on them. Furthermore by mounting the lid directly embedded into the sample mount we guarantee that it will be the coldest.

In our design access to the screws holding the trap into the cold finger are blocked by the imaging lens. The lens stage must be unbolted, moved backwards, and the lens itself entirely removed before the trap can be installed or removed. In future designs having dowel pins on the 90 K platform could help alleviate how disruptive this is.

As discussed in the trap cryopackage recipe the trap is always installed into the trap PCB before delivery to the experiment. Before attaching to the sample mount a small amount of thermal grease is applied to the top of the lid before the lid and sample mount are pressed together. The system is designed to maintain the copper package lid as the coldest part in the trap cryopackage, ensuring that the temperature of the carbon getter is as low as possible. It is important to use two people from either side to manage the installation because unwanted torques can cause the lid to disconcert if epoxy was not used. Once the dowel pins are engaged and all screws (2-56) tightened between the lid and sample mount, the PCB must be secured. The cold finger also has standoff mounts to which the PCB can be screwed down independently so that the trap ringframe joint does not need to be load supporting. This PCB is itself secured with a set of standoffs which will be used to mount a copper structure we call the cold bridge.

The purpose of the cold bridge is to directly cool the back of the trap CPGA. This structure can be seen best in Figure 4.4(a). The cold bridge has two solid copper
pieces connected by a thick bundle of copper mesh. One piece extrudes from the sample mount and simply has a clamp for the mesh. The other piece screws into the trap PCB standoffs with long spring-loaded screws and it also has a matching clamp for the copper mesh. This piece has a small pedestal which makes direct contact to the back of the trap CPGA. At this interface a small amount of thermal grease is used. We use the spring-loaded screws to keep this copper pedestal pressed against the trap during cool-down while still leaving some forgiveness in the entire stack for contraction. The mesh is clamped on tightly and the spring loaded screws are left slightly more than finger tight. Before installing the cold bridge it is wise to first plug in the DC connections and do a full connectivity test while the chamber is open. You can use the DAC box to run unique voltages to all pins and then measure that these voltages make it to the soldered CPGA socket pins. After installing the cold bridge these pins are obscured.
Figure 4.4: A view of the cryostat sample chamber photographed in reality is presented here with labels for various aspects not CAD modelled.
After the trap cryopackage, trap PCB, and cold bridge are installed we can connect the DC cables which are 50 pin flexi-PCB ribbon cables (Omnetics). The wires are made of phosphor bronze and the ribbon made of Kapton. Lastly we connect the RF helical resonator output to the RF pin on the trap PCB using a barrel connector. After making this connection it is wise to use a spectrum analyzer to ensure the trap RF resonance can be seen while open, and verify the Q parameter before beginning to close the chamber. For good thermal performance it is essential that the DC cables do not touch the helical resonator on their way to the trap, and other such touches between 90 K thermally lagged connectors going to the trap should be avoided. A thin, ribbon style, tooth floss such as Crest Glide can be useful for creating jigging to pull on the DC cables such that it avoids touching the resonator.

After completely installing the trap and checking the RF resonance the imaging lens can be replaced and either alignment by eye, or by the previously mentioned method can be used to place it in x/y. Alignment by eye is accomplished by first visually centering the lens about the trap from the top, and using the stages to move the lens into close proximity with the trap window. It needs to be visually flat and in the neighborhood of (0.5 mm - 1 mm) from the window. At this point you should be able to lower your head to the level of the imaging axis and look for the trap visually through the lens. Your head can be 1-2 feet away and most people will be able to resolve an image of the trap. The entire lens stage can be translated side-to-side by hand if needed while visually scanning the head. One should look that the trap image will form evenly as the head is scanned rather then more to one side. In this way the lens can be aligned roughly by eye. A more fine alignment using jigging and an imaging system as discussed before is also advisable.
4.3.2 Vacuum Pumping and Cooldown Procedure

Once the trap is installed we are ready to pump and cool down the system. After closing the sample chamber lid it is wise to double check that the cryogenic stages are connected and that they move. We should also check that the RF resonance still remains. The details of the cooldown procedure are crucial for good vacuum performance. Time must be provided for the trap cryopackage to evacuate fully through the meandering pathway. We do not use the Montana system vacuum pumping hose to evacuate the sample chamber. A resistor was installed into the control unit to fool it into thinking the sample chamber is always at full vacuum. We cap out the output from the back of the Montana cold head which is used to evacuate it by roughing pump. We find that a large leak source for the system when cold for long times is the valve and roughing pump system inside the Montana control unit. We have a KF25 flange installed directly into the sample chamber to which we attach a vacuum-T. The vacuum-T has on one port a cold cathode gauge for monitoring the sample chamber pressure. On the other port we attach a valve. We use a 2-ft bellows to attach a turbo pump system to the cryostat.

After closing the sample chamber we pump the system with a turbo for at least 24 hours, ideally 48 hours before any cryocooling is begun. This ensures that the trap cryopackage lid can be completely evacuated before it is cooled. Crucially still we begin the cryocooling early in the morning which takes approximately 10 hours to begin approaching its base temperature. We do not close the valve to the turbo until the system has gone below 70 K. Ideally we close the valve to the turbo pump in the neighborhood of 20-30 K. We do not want to leave the turbo open for prolonged periods of time in the < 10 K regime because we may experience back-flow as the power of cryopumping exceeds the power of the turbo. The cryostat should approach its
base temperature at the 10-12 hour mark however we usually still wait an additional 12 to 24 more hours before applying RF to the trap, we notice a longer slow decline in the base temperature of about 0.5 K to 1 K occurring over this timescale. After this cooldown procedure we can operate for six months to a year at a time without again opening the chamber and still see the same vacuum performance.

### 4.3.3 Helical Resonator and Impedance Matching

![Helical Resonator Diagram](image)

**Figure 4.5:** A cutaway schematic diagram of the helical resonator construction. Here we can see the ‘scorpion tail’ ground shield made of copper tubing. We see two machined copper endcaps which will mate with the copper tube making the main resonator ground shield. The square profile of the endcaps makes for easy mounting to the 90 K platform.

Here we discuss the design, thermal characteristics, and impedance matching for the helical resonator. The helical resonator in cutaway model view can be seen in Figure 4.5. The construction consists of two square machined copper plates which have grooves machined out to mate with a copper tube of OD=1.75 in and ID=1.62 in. The total length of the device is 2.7 inches. The main helical coil is made using
uncoated 9 gauge copper wire which is mechanically stabilized by a Teflon structure. The coupling antenna can be manipulated by a micrometer mounted into one end for impedance matching. We couple into this antenna with an SMA connector mounted to the bottom copper plate. Because of the square geometry of the endcaps they can be stably mounted to the 90 K platform surface using clamps held by screws. The ground shield is built from segments of copper pipe joints which have been soldered together. The output wire is run through these copper pipes and held centered by a few Teflon spacers along its length. This entire ‘scorpion tail’ ground shield is clamped to the larger body of the helical resonator, the wire running through the scorpion tail and the main helical coil are joined by a barrel connector which is embedded in a mating clamp, spaced by Teflon. The coil pitch is 0.197 in$^{-1}$ or 5 turns per inch with a coil length of 1.29 inches and the effective shield length (total length less the endcaps) is 2.1 inches. The helical resonator output is also attached to the trap RF pin via barrel connector. The unloaded Q-factor of the device is approximately 460 with an unloaded frequency of 130 MHz. For the HOA2 trap this leads to RF frequencies in the range of 47.5-48 MHz with a maximum Q-factor of 120 once base temperature is reached (Q-factor of 80 at room temperature). With the same resonator design the Peregrine trap (with a lower trap capacitance) has a resonant frequency of 51 MHz and a room temperature Q-factor of 100. Once cooled the Q-factor we measure is closer to 180.

During the cooling process there are small mechanical changes which cause changes in trap capacitance and the matching condition but these effects are small compared to the changes in series resistance during cooling. This change is series resistance will increase the Q-factor, and change the matching condition, compared to their room temperature values. During cooling the resonance frequency of the system will also increase by approximately 1 MHz consistent with a decrease in trap capacitance on
Figure 4.6: The helical resonator parameters drift as the temperature decreases. (a) The quality factor increases due to the lowered resistivity of materials at low temperatures. Here is an example trace from a cool-down of the Peregrine trap. We reach the maximum Q-factor of 180 at a base temperature of 8.0 K. (b) Here we see that as the temperature decreases the resonance frequency goes up. This behavior is consistent between traps and is hypothesized to be due to small repeatable changes in trap and stray capacitance.

the order of 1 pF. The dynamic behavior of Q-factor and resonance frequency are shown as a function of system temperature in Figure 4.6, where time flows from high temperature to low temperature, the shown traces depict the cool-down process.

In order to match the helical resonator to a new trap/CPGA system the thermal drift during cooldown must be taken into account. If the resonator is perfectly matched at room temperature then the reflected power can easily only be -10 dB when cold. For a new trap/CPGA system keeping track of the movement of the matching micrometer and the reflection ‘dip’ vs temperature trace can help to minimize the amount of cool-down iterations it takes to find good impedance matching.
Figure 4.7: The matching of the helical resonator to a new trap/CPGA system (the Peregrine system from the HOA 2.0 system) (a) We see three traces, one for each of a series of independent cooldowns with a different position of the matching micrometer which controls the coupling antenna. The resonance ‘dip’ is the difference between background and the minimum reflected power during a scan of the tracking source, which occurs at the resonant frequency of the trap-resonator system. The lower this dip goes the less power is reflected and thus the better the resonator is matched. The impedance will be best matched at a specific temperature for each different trace. (b) Plotting the temperature at which the match is achieved against the matching micrometer position allows one to fit a simple linear model to the matching process. The red point is our zero and where we begin counting and so we start by finding the room temperature matching condition, then we count how many times we turn as we detune from room temperature matching, the orange point is the 1st cooldown, the cyan point is the 2nd cooldown, with these three points a linear model was used to predict the position of good matching at 8.5 K (high voltage operating temperature). The final magenta point (and curve) is the result of this prediction, where we achieve a good enough match.

An example is presented in Figure 4.7 for changing the matching condition for our resonator from the HOA2.0 to the Peregrine trap/package system. The procedure is to first find the matching condition at room temperature and use that as your zero
position. Then the micrometer screw (without scale) is detuned from resonance while the turns are counted.

Intuitively we know that objects shrink at low temperatures so the two coils should tend to move away from each other when cold. We chose to detune by screwing inward, that is pushing the antenna coil farther towards the center of the resonator. We make this movement in hopes that when things shrink, they will move away from each other again and thus move closer to the matching condition. While the dominant reason for the matching condition shift may actually be a change in series resistance the change in matching necessary to compensate for this just happens to follow this intuitive mechanical logic, and so whatever the reason, we find in practice that we need to detune by moving matching the coil further inward when compared to the room temperature impedance matching condition.

We can pick a first movement and count turns we take and then collect data during the cooldown process, the resonance dip is tracked and we find at what temperature the best match occurs. This first cooldown is seen as a yellow trace in Figure 4.7(a) with a corresponding point in 4.7(b). We can use these two points to make an educated guess about how much movement must occur for a good match, although had we only used these two points we could have overturned by 5. It is best to use data from two cool-downs to more accurately predict the matching condition. We can fit a model to the two cool-downs (orange and cyan) in Figure 4.7 and then fit a linear model to these points. Drawing a line at our operating temperature (here 8.5 K with high voltage RF operation) we can predict how many turns from the room temperature match we require to match the resonator. In this example we predicted that 32 turns was correct (dotted blue line) and the magenta trace was our final match. The prediction was good enough for operation, however perhaps 1-2 extra turns would improve it further. At the -20 dB level or better of reflected power we
4.3.4 Imaging Lens Alignment Stage

Our imaging lens stage is essentially a large kinematic mount which is itself x/y translatable. To achieve this the design has three parts. The lowermost frame has mounted micrometer screws which are set to an axis 45° from the vertical. An upper stage piece which will be translated is loaded down by four springs which are mounted vertically. The micrometers push against the upper stage, the stage is constrained from rotating side-to-side by being pinned between two plates whose surface normal is orthogonal to the loading springs. This portion of the stage is seen in Figure 4.8(a) and is what is bolted to the 90 K platform. The upper movable stage has three Attocube cryogenic stages bolted to it. The imaging lens itself is pinned to a copper plate via screws going into a sleeve mount. This can be seen attached fully in Figures 4.8(b) and (c).
Figure 4.8: A diagram displaying the imaging lens translation stage design (a) We see the X/Y translation mechanism with one lateral constraining plate removed (so the loading springs can be seen), the micrometers push at a 45° angle with respect to the springs in orthogonal directions. (b) Attocube stages are then mounted to this translating platform and they all linearly push a plate to which the lens is mounted causing it to tilt about the lens center, translating all three stages will cause z-translation of the lens. The operation of the tip/tilt mechanism is identical to a 3-axis kinematic mirror mount. Springs act to pull the plate inward while the stages individually push against these springs to cause tilting. (c) The lens itself is mounted via strews rigidly to the front plate, and the springs are held to this plate by long pins.

The copper plate is attached to the cryogenic stages via a spring loading it inwards and a ball ended pin pushing outwards (4.8(b)). There is an attachment point at three corners of the plate (for each stage). The cryogenic stages move and they are able to tilt the lens with both pitch and yaw. By moving all three stages we can translate the lens along the imaging axis. Attached to the copper plate is also a copper mesh which runs directly down to the 90 K platform. This mesh serves as
an additional, stronger path to cool the copper plate and lens. Without this path it
would only depend upon the loading springs and ball pins as a thermal path.

4.4 Optomechanical Design

The optomechanical design paradigm for the system advanced with this iteration of
the project. Key in this was the change in orientation of the trap and trap PCB with
respect to the vertical. Previously the trap PCB stack and CPGA socket was oriented
with its plane normal pointing along gravity. The trap PCB was rigidly attached to
the cold finger sample amount and so this stack was critical for determining the
ion location. The sample mount in the second generation system instead determines
completely the ion location in all three dimensions and the PCB stack is unimportant.
These PCB/CPGA tolerances are extremely loose and thus unreliable for location
referencing. The design of the lid and sample mount thus enabled the removal of the
vertical degree of freedom on the plate heights. We could design the entire underlying
structure with a set height as long as enough space was provided for conceivable
optical elements which may be needed. This means that the entire optical breadboard
structure would have a very good parallelism with respect to the cryostat and optical
table. The height of our optical components are known be within machine tolerances.
Many more support structures could be included with little overhead to drastically
increase the rigidity of the entire underlying structure.

The quality of the quantum logic gates imposed on the ion qubits will depend
strongly on the stability of the laser beams seen by the qubits that drive the gates.
Given that all laser beams are delivered to the ions in the cryostat through optical
fibers, the optomechanical structure to accommodate beam path between the fiber

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and the ion requires care in reducing vibration and drift. The structure is placed in the main experimental enclosure and consists of a monolithic main frame onto which the various ‘downstream’ optical blocks are assembled. All optical blocks and base plates are machined from MIC-6 (Tool and Jig) cast aluminum plates. This material is chosen for its low stress and good flatness specifications. The granular structure of cast aluminum allows for high speed machining while minimizing distortion which can arise when working with rolled aluminum.

The entire Montana Instruments cryostat sample chamber and cryocooler module rest atop a 3.5 cm thick base plate which is 0.89 m long on a side. This makes up the foundation of the experimental setup (shown in Figure 4.9), serving as a mechanical substrate for the ion trap chamber and all downstream optics such that vibrations between them are common-mode. A series of evenly located 1.5 inch thick steel posts are used to suspend a second 1-inch thick optical breadboard at a height for the ion addressing optical block setups to operate. This custom optical breadboard has a grid of 1/4-20 tapped holes with a 2 inch pitch, and each tapped hole has two accompanying 2 mm precision holes for dowel pins. In the space between the base plate and the optical breadboard we have room for routing various cables and vacuum lines to and from the cryostat.

### 4.4.1 Optical Block Design Paradigm

Each unit of optical functionality is implemented as an optical block, which takes the form of a block of aluminum plate machined on top to locate and fasten optical components in space, similar to the unit shown in Fig. 4.11(b). Each optical block is installed onto the optical breadboard using 2 mm dowel pins for precision positioning. By fabricating the blocks from MIC-6 aluminum, good flat-to-flat interfaces between

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Figure 4.9: The optomechanical enclosure which creates the underlying structure upon which all experimental optics rests. Upon the custom 2mm dowel pin and 1/4-20 tapped hole gridded custom breadboard rests the CW beam delivery setup, the ion imaging setup, the MEMs individual beam optics, the Raman global beam optics, and the ablation delivery optics. All vibrations between the ion and these free-space optical modules will be common-mode.

the breadboard and the optical blocks are ensured. Optical components themselves are directly mounted on the aluminum modules with dowel pins for positioning. Some optical modules are elevated with steel posts, where a different height in the enclosure is needed. Custom-designed corner pieces are used to join the large base plate with the optical breadboard, and to attach side and top panels for the housing. Each optical block as well as the entire optomechanical structure can be enclosed with a metallic cover to eliminate air turbulence and reduce acoustic vibration.

This modular design strategy combines the stability of monolithic machined pieces
with the experimental flexibility needed to re-design and improve optical functions. Each optical block can be replaced and assembled into the optical breadboard with precision alignment with respect to other blocks. Fig. 4.9 shows an image of the optical blocks placed on the optical breadboard, labelled to highlight their optical functions. The optical blocks include: (1) an imaging plate for state fluorescence detection designed to image photons collected from each ion onto a multi-mode fiber in a fiber array, (2) a CW beam delivery plate which co-propagates all trapping lasers and brings them into the ion trap at a downward 45° angle elevated to access the window on the top cover of the cryostat, (3) two Raman optical blocks devoted to the counter-propagating Raman assembly, one for bringing in a global beam and the other to integrate the MEMs-based beam steering system for two individual beams, and (4) two optical blocks for the ablation laser (Q-switched Nd:YAG laser), one exterior to the enclosure upon which the laser head is mounted with steering mirrors and a periscope which brings this beam up to the second module which delivers the ablation beam to the atomic source in the trap cryopackage.

We make use of many compact custom and off-the-shelf glue-in mounts for our optical components. We prefer to use 302-3M fiber epoxy with a room temperature cure to bond our optical components to mounts. We have used many glue in mounts in enclosed spaces bonded in this way and haven’t had UV degradation issues arise after years of use.

**Dowel Hole Tolerances**

The optical block paradigm depends upon the proper choice of dowel hole tolerance. Our design is based off of an M2 (2 mm diameter) dowel pin system, and this is a convenient choice because the ThorLabs Polaris line of optical components uses M2 dowels for location. The proper designation of the dowel hole tolerances on a
mechanical drawing will help a machinist know how to ream dowel holes which will have the proper functional characteristics. Finding a copy of Machinery’s Handbook or an adjacent text [OJ14, PM00] and looking up the chapter on allowances and tolerances is essential for the designer to understand the proper way to size a dowel hole for a given purpose. For our purpose we desire a tight locational slip fit, not a press fit which will embed the dowel pins tightly. There is a system for notating hole and shaft tolerences on all drafting tools which you will find discussed in these texts. This system uses a letter followed by a number with increasing numbers being more loose. We have used hole tolerances H8-H11, with H11 being on the looser end and H8 having more risk of a tight fit. We know that H7 is too small and the fit will depend on the particular dowel pin. All of these ‘H’ fits are unilateral, that is they will call for holes no smaller than the nominal diameter with varying levels of looseness which will depend upon the hole diameter. Hole tolerances called in this range have been used with success in various projects to obtain holes with the property we desire. If for some reason a set of 2 mm dowel holes has been machined and improperly tolerenced and thus a true 2 mm dowel cannot fit without extreme force, an elegant solution to save this part is to purchase some 5/64” dowel pins which are 1.984 mm in diameter. These dowel pins will solve your problem and allow you to use the part. We typically chose to use approximately 6 mm long dowel pins, with a 4 mm deep hole in the larger underlying part, and a 2 mm extrusion into the smaller part to be located.

The x/y location of our dowel holes is usually allowed to be the standard 5 mil tolerance. It is sufficient, even for large optical modules, to utilize two dowel holes close to each other for locating the module and this practice will prevent binding or jams. When we install a plate we typically will just use two dowel holes which immediately surround a single tapped hole of our choosing. This will locate the
entire plate. Because of this we have never run into problems when specifying the standard 5 mil machine tolerances for the dowel holes locations. Locating large plates in this way still affords good repeatability of laser pointing to our ion location when removing and replacing the optical module.

4.4.2 Thermal Performance

The cryostat used in this work incorporates a Sumitomo SRDK-101E cryocooler. Sample space is rigidly attached to the base, and a unique cold finger design provides the thermal link between the cryocooler and the sample space while isolating the transfer of vibrations (details are Montana Instruments proprietary information).

Table 4.1: A summary of the temperature readings on the 5 K platform and in the sample mount for different experimental configurations. After measuring the base temperature without any heat load to the setup, we investigate the effects of extra heat loads in subsequent measurements. We require 96 DC wires for operating with surface traps in our setup. To minimize heat transfer to 5 K stage, these wires are made from low thermal conductance metal. Additionally, we need to make a RF connection between helical resonator and the trap PCB, which provides a heat transfer path from 90 K to 5 K stage. These two connections increases our operating temperature by about 1 K. In order to trap Yb+ ions and work at 2-3 MHz radial motional frequencies, we need to apply RF with 200-300 V amplitude.

<table>
<thead>
<tr>
<th>Temperature Readings (Kelvin)</th>
<th>Configuration</th>
<th>Platform</th>
<th>Sample</th>
</tr>
</thead>
<tbody>
<tr>
<td>Base temperatures</td>
<td>4.99</td>
<td>5.70</td>
<td></td>
</tr>
<tr>
<td>DC cables connected</td>
<td>6.05</td>
<td>6.56</td>
<td></td>
</tr>
<tr>
<td>DC and RF cables connected</td>
<td>6.20</td>
<td>6.75</td>
<td></td>
</tr>
<tr>
<td>RF on (EPICS trap)</td>
<td>6.51</td>
<td>7.05</td>
<td></td>
</tr>
<tr>
<td>RF on (HOA 2 Series)</td>
<td>8.16</td>
<td>8.79</td>
<td></td>
</tr>
</tbody>
</table>

In order to operate the trap cryopackage below 10 K to benefit from cryopumping by the carbon getter, the heat load to the 5 K platform should be kept as low as

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possible. The main heat load in our setup comes from the DC and RF wires, and the applied RF signal for trapping. Since the wires originate from room temperature, special care should be taken to thermally lag these wires at the 90 K stage to minimize heat transfer to the 5 K sample mount. A summary of operating temperatures for different configurations is given in Table 4.1. The measurement shows that the thermal dissipation from the RF signal is a significant contributor to the overall heat load. If we compare the last two entries in Table 4.1 we see a large difference between the EPICS trap and the HOA 2 series of traps. The single layer gold on fused silica fabrication of the EPICS trap has a lower trap capacitence than multilayer silicon traps which leads to a lower RF dissipation.

4.4.3 Vibration Measurement

Mechanical vibrations between the ions in the trap and the Raman lasers cause infidelity in quantum logic gate operations. Conventional closed cycle cryostats feature large mechanical motion of the cold finger sample mounts relative to the lab frame due to the expander motion of the cold head [VWB+13b]. We designed an optical interferometer apparatus for characterizing the relative vibrations of our optical modules mounted on the breadboard with respect to the cold sample mount. We constructed an optical module that includes both arms of a Michelson interferometer (similar to Ref. [DHH+21]), where the second arm reflects off of a mirror mounted in place of the ion trap on the cold finger sample mount (see Fig. 4.10(a) inset). We measured the vibrations in the horizontal plane with this apparatus. The direction is critical in maintaining optical phase of the Raman beams seen by the ions. We used 370 nm light to measure the interference signal on a photodiode and inverted it into a motional displacement in space. We observe that the displacement due to

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mechanical vibration is approximately 17 nm peak-to-peak with an RMS deviation of 2.4 nm. This displacement is small compared to the optical wavelength of the Raman beam (355 nm), and should provide stable beams at the ion location to drive high fidelity gates. The in-plane displacement and its power spectral density is shown in Figure 4.10(a) and (b) respectively. We see that the dominant peak is at 1.42 Hz corresponding to the mechanical pumping frequency of the cryostat, and some peaks at 120 and 130 Hz due to fans nearby.

**Figure 4.10**: Results of the vibration measurement. A Michelson interferometer is used to measure the vibration between optical modules and the cold finger. (a) In our system we experience approximately 17 nm peak-to-peak vibration with an RMS deviation of 2.4 nm as can be seen in this time trace. (b) We compute the power spectral density of vibrations, the strongest peak is at 1.42 Hz which is the pumping frequency. We also see a strong peaks at 120 Hz and 130 Hz due to fans nearby. This data was collected by Junki Kim and the apparatus designed by Zhubing Jia and appears first in [SIJ+21] © 2021 IEEE

### 4.4.4 Experimental Table Design\(^7\)

All upstream optics on our table are also designed as optical block module assemblies. Each plate includes a water-cooled base-plate with dowel pins and tapped holes for

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mounting individual optical blocks sub-modules. Each unit of optical functionality can be boxed up to create a controlled environment. The arrangement of these units on the optical table can be seen in Figure 4.11(a). The modules included in our setup are a CW-laser frequency stabilization cavity, a CW-laser modulation and fiber coupling plate, a Raman laser modulation setup, we call the Raman upstream. In principle these units could be stacked vertically into a rack-style enclosure. Once aligned their pointing, in practice, never needs adjustment. Our fiber couplings do not drift at all in years. In a recent case the entire frequency stabilization cavity module was disconnected from water cooling and moved to the opposite side of the table without losing alignment to the cavity mode.
**Figure 4.11:** Our optical table setup and system design philosophy. (a) The entire ion trap experiment easily fits onto a single 4 by 8 ft optical table. The setup includes a CW laser launch and modulation plate, a large Raman laser modulation setup which is integrated into a base for the Paladin 355 nm laser, a transfer cavity setup for frequency stabilization, and a large monolithic optical enclosure for the cryostat and main experimental optics. (b) As an example of our optical design strategy, we present the internal design of the continuous wave (CW) laser launch box. It has a larger water cooled base upon which a few optical blocks are mounted. We use 399 nm light for photoionization \([\text{VAS}^+19]\), 370 nm for \(^{171}\text{Yb}^+\) qubit operations, and 935 nm for re-pumping. Availability of fiber products at 935 nm, such as fiber modulators and beamsplitters, simplifies the design for this beam. We use free space splitters, EOM’s and AOM’s at 370 nm. These custom machined plates allow for compact and highly stable optical assemblies. This CW module measures 48.6 cm (W) × 36.2 cm (D) × 9.3 cm (H). Appearing first in [SIJ^+21] © 2021 IEEE
Continuous-Wave Modulation Module

Fig. 4.11(b) shows the CW lasers needed for trapping, cooling, initializing, and measuring the $^{171}$Yb$^+$ hyperfine qubits [OYM+07] used in our experiments. The 399 nm laser (AO Sense. Inc.) is used for resonant photoexcitation of the neutral Yb atoms for loading the trap. For efficient, isotope-selective ablation loading, a second photoionization step is required to excite the electron to the continuum. Here, we use the intense beam from the 355 nm Raman laser to achieve this goal. The 370 nm laser is resonant with the $^2S_{1/2}$ to $^2P_{1/2}$ transition in Yb$^+$, and is used for Doppler cooling, optical pumping for qubit initialization, and exciting resonant fluorescence for qubit state detection. The Doppler cooling requires a modulated sideband at 14.7 GHz to pump out the dark state in the ground state manifold, and optical pumping requires a 2.1 GHz frequency shift. These are provided by EOMs in the beam path. AOMs are used to provide shutter functionality and fine tuning of the cooling and detection frequency. These AOMs can also be used to intensity stabilize our downstream free space beams which are incident upon the ion. The beam is divided up and fiber-coupled to be sent to either the experiment, the frequency stabilization setup, or to a wavemeter for frequency monitoring. The 935 nm laser is needed to re-pump the atomic population from the meta-stable $^2D_{3/2}$ states back to the $^2S_{1/2}$-$^2P_{1/2}$ manifold. The optical elements are positioned precisely on the base plate using dowel pins according to a layout generated using computer-aided design (CAD) tools. The base plate is temperature-stabilized with a water chiller, and enclosed to minimize air turbulence along the optical beam paths. This optical setup is stable enough that the single-mode fiber coupling alignments have not required adjustment for over a year of system operation.

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8Section is adapted from [SIJ+21] where it first appears some parts are subject to © 2021 IEEE
Frequency Stabilization System

The frequency of the CW lasers driving atomic transitions (370 nm and 935 nm lasers) must be stabilized to the MHz level or below. This is achieved by taking some of the laser output and locking it to an optical cavity using the Pound-Drever-Hall technique [DHK+83]. The length of the optical cavity itself is stabilized to an absolute frequency reference, which is a 780 nm laser stabilized against Rb vapor using Doppler-free spectroscopy techniques. The setup is designed and constructed in a similar approach to the CW laser modulation plate shown in Fig. 4.11(b).

4.4.5 Raman Modulation Plate

Our upstream Raman modulation plate has a large underlying water-cooled base plate. The base plate was made larger than initially needed to accommodate future expansion. Attached to this base plate is a dowel-pin aligned base plate for resting and locating the Coherent Paladin laser. The laser has feet which are allowed to rest inside a V-groove on this plate so that the laser will come to a similar equilibrium between removal and placement. 3D printed adapters bring HVAC cooled air through a long duct directly into the lasers air-cooling port, another matching duct removes this air from the experiment and deposits By turning this waveplate we can either bring power to the experiment or to a beam-dump and photonic crystal fiber (PCF) curing station. Before dumping polarizing beam splitter (PBS) we sample some light which heads to a fast photodiode used to run the repetition rate lock. It goes then through an AOM used as a shutter for all beams. The diffracted mode from this AOM will be then split using a HWP and PBS between the global beam and the two individual beams. After this a second, optional, split can be made to send some light the opposite direction for expansion into another individual beam plate. This option
is not currently enabled.

**Figure 4.12:** A diagram of the Raman upstream modulation plate. We place the 355 nm Raman laser on a locating plate which aligns the laser by V-grooves which mate with the laser feet. The system has a large underlying water-cooled base plate to which we mount two optical block modules, one for modulating the global beam, the other for modulating the individual beams.

On the global beam side we then have another optional AOM which acts as a secondary shutter and could be used to run an intensity stabilization loop for the global beam. This AOM is not in use currently and we lock using the first shutter.
AOM. The global beam beyond this wraps around and is imaged through a telescope with a Brimrose AOM in the middle. The Brimrose was chosen for the global beam modulation because of its fast response time. The light then passes through a delay stage used for path length matching and is coupled into a photonic crystal fiber. On the individual side we split the beam into two beams each passing through its own double pass AOM setup. Both individual beams go through their own delay stage for individual path length compensation and finally exit to two fiber coupled ports.

### 4.4.6 Free-space Optical modules

Inside the immediate optical setup we have several previously mentioned free-space optical setups for beam delivery and imaging. The entire immediate setup with these free-space modules can be seen in Figure 4.9. Here the important and unique elements are described in more detail.

**Ion imaging Plate**

In Figure 4.13 we see the ion fluorescence imaging setup. The infinite conjugate light from the ion can be imaged to a chosen magnification with an off-the-shelf plano-convex lens. Here we use a lens with $f=250$ mm. First there is a mechanical flip mirror which is used to decide whether to end fluorescence into the 50x reimaging path or the CCD imaging path with 20x magnification. In the CCD imaging path we have a 1” pellicle which is used to send in diffuse 375 nm LED light for imaging the trap. We use a PointGrey CCD camera for imaging the trap and long ion chains at this level and this is mounted on a linear translation stage. The light which goes into the reimaging lens gets wrapped around and eventually goes through the reimaging which is mounted to a linear stage for focusing. The reimaging lens is a four-lens stack.
(power split doublet relay) with 2.5x magnification. Optionally there is a magnetic mounted 1/2” inch pellicle which can be installed to direct light simultaneously into either a CCD at 50x or a PMT/Fiber array with equal path length. This feature is useful for aligning the PMT and adjusting for aberration control. We mount a fiber array with 250 µm pitch to a 3-axis stage. The axes of this stage both along the ion chain and in the trap vertical direction can be automated with a closed loop stage.

![Diagram of ion fluorescence setup](image)

**Figure 4.13:** A diagram showing the path for imaging ion fluorescence. The infinite conjugate light from the experiment is sent through a simple plano-convex lens to be focused to an image plane. A 1” pellicle is used to send some diffuse 375 nm illumination light from an LED into the setup. A mechanical flip mirror is used to choose between a first image plane with 20x magnification where we place a CCD camera, and a beam path where the ion light is reimaged to 50x. The reimaging lens can be adjusted with a linear stage. Optionally a magnetic mounted 1/2” pellicle can be installed to split the 50x light between a CCD camera and the PMT leg. This is helpful for PMT alignment, but not always included. The PMT/Fiber array can be translated by a large XYZ stage. Closed loop actuators can be installed in place of manual micrometers.
Global Beam Plate

Figure 4.14: We have a simple global beam plate design. Shown here in blue the exiting global beam passes through full polarization control before encountering a 1,000,000:1 polarization filter which purifies it to a linear polarization. The polarization control waveplates are for maximizing the throughput here. There is a beam sampler which can direct sampled light onto a photodiode used for intensity locking. This is followed by a 1:1 telescope for pin-hole filtering, and optional cylindrical beam shaping optics. Then there is a telescope for expanding the beam before it reaches the final focusing lens with \( f = 150 \text{ mm} \). In this project we did not utilize the cylindrical beam optics. Shown on this plate is also an optional 370 nm beam path for future planned EIT cooling, but this function is yet untested as of now, and the CAD optics as shown here need to be adjusted (D mirror mounts reversed) in order to accommodate this beam path.

Our global beam plate was designed to shape the beam into a 50 \( \mu \text{m} \) by 10 \( \mu \text{m} \) elliptical beam focused onto the ion chain or to a 10 \( \mu \text{m} \) by 10 \( \mu \text{m} \) round beam waist. In this project we have not extensively used the elliptical beam capabilities except for a brief period of testing gates on 5 ions. As seen in Figure 4.14 we have full polarization control after the fiber collimator but before passing through a
1,000,000:1 polarization filter. This turns drifts in polarization into intensity drifts and the polarization control waveplates are there to maximize the throughput of the filter. There is a beam sampler (uncoated glass) which takes a small sample of the beam for sending to a photodiode whose signal we use to lock the beam intensity. Following this stage we have a 1:1 telescope with a 60 \( \mu \text{m} \) filtering pin-hole which can optionally be installed to clean up the mode. We have two cylindrical optics which can optionally be added to shape one axis of the beam into the 50 \( \mu \text{m} \) waist for longer ion chains. Following this there are two more plano-convex \( f=15 \text{ mm} \) and \( f=100 \text{ mm} \) lenses to expand the collimated beam before the final focusing optic which has a 150 mm focal length.

**Microelectromechanical Mirror based Individual Addressing Beam Plate**

In order to accommodate individual addressing of longer ion chains we utilize Microelectromechanical (MEMs) tilting mirror based individual addressing system. The functional principles of the beam steering system are outlined in more detail in [Cra16, KCF\textsuperscript{+}17]. Here we have adapted the optical design, originally prototyped with off-the-shelf optomechanical components, into our small form-factor optical block module format. The optical block has two stages, the first of which is the MEMs beam steering system and the second of which recombines the beams vertically as well as does elliptical beam shaping through cylindrical optics telescopes. At a high level the beam steering portions functionality can be understood by considering two individual beams which are vertically apart in space but horizontally aligned. They are brought through an initial focusing lens which images them onto a pair of 0.3 mm diameter MEMs mirrors which are displaced vertically apart by 666 \( \mu \text{m} \). After reflecting from this first pair of MEMs mirrors (first bounce) a concave mirror will reimage them back to a second pair of mirrors 2 mm displaced horizon-
tally (second bounce). A D-shaped mirror, as well as an X/Y translation stage for the MEMs chip itself are used to align the first bounce while imaging the beam (and mirrors) by replacing the concave mirror with a lens. The second bounce can be then aligned by tilting the concave mirror slightly while imaging after the second bounce. The z-position of the concave mirror is adjusted to make sure that as the first set of mirrors is tilted, the beam will not walk off the second set of mirrors. Finally a Fourier lens is placed which translates the small mirror tilts into linear shifts at the ion plane.
Figure 4.15: Our MEMs individual addressing plate is created to translate the optical design established in [Cra16, KCF⁺17]. The setup has two stages through which the two individual beams pass. The first stage is a folded MEMs beam steering system used in all MEMs setups such [WCF⁺20]. In this stage two beams are vertically offset (to utilize different rows of MEMs mirrors) and are imaged into two bounces on the MEMs each. If the system is setup correctly steering from one mirror will not cause the beam to walk on the next and tilts of the beams caused by the mirrors will be turned into shifts on the Fourier plane after being transformed by the Fourier lens. The second half of the individual beam optics have the dual purpose of transforming this Fourier plane into a second real image plane in which the beams will be elliptical beam and recombined to be copropagating.

We chose a MEMs configuration in which both sets of steering mirrors will tilt in the same direction to increase steering range rather than in opposite directions. The system will have a steering range of 10-12 ions depending upon how close to the HV mirror limit (150 V) we will approach and the ion spacing used. The second half of the MEMs plate includes a beam shaping telescope for translating the approximately 100 µm by 100 µm beam at the Fourier plane into an elliptical beam with dimensions
of 7.12 \( \mu m \) by 48 \( \mu m \) waist radius. The shorter waist being along the ion chain and the wider being along the trap vertical. We accomplish this beam shaping by forming a telescope with respect to an initial 250 mm plano-convex lens. We have two sets of two cylindrical lenses (power split into doublets) which are oriented orthogonally. Thus a different telescope is formed for the two orthogonal axes. Within the space between the two orthogonal lenses we must recombine the beams (which were originally displaced vertically to utilize two different sets of MEMs mirrors) into a copropagating beam path.

We accomplish this within a short distance by bringing the upper beam directly through a non-polarizing beam splitter after turning it with a fixed mirror. In order to bring the lower beam up to the height of the upper beam within the smallest distance possible, the mirror angles needed to accomplish this are beyond the range of our mirror mounts adjustability. We instead fabricate two custom pedestals which are already cut such that the mirrors will be tilted to the optimal angle for recombination within the allotted distance and then the mirrors must only be tweaked finely in the end. After these two cylindrical lens pairs the beams at real image plane is demagnified by the Raman projection lens by 4x before reaching ion image plane. This lens can be adjusted in the cryostat by a Smaract linear stage.

Two versions of the final elliptical doublet lens (which controls the narrow diameter) were created. One produces a 2.7 \( \mu m \) by 12.0 \( \mu m \) beam waist with a predicted 0.03\%-0.12\% crosstalk at 5 \( \mu m \) away. This design exists but has not been tested on the ions. The alternative design has an expected beam dimensions at the ion of 1.79 \( \mu m \) by 12.0 \( \mu m \) with a predicted crosstalk of 0.72\%-1.2\%. In practice when measured using the ion we have best case crosstalk figures of 1\% and the beam, measured by ion shuttling, has a waist radius at the ion of 1.78 \( \mu m \) in accordance with the design specification.
Continuous-Wave Laser Delivery

Our CW beam delivery optics are simple. We deliver all CW lasers through a single copropagated beam coming from one side. We identify dichoric beamsplitters which will serve to reflect the shorter wavelength light and pass long wavelength light and we arrange the beams in a cascade such that there are always enough degrees of freedom to completely copropagate the shortwave light with the previous transmitted longwave light. We use the longest wavelength light (935 nm) as the pilot beam initially and copropagate all of the shorter wavelengths with it by adjustment of the reflected dichoric and one previous kinematic mount. We aim to copropagate the beams roughly before passing through the final focusing lens. Our CW beams enter the chamber at a downward $45^\circ$ angle which we create with a separate mounted piece containing both the final lens ($f=150$ mm achromat) and translation stage for this lens. The beam is turned into the chamber by a final 1” kinematic mount. This kinematic mount is dowel pin aligned and can be removed at will to send the beam straight into another stage mounted beam profiler placed at the ion location. We copropagate all of the beams carefully between the beam waists and Rayleigh range using a beam profiler mounted to this second stage at the ion location. This setup provides for an easy beam pointing diagnostic. We find that upon removing and replacing the dowel pin mounted kinematic turning mirror mount we have good enough repeatability to still be pointed at the ion optimally.

Originally our design featured mounting locations for all cooling and repump beams for cotrapping ytterbium and barium ions. We have since upgraded to a new design which eliminates the barium capability in favor of elliptical 369.5 nm cooling beam and EIT cooling which are created using cylindrical lens telescopes. These features have not yet been tested and we use a 15 $\mu$m circular beam for the 369.5 nm
cooling beam.

**Ablation Setup**

We use a Continuum branded ND:Yag Q-switched pulse laser operating with 532 nm output for our ablation source. The laser head is mounted beside the setup but brought in in free-space after attenuation. A periscope is used to bring the beam to the top of our setup and it is focused through one lens to bring it to the ablation threshold. The details of ablation loading are discussed in Appendix A. The ablation beam passes through a beam splitter which allows the image of the ablation target to be split to a long-working-distance camera. We can image the ablation laser at a low power and high rep rate and make sure that it is hitting the target, 532 nm light makes this easy (we do not need to copropagate). The ablation intensity can be increased for trapping chains quickly, or operated at low fluence to minimize trap contamination for trapping short ion chains. For trapping long chains (with a single isotope) it is ideal to operate in the regime of low 399 nm photoionization power, tuned closely to the resonant line, and high 355 nm 2nd photon power.

**Photonic Crystal Fibers**

We utilize a modified version of the NIST Photonic crystal fiber (PCF) recipe and make then cure our fibers in-house. We have worked with NKT Photonics LMA-10-UV and LMA-PM-10 which were sent for hydrogen loading by O/E Land a Canadian company. In this recipe we use temporary fiber clamps to cure a long section of fiber whose facets have been prepared by simple cleaving then assembling using a jacketless version of the NIST PCF recipe. This allows use to avoid the use of UV cured epoxy and 5-minute epoxy, both potential sources of volatile organic contamination of the fiber facets. Once we cure a long stretch of fiber there is no longer any more
time urgency related to the hydrogen loss before the curing process. We then cut this fiber up into the desired lengths and connectorize it with a standard polished endcap, which we can then re-polish as needed. This recipe can be done with jacket or without, we have had good lifetime from our individual fibers which are unjacketed and entirely built with 302-3M as the only adhesive. The furcation tube fiber jackets are also a source of organic outgassing, the experimentalist should consider whether the convenience of these fiber jackets is orth the possible organic contamination when working with high power UV.

Additionally our original fusion splicer (which was from the NIST recipe) had broken and this part was recommended by the vendor to be succeeded by the Fitel S185-LDF Large Diameter Fusion Splicer. We have built a recipe for collapse using this new Fusion splicer, the key difference being that in this one both sides of the splider are under full stage control, and by default, open closing the lid to the system, a clamp will come down to rest on the far end. For successfully collapsing a PCF it is important that the stress is minimized between the clamped end and the far end (with the location to be collapsed in between). In order to use this new system we must obtain instructions from the vendor for accessing supervisor mode, which enables full control of all stages. We then must lift the far clamp, and align the clamped end such that there is a minimal flexing or bending of the fiber across the gap. Failure to both lift the clamp of the far end, and the properly align the stages of the clamped end will result in distorted collapsed regions with a large angle.
Chapter 5

Trapped Ion Experiments

5.1 Early Trapping Results

5.1.1 Four-rod macroscopic trap

In order to verify that the helical resonator was working, that the pressure of the trapping environment was reasonable and that the ablation geometry worked for our compact cryopackage we opted to begin working with a four-rod macroscopic trap version which was prepared into a similar compact structure. Four tungsten rods were held in place by Macor holders inside a copper cryopackage lid. The tungsten rods were simply fed through to barrel connectors for running RF and DC voltages. An RF power input of 0.5 W was used with a helical resonator Q=200 with a resonant frequency of 25 MHz. Once thermal equilibrium was reached the cryostat had a base temperature of 7.5 K. The trap depth was estimated to be 1 eV with a dark lifetime of approximately 12 hours. The design is interesting and could serve as a template for a modular replaceable macroscopic trap in itself. In Figure 5.1 we see an image of the four rod trap cryopackage.

5.1.2 Fully-Vacuum Sealed ColdQuanta Packaged HOA2.1

A fully sealed trap cryopackage was delivered by Cold Quanta. It consisted of a titanium lid with brazed sapphire windows. An HOA 2.1 was packaged onto a modified ringframe CPGA with trap spacer and wirebonded at Sandia labs. The lid was
Figure 5.1: An image of the four rod trap cryopackage. (a) We see the device from the side tungsten rods are fed through Macor holders then barrel connectors allow the RF and DC connections. We also see both viewports and a setscrew which will plug the cavities for storing either ytterbium or carbon getter. (b) We see a view of the top imaging window where the tapped holes for cold finger sample mount attachment are visible as well as the four rods. (c) We see a close-up view of the four rod trapping region through the imaging port.

bonded to the trap package using an indium seal completed inside a large packaging station at ColdQuanta. The completed trap cryopackage was shipped through standard means and promptly installed in our cryogenic system.

Trapping in the package was swiftly completed, chains were able to be trapped with long lifetime; however three major shortcomings of the closed lid package led to their abandonment for this work. The most significant was the emergence of unpredictable leaks after temperature cycling in the cryostat. On the second round of cooldown when trapping in our new package the leaks caused poor ion behavior from a high rate of collisions. We measured a leak rate between 0.0005 and 0.0001 TorrL/sec when pumping down the cryo system. This was measured by using the two parallel windows of the cryo package as an interferometer. A change in the pressure inside the package causes a change in the index of refraction between the two windows, which appears like a length change on an interferometer. This method was used to verify leaks both in a fully sealed package which we trapped in as well as a newly made test package without trap installed.
Another shortcoming of this fully sealed vacuum package design were the materials chosen due to coefficient of thermal expansion (CTE) matching considerations which had unforeseen consequences. This motivated a design with a titanium lid and sapphire windows. Titanium is a problematic choice of material because its thermal conductivity at 8K is approximately $1 \text{ W m}^{-1} \text{K}^{-1}$ while high purity oxygen free copper has a thermal conductivity of $10^4 \text{ W m}^{-1} \text{K}^{-1}$ [Eki06]. The lid to the cyropackage is the most important thermal conductor in this experiment because heat generated in the CPGA/trap package which heats the lid must be conducted away into the larger cold finger sample mount. This ensures that the carbon getter which traps hydrogen and helium remains as cold as possible. The choice of material and manufacturing process which focused on CTE matching also caused sapphire to be used as the viewport. Sapphire is a birefringent material; however its use as an optical feedthrough can be aided by using ‘Z-cut’ sapphire in which the crystal is cut such that the c-axis of the crystal aligns with the optical axis of the window. This is usually sufficient to remove aberration caused by the material, however for our high NA imaging we are collecting light from a wide angular distribution of rays. Our 0.6 NA imaging takes advantage of light within a cone angle of $36^\circ$ and for the edge of the cone, a significant angle still exists between the c-axis, leading to an aberrated image. An image of ions viewed through these sapphire windows can be seen in Figure 5.2 where the ‘rings’ around the ion cannot be aligned away. For this type of window two different focus points can also be found when scanning the imaging lens on the z-axis.

At the time of original testing the source of the leaks in the compact cryopackage was unknown. After removal of the original package for a couple months of shelf storage, the package had become so pressurized that upon cooling no ions could be crystallized in the trap, although a hazy fluorescent cloud of ions could be formed. Subsequent to these initial tests of the fully sealed package we had decided that we
Figure 5.2: An image of 8 ions trapped in the ColdQuanta fully sealed vacuum package. Rings can be seen around the ions which is an artifact of transmitting a high NA image through sapphire windows.

were going to embrace a leaky package by designing these packages which are differentially pumped by incorporating a controlled leak. While the fully sealed package was verified in principle, more engineering work needs to be done on reliability and materials choice before the concept is ready for full-time use. However this detour to an open lid design became the basis for a new accessible way of constructing ion trap experiments with high performance.

5.2 Trap Heating Rate Characterization and Optimization

The rest of this chapter will be about characterizing the behavior of our open lid trap cryopackage design. Some figures of merit that we will explore include the heating rate, ion lifetime, and pressure. For the heating rate key results from a few different traps will be presented which demonstrates how technical noise can greatly influence heating rate measurement. At this time the heating rate still seems to be limited
by technical noise, which given the already low base heating rate, is a good sign indicating that yet more progress is to come. The motional heating rate of ions in the trap is of paramount importance for two-qubit gate fidelity as most entangling gates utilize the motional degrees of freedom to mediate the qubit interaction. The gate fidelity and circuit depth attainable in our system will also be effected by the heating rate $[WCF+20]$.

In the effort to measure the pressure experienced by our ions we will utilize zig-zag chains. There will be a brief foray into zig-zag chain physics and the scaling of the zig-zag threshold as a function of ion number in surface traps which may be useful for experimentalists attempting to deal with long ion chains.

### 5.2.1 Sandia EPICS Trapping results

We utilized the Sandia EPICS traps for our first heating rate characterization of the open-lid trap cryopackage design. The EPICS trap is gold fabricated on fused silica. The fused silica substrate has been coated with a dielectric mirror for use in ion trap cavity QED experiments. On the EPICS trap we were putting 24 dBm (0.25 W) of RF power at 56.2 MHz into the cryostat for trapping. The helical resonator-trap system had a Q of 200 and we observed a maximum sample temperature of 6.94 K. Using an RF tickle measurement radial trap frequencies were observed to be approximately 2.0 and 2.3 MHz.

The exposed dielectric of this trap made is extremely susceptible to charging from our 355 nm Raman laser. This charging caused the trap frequencies to drift rapidly with exposure. This effect made good side-band cooling impossible and working with motional sideband flopping problematic. However the decay of counterpropagating carrier transitions depends partially upon the ion temperature. With the assumption
of the thermal distribution of fock states we can fit the carrier transitions after a wait time after Doppler cooling. The longer the wait time, the faster the oscillations damp out. There results were used to get an order-of-magnitude estimate for the heating rate of this trap and the results were in the 10-60 quanta/ms range as can be seen in Figure 5.3.

![Figure 5.3](image)

**Figure 5.3:** A plot of the heating rate as measured in the EPICS trap via the method of carrier decay fitting. The heating rate is extremely large both due to a lack of internal filters on the trap PCB and an excessive amount of charging.

One should note that two key contributions to the poor heating rate with this sample are known; the first being that at this time there were no filter capacitors on the trap PCB. In future system upgrades RC filters with 2 kHz cutoff were added to each DC channel on the trap PCB. There was a filter box for all DC lines placed outside of the cryostat plugged into the DC feedthroughs, thus the relevant noise is being coupled between the cryostat exterior and the trap CPGA socket. The second problem was the extremely severe surface charging which disturbs the micromotion compensation condition leading to uncontrolled motional noise which also increases the decay of the carrier. Thus this measurement cannot be considered a good measurement of the intrinsic heating rate of this trap, but rather a measure of the large technical noise present and poorly contained 355 nm Raman addressing
beams which clipped upon a trap with poor clearance.

Another possible contributor to poor heating rates in this system could have been higher pressures. Collisions with the background gas can contribute to heating rate [CS14]. At this time we had not developed a mature procedure (in Chapter 4) for pumping the system before beginning cryocooling. If the cryopackage is not adequately evacuated then there can be a lot of low-energy collisions from material adsorbing and desorbing from the trap surface.

5.2.2 Sandia HOA2.1a Heating Rate Measurements

The HOA2.1 trap has several sub-versions which have different fabrication imperfections. The original HOA2.1 had a narrower-than-intended slot with poor metalization due to fabrication errors. There was an updated version which alleviated this issue called the HOA2.1a. When measurements were done on the original HOA2.1 without RC filters installed, heating rates were found to be in the 500-1000 quanta/second range. The original HOA2.1 issue still struggled with a high susceptibility to charging which would cause the motional frequencies to drift with exposure, and would make the ion position drift.

The heating rates presented were the best results obtained from the HOA2.1 series of traps (measured in HOA2.1a) after adding internal RC filters to the trap PCB. By comparing the height of the red and blue sideband lineshapes we can extract the thermal occupation [TKL+00]. This trap still experienced enough charging from Raman operations that the sideband cooled transition would drift away from the transition as measured without sideband cooling. The 355 nm dose received during sideband cooling is much greater than most experiments. This made sideband cooling to the ground state difficult as the resonant frequency was always subject to a time
A demonstration of the magnitude of this effect can be seen in Figure 5.4. In this experiment the same side-band cooling pulses were sent, but with a far off-resonant frequency so that no cooling was completed, just the same 355 nm light exposure. This experiment was completed as a frequency scan where the resonant frequency was tracked under low illumination after a wait time. This ‘fake sideband cooling’ (fake SB) experiment is a useful diagnostic for characterising the timescale and magnitude with which charging of the surface trap effects the trap frequencies. You can see from the figure that compared to the resonance frequency when no fake SB cooling is applied there is a 4 kHz shift, this shift has a large component that will return to equilibrium on a 2 ms timescale, followed by a remaining 1.0-1.5 kHz shift which is much slower, even with wait times in excess of 30 ms there is still a 1 kHz shift (not shown).

Due to this large shift in the resonance frequency we opted to do full scans of the shape of the blue and red sideband peaks in order to find the peak heights and add error bars to this measurement. The sideband cooling we were able to achieve, while not reaching the ground state, did lower motional quanta below the Doppler limit. The scanned peaks can be seen in Figure 5.5 where longer wait times are coded in red fitting curves, the scatter of all points are shown in grey. With the height of the red and blue sideband peaks determined as $A_R$ and $A_B$ respectively we can calculate the motional quanta in the mode using the following equation:

$$\bar{n} = \frac{A_R}{A_B - A_R}$$  \hspace{1cm} (5.1)

With this measurement we have demonstrated a heating rate in line with the best room temperature HOA2.1 traps at cryogenic temperatures. The large susceptibility
Figure 5.4: A plot of the results from a ‘fake sideband cooling’ experiment. In this experiment all sideband cooling optical pulses are sent at a far-off-resonant frequency followed by a wait time which we scan. Each point comes from a curve fit to a following low intensity motional frequency scan. We can see that after exposure to the 355 nm light there is a > 3 kHz frequency shift which then damps out with waiting time on a timescale of 2 ms.

to trap charging causing low motional mode stability made doing good two qubit gates and sensitive operations involving the motional modes impossible.

5.2.3 Sandia HOA2.0 Heating Rate Measurements

Originally we sought to measure argon-ion-beam cleaned surface traps in a sealed ColdQuanta package. Argon-ion-beam systems are typically used for sputter deposition systems which to grow thin films[FFN+13]. Used in a deposition process a confined plasma is directed towards a ‘target’ material and collisions between argon and the target will eject the material to be sputtered towards the substrate. Utilized in a trap cleaning context the we direct the plasma beam towards the trap itself to

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1Section is adapted from [SIJ+21] where it first appears some parts are subject to © 2021 IEEE
Figure 5.5: In this experiment the ion is Doppler cooled, sideband cooled, a waiting time is allowed to elapse before pumping dark and doing a Raman interrogation experiment while scanning the frequency of this interrogation (a) A plot of the red sideband motional peak which was scanned after a heating wait time. Each set of data it fitted to a Gaussian lineshape from which the peak height is extracted. (b) A plot of the blue sideband motional peaks at different wait times.

Figure 5.6: The heating rate extracted from a 2.4 MHz higher frequency sideband of one ion trapped in the HOA2.1a trap.

sputter off the outer layers of the trap surface. This sputtering of the trap surface will clean the trap surface of contaminants.
The ColdQuanta trap packaging chamber includes the capability of in-situ surface cleaning. There has been some evidence that this treatment can improve heating rates [HCW+12b]. In order to facilitate this study we had old HOA2.0 traps recoated with a 500 nm gold layer. For comparison we first were going to install this new thick gold re-coated HOA2.0 into our standard open lid system and do good heating rate measurements. We were pleasantly surprised to find an almost order of magnitude improvement in the heating rate of this trap without doing the argon ion cleaning. Any further heating rate improvements would be marginal due to the other problems that we knew existed with the fully sealed trap cryopackages. We opted to abandon the surface cleaning experiment and the use of fully sealed packages.

Before presenting the findings it is important to explicitly note some changes to our process from the past which may have contributed to this conclusion. The process used for the HOA2.0 and Peregrine traps is the same (both of which have low heating rates) and for these packages we stopped doing any heating at all to the cryopackage after installing the lid. Previously after installing the lid the entire cryopackage was heated to soften the indium and help improve the mechanical bonding between the lid and the modified CPGA. This temperature was difficult to control precisely but is estimated to be in the 150-170°C range for as short a time as possible. For our final process there is no elevation of the trap temperature at any point in our handling of it, only a vacuum baking of the lid, getter, and constituent parts.

Another change implemented for both the HOA2.0 and the Peregrine traps was that our four 25-pin D-sub cables which deliver DC voltages from the digital-to-analog converter (DAC) box to the filter board PCB were wrapped tightly in thick ground straps. These ground straps were sealed to the cable using zip ties every 1.5 inch so that they form a complete extra layer of thick shielding. This adds to the theme that better shielding and filtering of the DC lines have a large helpful effect.
Finally it was noted as compared to the HOA2.1 series both the HOA2.0 and Peregrine were significantly less susceptible to charging.

The heating rate measurement scheme used is similar to that used in Mount et al. [TKL+00, MBB+13]. In surface electrode ion traps, one can generate the voltage solutions with arbitrary rotation of the trap principal axes. We measure the heating rates of the radial modes with a $45^\circ$ rotation of the trap axes with respect to the trap surface normal so that the Raman beams can interact with both radial modes. The ion motion is first cooled by Doppler cooling so that average motional excitation is $\leq 5$ quanta, followed by 30 rounds of resolved sideband cooling [LBMW03] to reach an average motional excitation of $\sim 0.09$ quanta in the radial modes. After waiting for a duration, we drive both the red and blue sideband transitions on resonance. We perform a fit to the experimental data with a simulated Hamiltonian of the blue side-band rotations with heating decoherence effects. We include an exponential decay envelope to simulate the motional dephasing effects. We use the fitted value of $\bar{n}$ of these sideband rotations to determine the average motional excitation (or temperature) at any given wait time.

We repeat this experiment for a variety of wait times, during which the surface noise influences the ion and causes heating. We scan the wait time to trace out the average motional excitation versus wait time and fit the results to establish a heating rate. The error of this heating rate is the standard deviation of the parameter estimate from the linear fit (Figure 5.8). We measured the heating rate at 15 different trap mode frequencies between 1.2 and 2.3 MHz, a plot of these results can be seen plotted in Figure 5.8. No apparent frequency trend was observed for this experiment. The lack of frequency trend could indicate that we are still limited by technical noise. The lowest measured heating rate was $5.0 \pm 0.8$ quanta/s at 1.82 MHz trap mode frequency,
**Figure 5.7:** Heating rate measurement at a 2.23 MHz motional frequency. Measured in the HOA2.0 trap with 500 nm gold recoating. We measure a heating rate of $14 \pm 1.1$ quanta/s with a maximum wait time of 90 ms. The error bars represent one standard deviation error on the parameter estimate from the non-linear curve fit to our side-band rotation simulation. Appearing first in [SIJ+21] © 2021 IEEE

**Figure 5.8:** A plot of all heating rate measurements completed on the HOA2.0 trap as a function of motional frequency. Over this range there is no clear trend in the data, no power law scaling was found. This may indicate that our heating rate is still limited by technical noise.
and the largest measured heating rate was $25\pm4.3$ quanta/s at 2.04 MHz. The mean heating rate of these measurements was 13 quanta/s with a standard deviation of 5.1 quanta/s. A representative plot at operational trap frequencies for two-qubit gates on the radial modes is presented in Figure 5.8, where we measure a heating rate of $14\pm1.1$ quanta/s. Improvements in the Raman intensity stability could reduce the measurement noise enough to detect a scaling trend with confidence; for this data the trend seems flat within the range of investigation.

### 5.2.4 Sandia Peregrine Trap Heating Rate Measurements

We have also tested the heating rates of the new Sandia Peregrine series of traps. Our work with the Peregrine focused more on optimizing Mølmer–Sørensen gates in the system, so we have not done a thorough survey of the heating rates as function of trap frequency. When checked at particular operating conditions we observe comparable or better performance than the average HOA2.0 trap.

We have measured the heating rate of the common mode of a single ion at 2.6 MHz trap frequency using the same method of fitting as with the HOA2.0 survey. We fit the blue sideband rotations via simulation to obtain a value for $\bar{n}$. An example of fitting to the rotations at two different wait times on tilt mode of two trapped ions and their corresponding fits can be seen in Figure 5.11. With this Peregrine data as compared to the previous work with the HOA2.0 we did rotations until we reached $6\pi$ times. Using longer rotations we were able to achieve a tighter errorbar on the simulation fits. The results can be seen in Figure 5.9.

We also measured the heating rates for two ion chains under the same sideband cooling scheme that was utilized for the MS gates. Meaning we sideband cool the tilt mode, which we use primarily for our gates, immediately before the heating wait time.
Figure 5.9: Heating rate measurement of one ion at a 2.6 MHz motional frequency. We measure a heating rate of $5.96 \pm 1.1$ quanta/s with a maximum wait time of 90 ms. The error bars represent one standard deviation error on the parameter estimate from the non-linear curve fit to our side-band rotation simulation.

Our sequential sideband cooling led to a slight heating in the previous mode cooled, such that in order to get as close to the ground state as possible in the mode used for gates we cooled it last. As can be seen in Figure 5.10(a) under these conditions we measure a $8.74 \pm 0.7$ quanta/s heating rate on the common mode of two ions, while in 5.10(b) we measured a $0.44 \pm 0.1$ quanta/s heating rate on the tilt mode. It is common that the higher order modes of ion chains experience lower heating rates as it is more difficult to excite these modes with spatially distant or uniform electronic field noise. We have also attempted in the past to measure the heating rate of the highest order zig-zag mode of a 6 ion chain. Long wait times in excess of 0.5 s are needed to notice a change.
Figure 5.10: (a) Heating rate measurement from the common mode of two ions which sat at 2.05 MHz, this is measured under the same conditions (sideband cooling routine) that we use for the gates, we chose to cool the tilt mode after the common mode, which gets effected by the cooling process leading to a higher base temperature. The measured heating rate is $8.74\pm0.7$ quanta/s (b) Heating Rate measurement of the tilt mode of a two ion chain at 1.98 MHz, where we center the FM solution for our two-qubit gates. The blue sideband rotations are barely effected until wait times of 200-500 ms are used. The heating rate is $0.44\pm0.1$ quanta/s.

5.2.5 Heating Rate Retrospective\textsuperscript{2}

Some of the bad heating rate results from earlier experiments were included to show that indeed simply going to cryogenics alone cannot solve a heating rate problem.

\textsuperscript{2}Section is adapted from [SIJ+21] where it first appears some parts are subject to © 2021 IEEE
The level of charging and technical noise experienced in the HOA2.1 seemed worse than in UHV chamber based room temperature traps in the same laboratory. We have not done a thorough enough survey of the Peregrine heating rates to determine if there exists an inverse power-law scaling as observed as in many other systems [DOS+06, BSC15]. As such we cannot conclude whether the peregrine heating rate
may be technical noise limited at this time. For the HOA2.0 despite much effort towards these ends, we were never able to find a clear inverse scaling relationship. Furthermore the landscape of heating rates observed indicates that technical noise still is the limitation. In either case we have seen an average heating rate around 13 quanta/s in HOA2.0 and have seen several Peregrine measurements less than 10 quanta/s heating rates.

In order to compare this to similar cryogenic systems, we need to normalize for the ion species and the trap frequency. Similar to Hite et al. [HCW+12a] we can calculate the electric field noise spectral density \( S_E(\omega) \) as

\[
S_E(\omega) = \frac{4m \hbar \omega}{q^2} \dot{n},
\]

where \( m \) is the ion mass and \( q \) is the elementary charge. When we multiply this by the motional frequency to obtain \( \omega S_E(\omega) \), we can compare our results to other low temperature surface traps at different trap-ion distances [HCW+12a]. Our noise spectral density \( \omega S_E(\omega) = 3.2 \times 10^{-7} \text{ V}^2/\text{m}^2 \), which lies an order of magnitude below the average line of preexisting experiments with 70 \( \mu \)m ion-trap distance at cryogenic temperatures. Based on careful considerations of the error contribution to entangling gates, we anticipate that this level of heating rate would not limit the gate fidelity up to 99.99\% [WCF+20].

Over time the largest single contributor to low heating rates over the course of many traps has been the filtering and shielding of electronic noise. Obtaining traps with less tendency to change with UV exposure helped stabilize the measurement. Additionally improvements in the chamber pumping routine before cryocooling is begun have played a role, although when pressure is an issue an abundance of F-state dark events can be observed.
5.3 Pressure Measurement and Zig-zag Chain Physics

It is desirable to measure the pressure inside our cryogenic package and compare it with other schemes. Measuring the pressure inside such a small package presents a unique challenge. Integrating an ion gauge or cold-cathode gauge is challenging, would be larger than the package itself, and negatively effect the pressures measured. The ion itself can be used as a pressure gauge in many different schemes. One possible method is to use the surface electrodes to establish a double well potential with an adjustable well height [AVN+20]. The double well method can be problematic when rare low energy collisions are sampled. The most probable collision energy of H$_2$ from an 8 K thermal distribution is 0.68 meV while at 300K this reservoir leads to collision energies of 25.8 meV. The double well potential must be manually balanced by adjusting compensation voltages until the probability of occupying either well is roughly equal. Furthermore these axial compensation voltages can drift overtime as the trap is subject to trap charging and discharging, electronic noise present on the DACs, and changing surface electric field noise due to adatom diffusion. This well balance can be made less sensitive by raising the potential barrier (which also decreases the hopping rate and makes experiments slow and difficult to calibrate), or by lowering the potential barrier (which makes the drift of well balance even more sensitive).

The rate of reordering can also be used to access the pressure. The potential barrier for reordering is on the scale of 1-2 meV based on both Aikyo et al. [AVN+20] and our own calculations. A correctly built trap cryopackage system does not see these reordering events at all. As such a method of using zig-zag shaped ion chains as a pressure measurement tool has been devised by Pagano et al. [PHK+18]. This method
is attractive because the potential barrier becomes insensitive to small changes in axial electric field noise and becomes a property of just the relative strengths of the radial and axial confinement potentials. Small changes in the axial electric field may translate the ion chain in space but has a negligible effect on the barrier height.

One complication of the method of [PHK18] is that they rely on a molecular dynamics simulation of the ion chain immediately after a collision event (and assume a fixed collision angle of $\pi$) to generate their probability of flipping as a function of threshold energy and temperature. We use a smaller ion chain and present a more simple analytical model to avoid this step of creating and calibrating a simulation of the chain cooling.

We also attempt to verify the zig-zag chain dynamics for differently numbered ion chains to see how the behavior diverges for surface traps from published literature. This measurement can also be used to understand the upper bound limits to good ion chain manipulation in surface ion traps.

### 5.3.1 Zig-zag Chain Physics

Ion trap quantum computing systems utilize linear crystals of ions confined to a Paul trap potential. Within the trap potential the central position of cooled ions will rest at an equilibrium determined by their mutual repulsion in balance with the confining potential of the trap. A linear ion crystal has well defined radial modes of motion which can be addressed with Raman beams and utilized for two qubit entangling gates. The work of [Sch93] showed that there exist sharp phase transitions between well-behaved linear crystals and bucked ‘zig-zag’ shaped crystals, then to helix shaped crystals with 3D structure, and finally to a disk shaped crystal. All of these transitions can be understood as being caused by and related to the relative strengths of axial
and radial confinement. Traditionally a new parameter $\alpha$ is introduced which is the ratio of the axial to radial confinement strengths \cite{Sch93}.

The harmonic portion of an ion trap potential can be characterized in simplified forms as:

$$U = \sum_{i=1}^{N} \sum_{j=i+1}^{N} \left[ \frac{q^2}{r_{ij}} + \frac{1}{2} m \left[ \omega_x^2 x_i^2 + \omega_y^2 y_i^2 + \omega_z^2 z_i^2 \right] \right]$$

Equation 5.3

Or more simply the sum of an approximately quadratic restoring force in each axis with a strength $\omega_{x,z,y}$ plus the force of each ion on every other ion. If we assume that the radial confinement forces $\omega_x \approx \omega_z$ we can reformulate the trapping potential as shown below in Equation 5.4

$$U = q^2 \sum_{i=1}^{N} \sum_{j=i+1}^{N} \left[ \frac{1}{r_{ij}} + \frac{1}{2d^3} m \left[ x_i^2 + z_i^2 + \alpha y_i^2 \right] \right]$$

Equation 5.4

$$\alpha \equiv \left( \frac{\omega_y}{\omega_{xz}} \right)^2, d^3 \equiv \frac{q^2}{m\omega_{xz}^2}$$

Equation 5.5

Scaled in this way it is claimed that the transition point between a 1D chain and a 2D zig-zag shape shall occur with a power law scaling behavior with the ion chain length, $N$, which goes as $\alpha_{\text{crit}} = cN^\beta$ \cite{Sch93}. In surface ion traps our potential is usually threefold anisotropic which implies that $\omega_x \neq \omega_z$. In such cases Schiffer et al. \cite{Sch93} claims that the same results are obtained for scaling; the smaller of the two radial axes shall govern the transition point and that the 2D plane of the zig-zag will face the weaker axis \cite{Sch93}. The work of \cite{ESG00} sought to experimentally verify this power-law which was predicted numerically and analytically in \cite{Sch93}. Enzer et al. used a macroscopic four-rod trap to measure the transition threshold for $N = 3 - 10$ ions. As of yet no one has measured an analogous scaling law for surface...
Paul traps to verify if the behavior is the same. We sought to verify the scaling of the Zig-zag threshold for our surface trap potential. The behavior may be different due to the much larger aharmonic component of the trapping potential, and the great flexibility of our axial DC potential.

5.3.2 Scaling of Zig-zag Formation in a Surface Trap

Our axial trap potential is realized in practice via the sum of a few DC solutions with a fixed strength. In this work we mainly utilized two different solutions on the HOA2.1, an approximately harmonic potential with minimal DC electrodes called ‘SimpleSol’ or the simple solution, and a more complicated solution with a more flat shape axially which helps to load and manipulate longer chains called the ‘FlatSol’. The trapping solution then in practice becomes a sum of this simple solution, a flat solution, and our micro-motion compensation voltages. The radial confinement is approximately set by the RF voltage level, while axial confinement is dominated by this SimpleSol parameter (which adjusts the strength of the harmonic component of the axial potential).

In order to translate this axial confinement strength parameter into a trapping frequency we have used a two-ion chain as a measurement tool for the axial and radial frequencies. We have trapped two ions into the potential used for zig-zag experiments and then varied the SimpleSol parameter and the strength of harmonic trapping, in order to increase the axial confinement. We do not have good axial coupling in this system and so we have used the common and tilt modes of the two ion system as a way to measure the axial confinement frequency and calculated the axial frequency as \( \nu_y = \sqrt{\nu_{x/z,comm}^2 - \nu_{x/z,tilt}^2} \) [Hay12]. Raman motional spectroscopy does not work well to measure the radial frequencies of a zig-zag chain. The results
of these measurements are seen in Figures 5.12, 5.13, 5.14. The axial common mode can be predicted identically no matter which pair of radial modes are used. There is one outlier point which is probably the result of error or temporary trap fluctuation.

\[ \alpha = \left( \frac{\omega}{\omega_{xz}} \right)^2 \]

Figure 5.12: A measurement of the radial common mode frequencies for a two ion system in a near harmonic potential as a function of the strength of that potential. Raman motional spectroscopy was used to find the peak values.

These measurements allow models to be fit for which we can map the SimpleSol parameter directly to the \( \alpha = \left( \frac{\omega}{\omega_{xz}} \right)^2 \) values. In older work on this scaling the radial common mode frequencies and the axial mode were measured for by the tickle method\[ESG^+00]\). In this method the axial and radial modes are found by coupling in a small extra signal near the RF frequency \( \nu_{\text{drive}} = \Omega_{RF} + \nu_{\text{tickle}} \). When this frequency comes to within a few kHz range of the radial modes the chain will melt but remain trapped and the frequency can be noted. DC filters prevented these frequencies from coupling into our DC electrodes which would allow the axial melting frequencies to be easily observed. In reality there is probably some difference in the radial and axial common modes in N-ion chains measured via tickle and the result obtained by measuring the confinement strengths using this two ion tool. Because of
Figure 5.13: A measurement of the radial tilt modes frequencies for a two ion system in a near harmonic potential as a function of the strength of that potential. Raman motional spectroscopy was used to find the peak values.
	his the exact values for the power law scaling parameters may be in error.

We have defined the phase transition point via first visible sign of radial dislocation, and an example of this can be seen in Figure 5.15. The procedure is to trap an N-ion chain at a particular RF strength. We then scan the SimpleSol parameter which we have previously calibrated to translate into an approximate $\alpha$ value. Once we see the the visible shift into a small radial dislocation, we check that this transition is repeatable by toggling the smallest increment of SimpleSol (at the $\frac{1}{100}$th level). We do this at larger and larger ion numbers until the process begins to break down.

One symptom of this is that the dislocations become asymmetric, e.g. some portion of the chain will buckle off-center first, and the actual axial confinement needs to be very weak for the chain to be linear once loaded. For a higher radial potential depth, the axial confinement can be larger while still being able to support a linear chain without buckling. However for low RF the axial confinement needs to be very weak, and then the chain becomes highly effected by axial surface noise which will distort
Figure 5.14: A measurement of the axial common mode frequency as calculated from the radial common and tilt modes in a two ion system. There is one outlier point which may be an error or trap fluctuation.

its potential and shift its position. For this reason macroscopic traps with deep radial potentials can support longer linear chains without buckling.

Here we present two examples of our scaling measurement which demonstrate this issue, in one case there is just the SimpleSol potential alone and in the second case an underlying flat axial potential which is held constant added to the regular near-harmonic SimpleSol potential. Both of these $\alpha$(SimpleSol) potential functions are characterized independently via the previously described two ion method. Usually only the results with $\alpha$ calculated using the weaker (lower frequency) radial mode are presented but in our case we show also the results for the higher frequency mode as well to demonstrate the different behavior.
Figure 5.15: (a) An image of a 12 ion Zig-zag chain in the linear crystal regime. (b) Right as the Zig-zag threshold is crossed the central ions begin to dislocate into a subtle zig-zag form. The Zig-zag shape will become more exaggerated until the threshold for a helical crystal is crossed.

Near harmonic simple potential scaling

We can see the results for the scaling study in Figure 5.16 we perform linear fits to both log transformed axes in order to extract a power law scaling. The fits are performed only until the point where the model breaks down. For the nearly harmonic potential we see that the scaling begins to break down at the 7-8 ion chain regime for the low RF case but for higher RF case it is maintained up until 10 ions but rapidly diverges at the 11 ion level. We do observe a strange behavior. In principle the scaling law should be governed by the lower frequency mode. In this case the power law model remains valid longer for the higher frequency mode; however, for these higher frequency modes they have a large RF amplitude dependent shift.
**Figure 5.16:** We conduct a measurement of the zig-zag threshold $\alpha_{\text{crit}}$ as it scales with ion number. In this plot we display two voltage levels for a purely harmonic potential. The scaling relationship breaks down by the 10 ion level even for the 200 V case. We see that the scaling relationship is close to RF voltage independent only for the lateral (x-axis) of the trapping potential which will have a stronger confinement (as expected). We also see that higher trap RF voltages will lead to a well-behaved scaling relationship for longer ion chains.

**Combined flat and near harmonic potential scaling**

For the flat potential plus the near harmonic potential the measurements can be seen in Figure 5.17. In case of the flat potential plus the SimpleSol potential we can see that the scaling remains valid for longer ion chains at the 200 V level. At this level it diverges at 13 ions while for the 250 V level it still has not diverged fully at 15 ions. This means that by adding a flat component to our radial potential we can keep much longer chains at a fixed RF level without buckling. With this potential the scaling law also overlaps much better for the low frequency case, but the actual values of
the exponents differ from models measured in macroscopic traps [ESG+00, Sch93]. Confirming and studying this result in surface traps could be interesting, are the power-law scaling parameters truly different for surface Paul traps? One may need a surface trap apparatus which allows the measurement of the axial modes via the tickle method, or alternatively have axial coupling via Raman motional spectroscopy. Furthermore for surface traps, which have a sharp upper RF voltage limit due to trap breakdown, the need for flat axial potentials to scale to long ion chains is highlighted. These flat potentials which allow long chains to be formed without buckling are also more susceptible to axial motional excitation. The problem of preventing zig-zag buckling while managing axial noise in long ion chains on surface traps is of the utmost import as the attempt to scale past the 10-15 ion chain regime.

5.3.3 Vacuum Characterization with Zig-zag Chains

The core trap cryopackage used in our system is too compact to permit the integration of an ion gauge or cold-cathode gauge to monitor local pressure. This motivates the development of methods to use the behavior of ions to calculate pressure. One method is to utilize a double well quartic potential of known well height and then monitor the hopping rate of the ion [AVN+20]. A double well potential measurement was problematic in our system due the low potential barriers necessary to probe the low energy collisions at cryogenic temperatures. The noise on the digital-to-analog converters (DACs) was comparable to this well height and it would create a drifting well balance. In order to estimate the vacuum level, we employed a different method utilizing the zig-zag shaped chain of ions [PHK+18].

A long linear chain of ions can be trapped in a surface trap and the axial confinement potential can be ramped up until a threshold is exceeded. Upon crossing

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Figure 5.17: We conduct a measurement of the zig-zag threshold $\alpha_{\text{crit}}$ as it scales with ion number. In this plot we display two voltage levels for a harmonic potential with an added flat potential component. As compared to Figure 5.16 we see that for the same 200 V RF level the scaling relationship is maintained to 12 ions. We see that the scaling relationship is close to RF voltage independent only for the lateral (x-axis) of the trapping potential which will have a stronger confinement (as expected). In this example we tested a higher RF voltage of 250 V to see that with the flat potential addition our scaling relationship is maintained until the 15 ion level. This is a realistic operating condition for our trap to operate with 15 ions without radial dislocations, allowing for good radial two-qubit gates.

At this threshold the chain buckles and collapses into a zig-zag shape (zig-zag phase transition) [Sch93]. Due to the slight asymmetry in the two radial modes of the trap, there are two energetically degenerate stable states that can arise in a zig-zag shape. The ion positions are the mirror images of each other about the x-axis (trap axis).

We capture images of 7 ion zig-zag chains and count the rate $\gamma_{zz}$ at which they flip between the two stable configurations. $\gamma_{zz}$ is related to the molecular collision rate $\gamma$ by $\gamma_{zz} = p_{\text{flip}} \gamma$, where $p_{\text{flip}}$ is the probability of a zig-zag flip event given a collision.
with the background molecules.

If \( p_{\text{flip}} \) can be estimated as a function of temperature, then this information can be used to translate the observed flipping events between the two zig-zag configurations into molecular collision events, and from these collision events a Langevin model can be used to estimate the pressure inside the cryopackage assembly. The energy barrier between the two zig-zag configurations determines the behavior of \( p_{\text{flip}}(T) \). The small energy barrier between the two degenerate modes in the potential can be estimated from geometric considerations.

We estimated the energy barrier between the two zig-zag configurations by considering the energy it would take to rotate a zig-zag chain by 180° about the x-axis between the two mode shapes, as this represents a path that the chain could take from one state to the other. We estimated this energy barrier by first calculating the zig-zag structure for seven ions in a 3D harmonic potential. The strength of the confining potential in each axis is measured via RF tickle spectroscopy [NVG+13] of the common modes of motion. This measurement is done on a zig-zag chain under the same trapping potential used for the experiments. A simplex optimization algorithm is then used to relax ions into a zig-zag shape in the potential. Once the lowest energy state is found we rotate the chain about the x-axis by 90° which will be its highest energy rotational position about this axis. We use our potential energy calculation tool to then calculate the energy difference between the 0° and 90° position to estimate the ‘barrier height’ between the two modes. In the type of zig-zag potential we used for this experiment, we calculated this energy difference to be approximately 0.02 meV.

The average kinetic energy imparted on a trapped Yb\(^+\) ion due to collision with background H\(_2\) molecules is estimated in Eq. (5) of Ref. [AVN+20]. For every collision with a background H\(_2\) molecule in the cryogenic environment, the ion receives 0.023
times the average kinetic energy of the molecule, as most collisions are glancing and do not transfer all of their energy to the ion chain. This is factored into the model as a larger effective threshold kinetic energy for the zig-zag transition. Using this scaled effective threshold energy of $E_{\text{th,eff}} = 0.876$ meV, we estimate the population of background molecules that could lead to the transition between the two zig-zag configurations as a function of temperature. Here the Maxwell-Boltzmann speed distribution is used to calculate the fraction of gas molecules with kinetic energy large enough to cause a zig-zag transition. Ultimately this analysis leads to a scaling behavior given by

$$p_{\text{flip}}(T, E_{\text{th}}) = \frac{\int_{E_{\text{th,eff}}}^{\infty} \left( \frac{m_H}{2 \pi k_B T} \right)^{3/2} 4\pi \left( \frac{2E_k}{m_H} \right) \exp\left( -\frac{E_k}{k_B T} \right) dE_k}{\int_0^{\infty} \left( \frac{m_H}{2 \pi k_B T} \right)^{3/2} 4\pi \left( \frac{2E_k}{m_H} \right) \exp\left( -\frac{E_k}{k_B T} \right) dE_k} \int_{E_{\text{th,eff}}}^{\infty} E_k \exp\left( -\frac{E_k}{k_B T} \right) dE_k, \quad (5.6)$$

where $E_k$ is the kinetic energy of the background molecule. Using this probability of transition between the zig-zag configurations given a collision, we can invert the zig-zag transition data to estimate the collision rate $\gamma$ as a function of temperature. Finally, we substitute $\gamma = \gamma_{zz}/p_{\text{flip}}$ and use a classical Langevin model to calculate the pressure $P$ as

$$P = \frac{\gamma_{zz}(T)k_B T}{p_{\text{flip}}(T, E_{\text{th}})e} \sqrt{\frac{\mu e_0}{\alpha_{H_2} \pi}}, \quad (5.7)$$

where $\alpha_{H_2}$ is the static polarizability of H$_2$ molecules [GYCD96], $e$ is the electron charge, $\mu$ is the reduced mass of the H$_2$ and $^{171}$Yb$^+$ system, $\epsilon_0$ is the permittivity of free space. Fig. 5.18 shows the results of this analysis. On the left-hand side we show the transition rate between the two zig-zag configurations observed (events per minute). On the right-hand side we plot this transition rate re-scaled by Equation 5.7 to estimate the pressure inside the UHV assembly. We see that the pressure
scales with temperature exponentially. At cryo-UHV conditions, the error of pressure estimation is large because the chain transition rate is extremely low (a few events over several hours). Elevating the system temperature and observing the increasing prevalence of zig-zag flipping events allows one to trace out an estimated pressure curve. This pressure curve is most useful to characterize the onset of the regime in which transitions between the zig-zag configurations are prevalent. This exponential scaling with temperature arises because of the exponentially activated desorption of \( \text{H}_2 \) molecules from the activated carbon getter. Our error analysis assumes that the worst case miscounting of flipping events is modelled by the shot noise error \( \pm \sqrt{\gamma_{zz}} \) of the observed flipping rate.

![Figure 5.18](image-url)

**Figure 5.18:** On the left axis we see the transition rate between the two zig-zag configurations measured in the experiment as a function of the sample temperature. On the right axis we see the data re-scaled to pressure using Eq. (5.7). The error bars represent the pressure calculated using an interval of \( \pm \sqrt{\gamma_{zz}} \) of the observed value. The largest possible error in counting the zig-zag flipping events is modelled as the shot noise error. Appearing first in [SIJ+21] © 2021 IEEE
5.3.4 Dark Lifetime and Chain Reordering

The dark lifetime is defined at the average lifetime that a single ion will stay trapped without Doppler cooling. Collisions with background gas can excite the motion of ions, and if this excitation exceeds the trap depth then the ion can be lost. We spent time measuring this thoroughly on the HOA2.1a. We leave the ion in the dark and check after a time. This was repeated 6 times for each time interval and a survival probability was calculated. The results can be seen in Figure 5.19. The time investment was never made to measure it on the HOA2.0 system.

The cryopackage lid was reused for both of these HOA packages, and for HOA2.1a construction we did elevate the temperature of the trap cryopackage to increase the indium bonding strength. We had never rebaked the carbon getter package between these traps, and effort was made to simply minimize air exposure when swapping it to the new trap. Better performance would be seen by making a new cryopackage lid entirely for the Peregrine.

For the Peregrine series of traps we designed a new lid with a fresh carbon getter. After baking this getter initially and installing it into the cryopackage lid, the temperature of the system was never raised above room temp again. We have observed immeasurably long dark lifetimes in our Peregrine trap. We keep the same ion in the dark every night for weeks at a time without loss. This is an amazing performance for a surface trap which at room temperature may have a dark lifetime closer to 30 minutes in a well prepared chamber UHV system [MBB+13]. The Peregrine cryopackage has been temperature cycled many times (and reused) and most recently a single trap cryopackage has been continuously in operation for 6 months while retaining dark ions for entire 72 hour periods. Two ion chains can survive in the dark without loss for up to 20 minutes, and longer chains for a few minutes.
Figure 5.19: A dark lifetime measurement was performed for the HOA2.1a trap cryopackage. The points are the survival probability after leaving a small population of ions in the dark for a specified amount of time and then checking them. For the Peregrine trap this could not be measured because the dark lifetime for a single ion can exceed 72 hours. Every morning when we turn on our lasers that ion remains trapped from the previous night.

Chain reordering can be observed if the system temperature is too high. In a well prepared trap with an optimized base temperature of 9K or below another isotope can be trapped in a linear chain and reordering events will not be observed. The threshold for reordering is larger than the threshold for zig-zag flipping events while the collisions are low energy compared to room temperature.

5.4 Long-Lived Dark Events

Two problematic long lived dark events have been observed in this trapping system. The first is unrecoverable and was disruptive to working with long ion chains. The mechanism is not yet fully known, but it bears mentioning because we have found out how to stop it. The second is irritating and can lead to 15-20 min periods of
downtime, but its incidence is fairly rare. We explore these two events: doubly ionized ytterbium and YbH⁺ in this section.

5.4.1 Observation of Doubly Ionized Ytterbium Production

We observed the production of the second ionization state of \(^{171}\text{Yb}^+\) to make \(^{171}\text{Yb}^{2+}\). The symptom of this event compared to the most usual event causing an ion in a chain to go dark is that the dark event will also change the spacing of the ions simultaneously. Furthermore an ion once doubly charged will never return to usability, and these events can be often enough that they will eventually consume the entire chain of ions. In this system we used a wavemeter locked 638 nm laser, which we shined into the chamber with 2-3 mW power. Because of a large beam size it also would scatter a lot off of the trap. We found that these double charging events have two properties, one is that the rate of double charging is proportional to the RF amplitude as seen in Figure 5.20 and the second is that without the 638 nm light these events do not occur. Furthermore their rate also decreases when 638 nm is unlocked such that it is not resonant with the transition. We were able to remove these events by ceasing to use the laser for our F-state repump, instead we used 355 nm light.

5.4.2 Possible Observation of YbH⁺ Production

One final occurrence which has been noted is a long-lived dark event which does not change the spacing, and can be removed by application of intense 355 nm light. Usually when we see an F-state dark event we immediately shine our 355 nm global beam and the ion is removed from the F-state within a few seconds. However sometimes, especially if we are slow to begin shining 355 nm, a light a long lived dark state is seen which can take 15-20 minutes at least to become bright again with constant
Figure 5.20: The rate of doubly charged Ytterbium production is observed to depend upon the RF amplitude. Double charging goes away without 638 nm light applied. This data was collected by Ismail Volkan Inlek.

![Graph showing the rate of double charging vs. trap voltage amplitude]

Application of the 355 nm global beam even at full power (80 mW focused to a 15 μm beam waist). These events are very clearly different than the regular F-state events with a large difference in the rate at which a 355 nm beam will bring the ion back. If one waits long enough then they will always come back. We hypothesize that these events are caused by hydrogen molecules colliding with already dark F-state ions. The existence of this mechanism has already been established [HJOS20, SY95]. Furthermore these events only begin to become a prominent occurrence once the system has been operated with full RF for a long period of time after first cooling down (after a few months of daily operation). Even then these events happen only a few times a day, but over time can become more prominent. If the system RF is turned off completely and allowed to sit at base temperature for a 24-48 hour period, the incidence of these events will decrease again for a few weeks. This evidence is admitted anecdotal yet is repeatable enough that it has become the response to a large amount of these events which does work. Achieving even lower base temperatures
with RF at full amplitude may remove this issue. More systematic investigations of this phenomenon are needed.
Chapter 6

Qubit Gate Experiments

In this chapter we explore the performance of quantum gates in our system. For this chapter we have only focused on more recent work completed with the Peregrine series traps. While the single qubit gate performance was similar for the HOA 2.0 series of traps the long-term motional mode stability has been better on Peregrine. We have been able to do more fruitful work with two qubit gates on the Peregrine trap. We first will characterize the microwave and counterpropagating single qubit coherence times using the Ramsey method. We will then present the results of single qubit gate set tomography (GST) used to characterize the state preparation and measurement errors (SPAM) and counterpropagating gate fidelity, while also determining the prevalence of systematic errors. We present a case study for extremely low frequency (ELF) noise which presents itself only on the motional mode. The International Telecommunications Union defines ELF frequencies as being between 3 Hz and 30 Hz, but here we use the term to include also frequencies below 3 Hz, sometimes this is referred to as tremendously low frequency (TLF). We explore the unique signatures of ELF coherent noise and a method of characterizing such noise. We speculate on mechanisms which can generate such a noise and how the noise source was removed.

Lastly we present our work on FM two qubit gates. We provide a brief introduction to FM gates, explore the unique motional mode noise in our system which limits our gates. We present a gate survey featuring FM gates with diverse parameters (such as gate time, predicted displacement error, and Rabi frequency). Our gate
times range from 120 $\mu$s to 450 $\mu$s and we sample 32 robust-FM gates [LLF+18] and 5 $f$-robust-FM gates [KWF+22]. We discuss the interpretation of this gate survey, which is dominated by parity errors, key in which is a motional dephasing process which is dominated by low frequency (30 Hz - 3 kHz) coherent noise which could come from spurious currents flowing through the shared system grounds. We examine the spectrum of ground noise present in the setup and use it as a hypothesized power spectrum of frequency noise. This PSD of frequency noise is scaled by simulating the motional Ramsey experiment and matching our observed data. We utilize this scaled noise PSD to explain a large fraction of our gate error. By using this noise PSD in the filter function formalism we find that a large fraction of the gate error is predicted by this noise PSD alone. We present a master equation simulation of the entire gate survey completed using the Duke computational cluster. We apply our fitted noise PSD and generate phase noise according to this spectrum with a quantum Monte Carlo phase randomization we repeat our gate survey in simulation. We find that additional errors necessary to fit the data can be explained by a DC stochastic offset in the motional frequencies and the Rabi frequency. We also find that the white noise motional decoherence time necessary to explain our odd population infidelity is much longer than would be expected by just using the decoherence threshold time, our motional decoherence process is unique to this noise source.

6.1 Single Qubit Operation Characterization

Using microwave Ramsey, counterpropagating Raman Ramsey, and single-qubit GST (form counterpropagating Raman Ramsey) we characterize the single qubit coherence, optical coherence, gate fidelity, and SPAM errors.
6.1.1 Microwave Ramsey Measurement

We characterize the coherence for single qubit microwave manipulations. Here we directly drive the 12.6 GHz qubit resonance by applying radiation with a microwave horn external to the chamber. In this experiment we apply a microwave $\pi/2$ pulse to prepare a $(|0\rangle + |1\rangle)/\sqrt{2}$ state. We allow this state to evolve for a wait time before applying another $\pi/2$ pulse. We scan the phase of the second pulse to trace out a parity curve. We repeat the experiment for 11 phases between 0 and $2\pi$. We fit the curve to obtain a contrast for a particular wait time. Plotted in Figure 6.1 is the result for wait times up to 1 s. We observe almost exactly a 1 s coherence time for microwave operations (without spin-echo) typical for the ytterbium system whose qubit states are first order magnetic field insensitive.

![Microwave Ramsey Contrast](image.png)

**Figure 6.1:** The microwave Ramsey experiment to measure the qubit coherence time. We measure an approximately 1 s coherence time for microwave rotations.

6.1.2 Counterpropagating Motional Raman Ramsey Measurement

We then use our counterpropagating Raman beams to perform the same experiment. This measurement will add to the qubit decoherence the effects of optical decoherence
from phase instability of the Raman lasers. Some decay from heating will also be included. By comparing the counterpropagating Raman coherence time to the microwave qubit coherence time, we can get a sense for the effects of beam instability of our optics. In this system we achieve a 527 ms coherence time which can be viewed in Figure 6.2. Without any triggering to the cryocooler’s pumping this is an impressive testament to the interferometric stability of our beams. The counterpropagating coherence time is comparable to our qubit coherence. This result is impressive because we use a tightly focused beam and we do no triggering to the pumping frequency. Many cryogenic setups struggle with the cryocooler’s mechanical noise hindering their optical coherence.

![Graph](image)

**Figure 6.2:** The counterpropagating carrier coherence is measured by the Ramsey experiment (without spin-echo) on our single qubit gate operations. We find a coherence time of 527 ms measured on the Peregrine trap.
6.1.3 Single Qubit Raman Gate-set Tomography Results

In order to more rigorously calculate the fidelity of single qubit operations, under a typical system calibration, we performed gate-set tomography (GST) [NBKGR16, NGR+21, Gre15]. The method of GST defines a set of ‘germs’ which are preset sequences of quantum gates chosen to amplify some deviation from the target gate [NGR+21]. Long sequences of these germs can be concatenated to enhance these deviations. We utilize pyGSTi (a python implementation of gate set tomography created at Sandia) for generating our gate sequences and analyzing the resulting data. For our purposes because there is no natural $Z$ gate (but a virtual one) in our system we restrict ourselves, to testing the identity gate $I$, and the Pauli $X$ and $Y$ gates. We utilize the ‘SK1’ sequence with a Ramsey calibrated qubit carrier frequency to implement these GST sequences. The SK1 pulse sequence is is part of a family of Solovay–Kitaev sequences which is robust to variations in pulse length, ion Rabi frequency, and single qubit crosstalk errors [BHC04, MB14]. Our Raman intensity was locked (discussed in the two-qubit gate section), however the intensity lock duty cycle was not integrated into our control sequence optimally. We have since upgraded our control software and improved the Raman intensity lock. In Figure 6.3 we show the results of two similar GST executions which have been collected continuously in one shot after one initial calibration sequence. The difference between them is that in 6.3(a) we implement the $I$ gate as an instantaneous pause whereas in 6.3(b) we implement the $I$ gate as a full pause for one $\pi/2$ time with no wait.
Figure 6.3: We compare two model violation visualizations for our single qubit counterpropagating Raman GST. (a) In this GST sequence we implement the $I$ gate as an instantaneous pause (b) in this second sequence we model the $I$ gate as a pause for a full $\pi/2$ time making it the same duration as our gates. The bar graph represents the model violation (measured in standard deviations) per iteration for sequences up to sequence length $L$. The second plot (a histogram) shows the distribution of model violations for each circuit which should conform to a $\chi^2$ distribution, violations of the model at the 0.95% significance level would appear in red.

We see that in Figure 6.3 we have two sorts of plots which compare the observed data to the behavior under a Markovian error model. The first plot shows the model violation in units of standard deviations of actual log-likelihood which exceeds the expected value from the Markovian model. This is shown up to a subset of gate sequences up to length $L$. The bars are colored by their ‘star’ rating by pyGSTi and here we see only dark green (5-star) and light green (4-star). The plot on the left shows a histogram of model violation on a per-circuit basis, which should obey a $\chi^2$ distribution. Any violation of this distribution to the 95% confidence level would be
shown as a red histogram bar, but no such bar exists.

Figure 6.4: A table of model violations per germ, significant model violations would appear as red squares in the heatmap. We see a roughly even distribution of grey squares which get darkest at the 128 sequence length.

In Figure 6.4 we see a heat-map which will display the same model violation metric but with the detail of which germs are being tested. Any clear model violations will be shown as red boxes. We can see that for these two counterpropagating Raman gate GST executions we have a fairly regular distribution of light and dark grey, and a complete absence of red model violations. This indicates good performance.

Some metrics of interest to compare are the infidelity, the trace distance, and the diamond norm. These are all different metrics for comparing the distance between the ideal quantum state and the observed state [GLN05]. The pyGSTi package will
prepare a calculated infidelity for our gates $X,Y,$ and $I$ by these three metrics. Some discrepancy between the infidelity and the trace/diamond norm can indicate there are still some systematic errors present in the system. In Tables 6.1 and 6.2 we see the results of these two GST runs for datasets from Figures 6.3(a) and (b) respectively.

**Table 6.1:** For the GST run of Figure 6.3(a) we present the calculated distance between the desired gates $X,Y,$ and $I,$ and the observed gates. We show the infidelity, $1/2$ trace distance, and the $1/2$ diamond-norm distance. In this GST run we used an instantaneous pause to implement the $I$ gate.

<table>
<thead>
<tr>
<th>Gate</th>
<th>Entanglement Infidelity</th>
<th>$1/2$ Trace Dist.</th>
<th>$1/2$ Diamond Dist.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$I$</td>
<td>0.000053</td>
<td>0.00074</td>
<td>0.000757</td>
</tr>
<tr>
<td>$X$</td>
<td>0.001305</td>
<td>0.012622</td>
<td>0.01645</td>
</tr>
<tr>
<td>$Y$</td>
<td>0.00164</td>
<td>0.012259</td>
<td>0.015844</td>
</tr>
</tbody>
</table>

**Table 6.2:** For the GST run of Figure 6.3(b) we present the calculated distance between the desired gates $X,Y,$ and $I,$ and the observed gates. We show the infidelity, $1/2$ trace distance, and the $1/2$ diamond-norm distance. In this GST run we used a $\pi/2$-time wait time to implement the $I$ gate.

<table>
<thead>
<tr>
<th>Gate</th>
<th>Entanglement Infidelity</th>
<th>$1/2$ Trace Dist.</th>
<th>$1/2$ Diamond Dist.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$I$</td>
<td>0.00006</td>
<td>0.001456</td>
<td>0.00152</td>
</tr>
<tr>
<td>$X$</td>
<td>0.001062</td>
<td>0.012284</td>
<td>0.016517</td>
</tr>
<tr>
<td>$Y$</td>
<td>0.001233</td>
<td>0.01196</td>
<td>0.016884</td>
</tr>
</tbody>
</table>

In Table 6.3 we see the SPAM errors calculated by the GST method. We observe a $0.55\%$ SPAM error for a standard calibration procedure. We currently lack a precise *in-situ* $X/Y$ degree of freedom for lens alignment. This limits our ability to perform precise diffraction limited lens alignment in practice. We also do not have a robust passive alignment scheme in place for the $X$ and $Y$ degrees of freedom. This currently limits the minimum attainable SPAM error. A slight redesign of the imaging lens stage could help improve this figure. Another limiting factor for reducing SPAM error
is polarization drift on our state detection and pumping beams.

**Table 6.3:** The average of state preparation and measurement (SPAM) errors through the course of our single qubit GST. We compare the target states of 1 and 0 to the prepared states and we find a 0.55% SPAM error. Improvements in SPAM error could come from improving our lens alignment procedure to increase the detection counts.

<table>
<thead>
<tr>
<th>Target $\rho_0$</th>
<th>1</th>
<th>0</th>
</tr>
</thead>
<tbody>
<tr>
<td>Estimated $\rho_0$</td>
<td>0.994478</td>
<td>0.005522</td>
</tr>
</tbody>
</table>

Lastly we show the results of 6 different GST runs where we generate a GST sequence based upon only the $X$ and $Y$ gate. This type of GST is less time consuming and only will take approximately 30 minutes from start to finish and so we use it to further characterize our $X$ and $Y$ rotations without the influence of idling errors. The short duration also means that it can be more easily run repeatedly. We complete a standard calibration of both the state detection and the $\pi$-time for gates and then GST up to sequence lengths of 128. The results of all 6 of these runs averaged together is shown below in Table 6.4 where the standard deviations of the measurements are shown beside as errors. These results show that we are performing single qubit gates with approximately a 99.9% fidelity over many circuit-wise repetitions. Our SK1 pulse sequence pulse calibration could be improved and our optical intensity lock would later be improved but GST has not been attempted again.

**Table 6.4:** The mean gate quality metrics of six GST runs where only $X$ and $Y$ gates are implemented.

<table>
<thead>
<tr>
<th>Gate</th>
<th>Entanglement Infidelity</th>
<th>$1/2$ Trace Dist.</th>
<th>$1/2$ Diamond Dist.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$G_x$</td>
<td>0.000632 ± 0.000294</td>
<td>0.0117 ± 0.00290</td>
<td>0.0156 ± 0.003372675</td>
</tr>
<tr>
<td>$G_y$</td>
<td>0.00103 ± 0.000237</td>
<td>0.0117 ± 0.00204</td>
<td>0.0158 ± 0.00226</td>
</tr>
</tbody>
</table>

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6.2 Extremely Low Frequency Motional Mode Noise

Here we discuss the manifestation of an extremely low frequency coherent noise source appearing only on our motional mode. The signatures of this type of noise source can be confusing to interpret. Such a noise source was first noticed as a motional coherence scan with a strong oscillating envelope. An example of this can be seen in Figure 6.6(a). One can notice that the first collapse occurs at roughly 0.3 ms and the pattern continues from there. This is deceptive. The nature of this motional mode noise was made apparent only when we decreased the optical power significantly below where we would typically operate for gates. Thus our motional mode $\pi$-time was on the order of 2 ms. Decreasing the optical intensity at the ion will narrow the motional mode peaks and then we can resolve the issue. In Figure 6.5(a) we see a fine frequency scan of the motional mode. We notice that for a single motional mode we resolve two prominent peaks, displaced by approximately 700 Hz. Over the course of the frequency scan, we see a series of regular sharp dips in prominence.
Figure 6.5: We interrogate the motional mode with a frequency scan at the $\pi$-time (no sideband cooling) under low optical intensity to narrow the peak such that we can observe two distinct motional mode locations about 700 Hz apart. (a) We perform this experiment at the natural 4 Hz rep rate per experimental point. The noise is slower than the repetition rate of datapoint generation. Sometimes the motional mode is at the opposite location and we have low brightness. At some points the mode moves back in resonance with the scan and we see a peak. Because of the interplay between the datapoint generation repetition rate and the slow coherent noise, with a fine scan, we see a sharp series of peaks. (b) We introduce a waiting time bringing our experimental repetition rate below the noise frequency. The noise fluctuations will be averaged away. We still see what appears to be a bifurcated motional mode peak but now the sharp oscillating behavior is smoothed by under-sampling in time.

These dips and peaks are a results of the interaction of a strong low frequency noise with an experimental repetition rate which exceeds this noise frequency. Each point on the scan plot is the result of 100 single shot experimental repetitions. In
order for there to be such a sharp rise and fall of the data points, we require that the mode frequency has moved, then stayed stationary for the full duration of the 100 experiments. In this case our total repetition rate per data point is about 4 Hz which includes all single shot experimental timescales repeated 100 times. We generate a data point at a rate of 4 Hz, and for these purposes, we call this the experimental repetition rate. We know that the motional mode frequency has bifurcated and is moving sharply between two states at a rate lower than this experimental repetition rate. In 6.5(b) we confirm this suspicion by artificially lowering the experimental repetition rate by including a 15 ms wait time per experiment. This lowers the total repetition rate for the generation of datapoints to 0.57 Hz. The new repetition rate per point is low enough that the coherent noise is no longer visible point-to-point. We still see a bifurcated peak but we undersample the coherent noise and the sharp oscillations disappear. This noise is so slow that the experimentalist must consider the total repetition rate of the experimental points generated during a scan.

The way in which the frequency scan can depend upon the experimental repetition rate has allowed us to pinpoint that this noise is strong and results in two fixed motional mode frequencies 700 Hz apart. Even if the experimental repetition rate is 1.3 Hz, we can still see signatures of sharp oscillation but this plot is not shown. Thus motion between these two states then must happen at a rate between 0.5 Hz and 1.3 Hz. The mode must be fairly stable in between the movements. This type of motional mode noise makes for extremely confusing motional Ramsey experiments. In Figure 6.6 we see two examples demonstrating this by comparing Ramsey experiments with different Ramsey wait time intervals. The experimental points are generated from a phase scan to obtain the Ramsey contrast. These Ramsey experiments do not need to be done with low optical power to obtain this behavior, and in fact we noticed 6.6(a) using Rabi frequencies that would be typical of our experiment. We see in
6.6(a) a repeatable oscillating envelope for the Ramsey experiment which we would observe without going to special efforts to finely scan the wait times. We see that the first collapse is seen at approximately 0.3 ms and then there is a periodic decaying pattern. In 6.6(b) we see that this slow decaying envelope is itself a product of undersampling; in this figure we do a finer scan at lower wait times. With a sampling rate of $t_s = 0.025$ ms, we can see repeatable sharp losses in contrast and then returns in contrast. When the sampling rate is brought to $t_s = 0.05$ ms we see that, by chance, we are on the top of this previous pattern and we miss the contrast loss regions. Finally when we return to the sampling rate of $t_s = 0.1$ ms we recover a slow oscillating envelope as in 6.6(a).
Figure 6.6: This type of coherent ELF noise on the motional mode produces confusing motional coherence plots. Each point on these plots is from a full phase scan with a contrast fitting at the particular wait time. (a) We see the Ramsey experiment conducted on the motional mode where the wait times are scanned in a fixed 0.1 ms sampling interval ($t_s$). We observe what appears to be a motional mode decay with an oscillating envelope. The signal collapses by around 0.3 ms then returns. (b) We find that if we vary the sampling interval, we will see different manifestations of this behavior. In this scan we initially have a fine $t_s = 0.025$ ms and the signal will collapse and return very sharply at different times. We find that by chance when we change to a $t_s = 0.05$ ms we are always on the crest of this fine oscillation and then when we return to the original $t_s = 0.1$ ms, we recover a similar oscillation period as (a) by undersampling. Depending upon the experimental rep rate of your phase scan experiment, which depends upon the Ramsey wait time sampling interval (from which the points are generated) we can recover very different unique behaviors in the final motional Ramsey experiment. This highlights the confusing nature of strong coherence ELF noise.
The particular behavior of this is determined by an interplay between the total time of the phase scan experiment as 11 points are generated (a rate of about 0.7 Hz) to the experimental repetition rate per data point (at about 7 Hz), each having 100 single shot experiments per point. As we change the wait time sampling rate ($t_s$) for the Ramsey experiment, this will change the increment by which these different sampling rates change as the wait times increase. For this example a $t_s = 0.1$ will change the repetition rate by close to 0.4-0.3 Hz out of 7 Hz, while for $t_s = 0.05$ the change will be by 0.3-0.2 Hz, and $t_s = 0.025$ will change it by close to 0.1 Hz. These changes in the experimental repetition rate can bring the repetition rate into resonance with a harmonic of the underlying 0.88 Hz coherent noise. The repetition rate of the phase scan itself is close to this ELF noise frequency. Situations can arise where the experimental repetition rate has a consistent relationship with the noise frequency. Even though the phase at which it starts is random, the phase relationship should be maintained during the scan. The phase relationship can be maintained because the jitter caused by control computer is still small compared to a Hz level oscillation. We later found out by direct electrical measurement that the ELF behavior is closer to a square wave which makes the effects of a random starting phase even less apparent. Noise at this frequency scale can interact with a control system repetition rate in unique and surprising ways!
Figure 6.7: (a) We repeat a sideband Rabi rotation at the $\pi$-time with a frequency fixed at one of the two observed peaks in Figure 6.5. As the motional mode moves between its two states, we observe the peak move in and out of resonance with the addressing beam. We display this repeated experiment plotted against the real lab time-tag for each experiment. (b) A Fourier transform of the time trace of (a) reveals a sharp peak at 0.88 Hz; the small peak appears to be an artifact of undersampling because the location of this small peak in frequency space depends upon the particular experimental repetition rate while the 0.88 Hz peak remains stable across experimental settings and time.

We can use the ion to directly measure the noise frequency. The sharp oscillation we observe on the frequency scan of Figure 6.5(a) will also appear in a time-scan. In this frequency scan we are doing a $\pi$-time rotation (1.5-2 ms) at the interrogated frequency. If we instead pick either peak and repeat the scan again and again at the same frequency, we will see an analogous rise and fall behavior. We repeat the same scan at a fixed frequency and time of the motional sideband, time-tagging the results in real time. We can see the results of this measurement over 175 seconds in Figure 6.7(a). We then perform a Fourier transform of this data which is shown in Figure 6.7(b). We see a sharp peak which is at 0.88 Hz. There is always an additional small peak whose location in frequency space seems to change with the particular experimental parameters which influence the repetition rate ($\pi$-time chosen for the investigation influences this). Thus it may be an artifact of the sampling rate. We
have not studied this thoroughly and so it could be an additional weak component which changes in time.

In order to eliminate this undesired behavior in the ion, we look for its signature independently in our setup. We have a capacitive pick-off which can be used to monitor the RF voltage amplitude at the trap. It exists so that we can lock the RF amplitude. We have had a long design struggle to produce a low-noise rectifier readout system for our pick-off. The original design was too noisy to resolve this $2.8 \times 10^{-4}$ level (700 Hz / 2.5 MHz) mode frequency drift. Upon an upgrade to the readout circuit, we could now clearly resolve this previously hidden RF intensity drift which appeared as a nearly-square wave with an 0.88 Hz frequency. We used an oscilloscope to perform ELF spectrum analysis. An oscilloscope trace 40 seconds long was obtained with a 4 ms sample spacing, a 10 second portion of which is shown in Figure 6.8.

A Fourier transform of this rectified pick-off signal is shown in Figure 6.8 as well. We observe this 0.88 Hz and its higher harmonics on the oscilloscope trace spectrum. Also in Figure 6.8 we see in red Fourier transform of a 15 minute audio recording of the lab environment. We zoomed into the frequency region of interest. The responsivity of a microphone to these low frequencies is small and so this signal is far weaker than the audio frequencies recorded. We still see sharp peaks at the even higher harmonics of 0.88 Hz. This result will be returned to later as we speculate on the source of this noise.
Figure 6.8: We take the Fourier transform of an oscilloscope trace of our RF pick-off circuit independently confirming the presence of our 0.88 Hz noise. This was enabled by a new low noise RF rectifier circuit which had a large enough SNR to see this coherent noise. We also take the Fourier transform of a 10 minute audio recording of the lab environment and we see the even higher harmonics of 0.88 Hz picked up by this electrically isolated sound sensor.

6.2.1 The ELF Noise Source

Typically when we talk about 1 Hz level noise in a cryogenic system, the first instinct is to blame the GM cryocooler itself; this is an obvious first candidate. The mechanical pumping frequency of the cryostat is 1.42 Hz which does not divide evenly into 0.88 Hz and thus our noise is not a harmonic of this pump frequency. We have also earlier established that the RMS vibration level of our cryostat is at the 2 nm level between objects on the optical breadboard and the sample chamber mount.
Furthermore we would expect that this effect would be more apparent on our carrier coherence than our motional mode frequency if vibration was to blame.

The noise is affecting the RF amplitude in such a way that it is independently detectable at the trap pick-off. For electrical background noise, 0.88 Hz is a strange frequency. Noise from appliances and man made electromagnetic interference (EMI) tends to be at the 60 Hz level or above. A survey of electrical engineering literature will not reveal many man-made sources of coherent 1 Hz level EMI. Entire power systems on a multi-city scale can have fundamental oscillation modes at the < 1 Hz level [LJC+15, BBDBLS07]. However power systems are designed to discourage this sort of oscillation and the excitation of these modes is transient and temporary.

Noise in the ELF range tends to come from atmospheric sources such as the Schumann resonance, which are oscillations on the cavity formed between the earth’s surface and the ionosphere [Sen17, Nic97]. However the frequencies of this noise are still too high, and the intensity varies dramatically depending upon the weather and seasonality. By contrast our noise is regular in frequency and varies in amplitude only slightly. Additionally picking up any noise in this range by the unintentional antenna mechanism seems unlikely due to the tremendous wavelengths involved.

We have earlier noted that we can pick this noise’s even overtones up on an audio recorder. This audio recorder is small and battery powered and is disconnected from the power system. The audio recorder will however pick up mechanical noise which is allowed to interact with its capacitive microphone. The audio recorder can detect these tones whether it sits directly on top of the cryostat, or if it sits on the rack above the experiment. It can also pick up the same tones on another experimental rack on the opposite side of the lab which houses a UHV chamber based experiment. The source is likely to be some mechanical commonality between experiments (not the cryocooler). Some candidates are the oscillation mode of the building, the oscillation
mode of the optical table, or the oscillation of the large interconnected rack system spanning the entire lab.

We used our pick-off signal to inspect our experiment to determine by what path the noise couples into the RF amplitude. We found that the singular source of coupling is the cables used to drive our cryogenic stages. These cables are connected to the experiment each by a 5-pin header connector. The header inputs are located on an electronic feedthrough PCB which is between the Montana system vacuum chamber and the bottom of the sample chamber. This PCB also acts as sealing surface for two O-ring seals at the bottom of the sample chamber. These header cable outputs can be seen in Figure 4.4(b) plugged into this PCB. The cables themselves are heavy and thick and lead to D-shaped connectors at the opposite end. These cables terminate on the rack above at the stage controller. The cables create a loose mechanical connection between the optical table and the rack system which is mounted to the room ceiling. The 0.88 Hz noise is unaffected whether the stage controller is on, or off, or even unplugged entirely from the power system. Unplugging the cables entirely from the stage controller box will not effect the noise. If the then-unplugged cables are left to rest, and then touched with a human hand, the pick-off level can be effected by just this touch. The amplitude of the 0.88 Hz noise can be decreased by unplugging the header connectors at the cryostat side, one at a time. When none are connected, the noise is not detectable. The noise contribution from each cable is not equal; one is particularly bad, and it seems to depend upon the mechanical details of the cables route.

The experiment from Figure 6.7 is repeated at a later date. This time the experiment was repeated with the stage cables either plugged or unplugged. It is clear the noise coupling can be removed by unplugging these cables when the cryogenic stages are not in use.
Figure 6.9: We repeat the experiment of Figure 6.7 at a later date, twice in a row, with and without the cryogenic stage control cables plugged in. We see that removing these cables will remove the path by which the 0.88 Hz noise can couple.

One theory for the mechanism of noise coupling is that the header connector is a loose connection which is subject to mechanical jiggle. As the cables sway, the header pins are jiggled and they change the trap capacitance. This time dependent change in trap capacitance changes the matching condition of our helical resonator very slightly, modulating the trap RF amplitude. While no scheme to measure the mechanical spectrum of these cables was attempted one can observe a natural swinging frequency at the 1 Hz level when perturbed. As a mechanical resonator with a very low Q-factor the bandwidth of the mechanical resonance would be broad. Upon consistent forcing with a 0.88 Hz frequency it is plausible that they could act as conduits for the 0.88 Hz mechanical forcing. We have not yet found the source of this very consistent forcing frequency. In the final section we discuss more extensively the unique electrical situation in a cryocooler and the Montana cryostat which may cause this sensitivity.

In short the helical resonator ground is connected to the cold finger sample mount (at RF frequencies) most strongly through the chassis of the Sumitomo cryostat. This is because the helical resonator sits on the sample platform at 90 K which
is connected to the stage 1 cooler. The cold finger sample mount (to which the trap is mounted) is connected to the stage 2 cooler. The stage 2 cooler has limited thermal and electrical connections which are only made near to the Gifford-McMahon heatpump mechanism. There is a connection with a small DC resistance but also some capacitive coupling at RF between the coldest stage (stage 2 at 7 K) and the first stage (stage 1 at 30 K). Thus looking into the trap from the helical resonator’s perspective it sees this capacitance between the stage 2 and stage 1 coolers. If that capacitance is changing in time this will change the matching condition and it could modulate the RF amplitude. A schematic to help visualize the situation is depicted in Figure 6.10.
Figure 6.10: A Figure visualizing how electromechanical noise can couple into our system capacitively. The resistive connection at warmer temperature between stage 2 and stage 1 and the capacitive connection between these stages are both seen by the helical resonator / trap system. Changes in the system capacitance could effect the RF amplitude.

At a later date there was a short episode where our cryocooler begin to exhibit some mechanical interference during its pumping cycle. An audible ‘clunk’ could be heard for each cycle of the cryocooler. And a similar signal on our pick-off but at 1.42 Hz could be seen. Furthermore we could detect it by the same method of ELF motional mode noise analysis which was used in Figure 6.7 to detect the 0.88 Hz. In
this case these cables were already unplugged. By a similar mechanism the change in capacitance (seen by the helical resonator) due to the mechanical interference of the heat-pump mechanics could be to blame. This time the noise was transmitted at the cryostat pumping frequency. The amplitude of the effect was less but this is further evidence that our system is sensitive to electromechanical noise. This problem went away on its own as did the ‘clunk’ sound and has not since returned. More direct evidence that the mechanism is due to a changing capacitance would be desirable; experiments could be designed to test this theory.

### 6.3 Frequency Modulated Two-Qubit Gate Optimization and Low Frequency Coherent Noise

Before we move on to algorithms and long chain operations, we require high fidelity two-qubit gate operation. The main struggle for high fidelity gates in this system has been motional mode stability. Reducing long term drifts, which on an experimental timescale, are seen as DC drifts experiment-to-experiment is important for maintaining the calibration during operation. Coherent noise on the motional mode had unique effects in our system which will be explored in detail.

MS gates in general utilize two off-resonant motional sideband drives to create a spin dependant interaction. After an interaction time $\tau$, also called the gate time, the ion spins are left in an entangled state. The motional modes traverse phase space during the interaction. Ideally this is in a closed loop such that there is no remaining entanglement with the phononic states. A perfect MS gate is described by a unitary operator $\hat{U} = \exp \left( \sum_{j_1,j_2} \sum_k (\alpha_{kj} \hat{a}^\dagger_k - \alpha^*_{kj} \hat{a}_k) \hat{\sigma}^x_{j_1} + i \Theta \hat{\sigma}^x_{j_1} \hat{\sigma}^x_{j_2} \right)$ where the $\hat{a}^\dagger_k$ term is the creation operator and $\hat{a}_k$ is the annihilation operator for the $k^{th}$ motional mode and the $\alpha_{kj}$ terms are the displacement in phase space of these modes [WWD18]. The $\hat{\sigma}^x_{j}$ terms are the spin flip operators for each ion. During the interaction they will
undergo a $\Theta = \pi/4$ rotation with respect to the $XX$ axis. Deviations from the ideal MS interaction can be characterized by a displacement error $\mathcal{E}_\alpha = \sum_k (|\alpha_{kj1}|^2 + |\alpha_{kj2}|^2)$ and an angle error $\mathcal{E}_\Theta = (\Theta - \pi/4)^2$. The total gate error being $\mathcal{E} = \mathcal{E}_\alpha + \mathcal{E}_\Theta$. This is related to the gate fidelity $\mathcal{E} = 1 - F$ and so here we often call it the infidelity.

Drift in system parameters during the interaction causes imperfection in the two-qubit gate. A leading source of error is drift and fast noise in the motional mode frequencies [WCF+20]. Although drift in the intensity of the addressing beam, optical dephasing, imbalance of the intensity on either ion, imbalance of the tones on either motional sideband, and many other experimental nonidealities can all contribute to gate error. Some of these issues are in finding the correct system settings (calibration) and some of them are due to stochastic drift (noise). Some could also be due to a noise source with strong correlations in time (coherent noise).

6.3.1 Active Experimental Stabilization

In order to reduce the amplitude of our stochastic noise we require a series of stabilization locks. The first is to stabilize the comb-tooth distance in frequency space to be resonant with the qubit transition. We monitor the repetition rate by sampling light immediately out of the laser as seen in Figure 4.12 and discussed earlier. We used the observed repetition rate to feed forward a tone to our counterpropagating individual beams using a scheme similar to [MGV+16] to keep the frequency combs resonant around a chosen global beam driving AOM frequency. Instability in this lock will result in optical dephasing and reducing carrier coherence, and lead to carrier frequency drift. Good carrier coherence is one indication that this lock is functioning properly.

The second major locking system, which will effect our MS gate quality, is the
Raman beam intensity lock. This lock is sometimes referred to as a ‘noise eater’. For this system we sample light after the MEMs mirror optics but before the chamber for each individual beam, and we also sample light from the global beam fiber output after the polarization filter. This light is directed at a photodiode (ThorLabs SM05PD3A) whose output is amplified (ThorLabs AMP120 transimpedance amplifier) and sent to our proportional-integral-derivative (PID) lock box analog-to-digital converter (ADC) so that it can be used to stabilize the Raman intensity. We can stabilize this intensity to approximately the 1% level. Because there may be a wide variety of 355 nm doses and powers during calibration and gate routines, we must choose to lock it at a particular consistent power setting. We run the lock for 1 ms during Doppler cooling by turning the Raman beams on and also triggering the intensity lock to run for this duration. The optical power is adjusted by variable RF attenuators in line with the RF amplifier which drives the AOMs for the Raman beams. After running the lock for this Doppler cooling duration, we hold the variable attenuator constant until the next experiment. By only looking at Doppler cooling, we tweak the power up to a consistent level before any coherent Raman operations are performed. This strategy allows the lock to remain stable over the course of many different calibration procedures and circuits.

The third experimental parameter which we desire to stabilize is the motional mode frequencies. The motional mode frequencies drift primarily due to changes in RF amplitude which drives the trap. Thus if we detect the RF amplitude, we may be able to feed back on the inputted RF power to stabilize the motional frequencies. It is challenging to create an active RF amplitude lock such as this without injecting fast noise into the system and actually lowering motional coherence. Obtaining an RF amplitude detection circuit with a signal-to-noise ratio sufficient to stabilize the frequencies to the 10’s Hz level will be difficult. So in the data present here we have
not stabilized the RF amplitude actively. Instead we stabilize the temperature of the RF amplifier. Amplifier temperature drift is a major source of long term drift of the motional modes.

![Graph of Amplifier Temperature vs. Motional Frequency Change](image)

**Figure 6.11:** We scan the lock point of our amplifier temperature and do Raman motional spectroscopy to observe a frequency shift with amplifier temperature. We observe a 5.35 kHz/C slope in motional frequency. The standard deviation of the residuals about the linear fit is 300 Hz which corresponds to a 0.056° C stability, although some of this frequency drift is likely from other sources.

Our locking system consists of a machined aluminum box with an open top which houses the amplifier. The amplifier rests on top of a Sorbothane bed inside the box to thermally isolate it from the box itself, and vibrationally isolate it from the lab environment. The box is ringed with tapped holes which are used to mount a large water cooling plate as the lid. This water cooling plate is thermally isolated from the box and mounted with Delrin screws, again to aid in thermal isolation. Mounted on top of the amplifier is a small aluminum block in which we mount a screw on thermistor. Machined into this aluminum plate is an indentation for an oblong TEC to rest. There is a gap between the water cooling plate and the machined aluminum housing, and we pin a TEC in between the amplifier (with plate) and the water cooling plate. We then run a PID temperature lock using the signal from the thermistor to
adjust the TEC. The temperature is stable to better than 0.1° C once locked. We demonstrate both the sensitivity to temperature drifts and the capabilities of our temperature lock by moving the lock point and tracking the motional frequency drift. This can be seen in Figure 6.11. From this data we can see that the motional frequency responds to changes in amplifier temperature by 5.35 kHz/K. If we calculate the standard deviation of the residuals (distance between the linear model and the data), we find 300 Hz, which corresponds to a 0.056° C instability. This is likely an overestimate because not all of the mode frequency drift is caused by amplifier temperature instability. However this number is a reasonable upper bound for the effectiveness of the lock and the 0.1 C° digit on our thermistor readout does not change when locked.

With this amplifier temperature lock in place, we observe good long term stability of our trap mode frequencies. Each day the modes remain stable to within approximately the 500 Hz to 1 kHz level when returning to the lab. However we seek to get a sense for the magnitude of long term drifts between experiments. We setup a low power motional spectroscopy to narrow the transitions. We set this experiment up to run repeatedly over an hour and collect the fitted common mode frequency. We can see the results of this long term study in Figure 6.12. Two features are immediately noticed. There is a sharper monotonic mode instability at the beginning of this scan, followed by a flatter region of stochastic noise. The initial mode drift is likely caused by charging. The experiment was started from a long period of darkness. Then we apply UV light with a uniform duty cycle. The UV light is low intensity so it takes a while for the charging to saturate. This level of charging is less severe than in past traps where the effect has been greater than 1 kHz in magnitude at least. During real gate operations, the charging should be mostly saturated by the long period of sideband cooling preceding the coherent operations. This period of sideband cooling
remains constant between our gate calibrations and our gate execution which will reduce susceptibility to this drift source. Nevertheless it is expected that the charging will effect frequency stability somewhat in excess of natural stochastic drift. This is because throughout the calibrations sometimes there are variable wait times, and different duty cycles of operation which may cause this charging to randomly walk. This is difficult to characterize exactly but we can put an upper bound of the effect as 400 Hz based upon Figure 6.12.

![Figure 6.12: Mode Drift vs Time](image)

**Figure 6.12:** We study the long term drift in our motional mode by repeated mode spectroscopy at low interrogation intensity. The results of a fit to reach mode line over 55 minutes are seen here. We have a uniform dose and duty cycle of UV exposure so there is an initial slow saturation period for UV charging which should put the upper bound of charging-induced frequency instability at approximately 400 Hz.

After the charging saturates, we see a flatter region of stochastic and long term drift. We can calculate the frequency instability due to stochastic drift from this region by excluding the charging data and computing the standard deviation which we find to be about 120 Hz. This experiment was conducted with the amplifier temperature lock on and stabilized.
6.3.2 MS Gate Calibration and Characterization Procedures

In order to calibrate MS gates in our system, we utilize a similar routine as was employed in \cite{WCF20}, based upon insights from \cite{Hay12}, other previous literature, and colleagues. We have adapted this calibration routine to the specifics of our system and workflow needs.

**Ion Intensity Balance**

We must balance the beam powers on each of the two ions. At the start of the day upon powering on the laser, we first measure the power on each individual beam using a power meter to ensure that at maximum operating amplitude, it does not exceed the power limitation of our MEMs devices which is 10 mW. At this point we manually roughly balance the optical power exiting the fibers using a wave-plate in the Raman upstream 4.12 which divides the power between the two individual paths. Once we ensure the power maximum is safe and the beams are roughly balanced, we lock the intensities. We check the global beam power and also lock its intensity. After this has been verified, we trap two ions. Once we have the ions, we first finely align the beam to each ion by maximizing the Rabi frequency by beam steering alone.

Once each beam position has been maximized individually, we perform Rabi rotations on each ion independently. We pick the faster ion and then seek to equalize the Rabi frequency of the other ion to the first. We do this first by adjusting the lock point for the beam until we equalize the rotation times between them. We do this first at one $\pi$-time, then at $3\pi/2$-times, then at $5\pi/2$-times. Once this is done the two time traces are difficult to distinguish and we move on. This calibration is checked roughly hourly to ensure it has not drifted. This calibration is done manually. The remaining calibrations are mostly automated via control software, with a few human
choice points.

**Rough Frequency Calibration**

We perform scans of the needed motional mode frequencies. These scans are done with low optical power to narrow the lines. We scan all four red sideband motional modes to calibrate for sideband cooling. We scan the tilt mode of the blue sideband which will be used for the gates. This is used as a starting point for the more careful sideband cooled gate calibration of the blue sideband.

**Two-Tone Fine Frequency Calibration**

Here we perform the fine sideband cooled frequency calibration for our MS gates. These calibrations must be performed at the same optical intensity of the gate so the AC Stark shifts are the same [Hay12]. We pick the laser power setting at which we will complete the gate. We must use both a red and blue tone so that the AOM power will be the same. We then fix the tone (which we are not scanning) at some value 100 kHz away from the its motional mode. We scan the other tone in the neighborhood of the motional peak after sideband cooling. We want the interrogation time to be one $\pi$-time of the mode under the conditions of the gate.

Initially, for a particular power setting, we will do a rough frequency scan. This rough scan is followed by a time scan under these conditions (and we fix the fitted frequency). Next we complete another frequency scan in which we fix the interrogation time to the $\pi$-time. Then with the correct $\pi$-time, we scan the frequency again fitting to the final value. For recalibrations at similar power settings, the scan only need be run once then automatically fitted for frequency. We must obtain the frequencies of the two red sideband modes (common and tilt). We prepare the ion into a $|1\rangle$ state so that we see a full peak (or trough). These two red sideband mode frequencies are
needed as an input for the FM gate solution optimizer. We also scan the blue tilt mode, which along with the red tilt mode, we will utilize to calibrate the asymmetric Stark shift correction and the red and blue tone balance.

**Stark Shift Calibration**

After fitting the sideband cooled modes of the gates, we perform a calibration which provides a more fine correction to the frequencies for the asymmetric Stark shift. This routine is based on a simplified version of a scheme from [Jun17]. The asymmetric Stark shift is optical intensity dependent and will also drift in time due to the Raman laser’s repetition rate drift [Hay12]. We are fitting for a small correction (not the absolute Stark shift) to the previously fitted frequencies which will be added to both modes such that they move in the same direction. In this scan we enact the MS Hamiltonian with a long time and large detuning using the frequencies we have previously calibrated (red and blue tilt modes).

We first run this as a time scan where we find a detuning which provides a clean scan of the large detuning MS interaction. We scan past the point of entanglement generation until we swap the \(|00\rangle\) populations for \(|11\rangle\) or \(\Theta = \pi/2\). We can see an example of this timescan in Figure 6.13. If the initial guess of Stark shift is far off then the interaction will not appear. Once we know the rough \(\Theta = \pi/2\) time for a particular detuning, we can set it and scan the Stark shift until we see MS interaction turn on at the correct value. The larger the detuning we use, the more narrow the peak. Also the \(\Theta = \pi/2\) gate time will be larger. For our system we were using detunings between \(35 - 125kHz\) depending upon the optical intensity level that we sought to calibrate gates for. In the example given in Figure 6.13 we are using a 85 kHz detuning, and we end up deciding that the correct interaction time is 2.5 ms. We call this the ‘long time large detuning MS’ Stark shift calibration technique.
Figure 6.13: We see the long time large detuning native MS interaction evolution in time. This was done with an 85 kHz detuning. Obtained from this scan we would then use a 2.5 ms MS interaction time when we then scanned to fit the Stark shift frequency.

Once we have setup this calibration experiment for a particular optical power setting, we can utilize this as a quick calibration scan to correct for the temporal drift in the asymmetric Stark shift. An example of this scan can be seen in Figure 6.14 where we fix the parameters obtained in Figure 6.13 and then scan only this asymmetric Stark shift added MS frequencies. We can then fit for the asymmetric correction to the addressing mode in which the MS interaction is optimized. Here we fit a -182.33 Hz correction.
Figure 6.14: We scan the Stark shift frequency for a detuning of 85 kHz with a time of 2.5 ms. Upon scanning we obtain a -182.33 Hz Stark shift which will be added as an asymmetric (about the carrier) correction.

While this method does not give the absolute asymmetric Stark shift, it will give a more precise correction to our previously fitted frequencies. This correction should also depend upon the Rabi frequency. We see plotted the results of this Stark shift calibration for a population of gates (which are studied at the end of the chapter) which we surveyed over a range of optical intensities and gate times. We can see that the magnitude of this asymmetric Stark shift correction will increase linearly with Rabi frequency. The direction that the parameter will move (positive or negative) will depend upon how the frequencies in the control software are set up, the laser repetition rate [Hay12], and the details of the AOM setup driving the Raman beams. Here we plot the direction it moves in our control software, which is a actually towards the blue side. In our case a negative shift is actually a shift in the blue direction relative to the carrier. In Figure 6.15 we see for our experiments over a range of Rabi frequencies that there is a linear relationship with a slope of -12646 Hz/Mhz which gives the ratio of asymmetric Stark shift to Rabi frequency (here we
use real not angular frequencies). We calculate the standard deviation of residuals between the linear model and the experimental data to be 97 Hz. This could indicate a rough measure of the level of Stark shift drift we experience on hour long timescales.

![Figure 6.15](image)

**Figure 6.15**: A plot of the calibrated asymmetric Stark shift correction for each of the gates in our gate survey. Here we see a linear relationship with a slope of -12646 Hz/MHz. The standard deviation of the residuals (distance between linear model and experimental results) is 97 Hz, which could be a good indicator of the level of Stark shift drift. There are two outliers here which are not due to fitting error; the cause of these outliers is unknown.

**Tone Intensity Balance and Gate Time Calibration**

After we finish the Stark shift calibration, we can check again the sideband $\pi$-time for rotations under the MS Hamiltonian. We require a rough measure of this time in order to calibrate the frequencies. After all frequency calibrations are completed, we can complete another motional mode time scan using both MS tones. Like the frequency scan, we detune the other tone by 100 kHz, keeping the optical power delivered by the AOM similar but far detuned (of the mode not being interrogated). We drive the tone to be interrogated on resonance and rotate $2\pi$ to $5\pi/2$-times and note the $\pi$-time. We then scan the opposite (red/blue) tone and seek to make fine power adjustments to match the rotation times well at $3\pi/2$-times. Once the red and blue tones are balanced in their rotation times, we are ready to search for a gate.
solution.

**Gate Solution Search**

We have noted the blue sideband \( \pi \)-time and from this, we pick a gate time in its neighborhood. We utilize an FM gate optimizer for robust-FM and f-robust-FM gates. These use the optimization schemes of \([LLF+18, Leu20, KWF+22]\). Two free parameters exist: the gate time and the Rabi frequency limit. We tweak these two parameters and search for an FM gate solution with low predicted error which has a prescribed BSB \( \pi \)-time close to our established and calibrated experimental time. This may take a few tries but it is easy to converge on a good solution in less than 10 (usually only a few) iterations of the gate optimizer. We find that it is more straightforward to search for an FM solution which fits our sideband \( \pi \)-time rather than adjust our sideband \( \pi \)-time to fit the FM solution. In this way less time is spent re-calibrating and performing amplitude calibrations for the gate. This work could, in theory, be automated by a meta-optimizer which will perturb the input parameters to the FM gate optimizer (and could parallelize the optimization) to find the ideal gate solution for your experimental parameters.

**Gate Detuning Scan**

Once we have picked a good FM solution for our calibrations, we then run a scan of the FM gate detuning. We use our controller to play our FM solution with a scanned total detuning and observe the behavior of the populations. We need to pick the correct detuning point at which the ideal gate is performed. We pick detuning points where the \( p_{00} \) population crosses over the \( p_{11} \) population while the \( p_{01} + p_{01} \) populations are suppressed\(^1\). If the amplitude is calibrated properly, the crossover point for the

\[^1\text{See the footnote at the end of Section 3.3 for a reminder of the definitions of } p_{00}, p_{01}, p_{10}, p_{11} \]
even states should be aligned to the minimum of odd population. Because of the robustness of FM gates, we often see a wide region of low odd population. In Figure 6.16 we can see an example of the detuning scan at the one gate level (a) being a broad scan and (b) being a more fine subsequent scan. We chose our detuning point in the center of the region where both population conditions are met. These scans will become more confusing as more gates are concatenated; traditionally the angle between the two even population curves should become more steep as we ramp up gate number. In our case this is true; however, due to our unique coherent motional mode noise and drift, we also see regions of many cross-over points which change scan-to-scan. This can make simple algorithms based on linear fitting behave unstably and we must choose the detuning points by hand.

Figure 6.16: A plot of the detuning scans for 1 robust-FM gate with a 175 µs duration. (a) We first perform a broad detuning scan. (b) We can then do a fine detuning scan, which for one gate is fairly flat in this range.

Parity and Population Scan

Once the calibration phase is complete, we can collect data on the gate performance on both a per-gate basis and in a series of concatenated gates. In order to extract the true infidelity as a result of the gate repetition and separate it from SPAM errors,
we characterize the slope of repeated gates. For each gate number, we require a scan of
the parity contrast and the odd population. A good discussion on how the parity arises
from the calculation of gate fidelity ($F = \langle \psi | \hat{\rho} | \psi \rangle$) with a density matrix is given in section 4.2.1 of [Hay12]. For the parity scan, we perform a global $\pi/2$
rotation on the qubits with a scanned phase ($\phi$), thus enacting the rotation operator
$\hat{R}^{(1)}(\pi/2, \phi) \otimes \hat{R}^{(2)}(\pi/2, \phi)$.

As discussed in more detail in [Hay12] we can calculate the fidelity as $F = \frac{1}{2}(p_{01} + p_{10}) + \frac{1}{2}c$. Here we define the $c$ as the amplitude of the parity curve traced out by
scanning the phase of this interaction. We refer to it as the parity contrast (in [Hay12]
is defined as the total peak-to-peak contrast and thus it has a factor of $\frac{1}{4}$). In order
to find ($p_{01} + p_{10}$) we repeat the gate 30 times and measure the populations of these
odd states, averaging the result.

![Figure 6.17](image_url)

**Figure 6.17:** A comparison of typical parity contrast curves. We compare a 175 $\mu$s
gate at 1 and 17 repeated gates and a 300 $\mu$s at 1 and 13 repeated gates. The 175 $\mu$s
gate is sampled at 61 phase points while the 300 $\mu$s gate is sampled at 21 points.
In Figure 6.17 we see some typical examples of parity contrast curves in our surveyed gates. We show a 175 $\mu$s gate at 1 and 17 concatenated gates and a 300 $\mu$s at 1 and 13 concatenated gates. The 175 $\mu$s gate is sampled at 61 phase points while the 300 $\mu$s gate is sampled at 21 points. For our gate survey 9 of the gates utilized this 61 points sampling and the rest used 21 points. By using fewer points, we can be less subject to drift after calibration, but more points can give a better sense of the true parity noise. Any difference in these two approaches is washed out by the true source of parity error, and we opted not to use 11 points (as is often done) because it would not always fairly sample the crests.

As noted in [Hay12] we must make sure that there is a stable phase relationship between the carrier used for the parity scan and the blue and red sideband frequencies used for the gate. With our old DDS-based system used for applying RF for the gates, if a rounding error existed such that the DDS tones used for the gate do not average to precisely the tone used for the single qubit rotation (for the parity scan) then the parity curve will have an unstable phase (totally random scan-to-scan). If the points are randomized than the signal is washed out completely.

In 6.17 we see that the 175 $\mu$s and 300 $\mu$s gates are phase shifted with the longer gate time closer to ideal. This is an artifact of the capabilities of our RF system-on-a-chip (RFSoC) control system and the choice to use the Ramsey calibrated underlying qubit frequency (rather than the midpoints between the red and blue sidebands) as the frequency for single qubit operations. This is desirable during circuit operation because SK1 pulse sequences can be used, and the RFSoC allows the experimentalist to set the phase relationship before our gate (to this qubit frequency) such that is it consistent and does not walk. However this still means that there will be a Stark shift derived difference between the parity frequency and the center of the motional side-bands. The ion will incur an optical intensity dependent virtual $Z$-rotation per
gate (due to phase accumulation). The phase shift is consistent for each optical intensity / gate time setting and repeatable and the effect becomes small for low Rabi frequencies.

**Concatenated Gates**

In order to separate the true infidelity per gate from the SPAM error, we run concatenated series of gates and fit the slopes to find the fidelity. We see an example of this from the same 175 $\mu$s and 300 $\mu$s gates of Figure 6.17 in Figure 6.18. In these two examples, we see a differing level of apparent SPAM error (y-intercept of infidelity curve) in excess of our expectations from the single qubit GST result. In this system when doing qubit gates, it is not unexpected to see SPAM errors on these gate scans at the 2% level [WCF+20]. Over the course of the 32 gates used in our main gate survey the average fixed SPAM contribution to our infidelity is 2.2%. This figure is likely driven higher by the linear fits to the large parity error whose intercept is noisy. If we compare the odd population y-intercept to the parity y-intercept, we see a mean value of 1.8% and 2.7% respectively.

A noted weak point in this system is the stability of the polarization and power level of the dark pumping and detection beam. These beams are both coupled into the same fiber as the cooling beam which is combined with a non-polarizing beam splitter before coupling. They pass through a single mode fiber (not PM) on their way to the experiment. On the fiber output, we have two waveplates which we can use to adjust the polarization to provide a good balance between the cooling and Doppler cooling. Polarization drifts in the fibers are uncontrolled; furthermore, thermal hysteresis in our free-space EOMs can cause large polarization drifts in these beams. At the beginning of the experimental day, the state detection and dark pumping are calibrated carefully. During operation we monitor our detection histograms and car-
rier/sideband rotations to ensure the SPAM has not drifted much; however it needs
to be readjusted frequently and may also be subject to drifts which are not easily
noticeable in these simple calibrations. Furthermore state detection itself could be
improved by fixing the lens alignment procedure. This would also decrease SPAM
error and population leakage. In practice more attention was put on gate frequency,
amplitude calibrations, and infidelities per gate rather than having perfect and consis-
tent SPAM errors across the entire dataset. Fixing these drifts which cause changing
SPAM error should be a top priority when system improvements are considered.

![Figure 6.18](image)

**Figure 6.18:** (a) A plot of the 175 μs gate from Figure 6.16 concatenated at
1,5,9,13,17,and 21 gates. (b) A plot of the 300 μs gate concatenated at 1,5,9,13,
and 17 gates. We fit the slope of to find the infidelity per gate.
6.3.3 Frequency Modulated Gates and Their Filter Functions

![Diagram](image)

**Figure 6.19:** (a) A plot of a typical robust-FM gate used in this study. We use FF solutions with 20 segments centered about the $\nu_{\text{tilt}}$ mode. (b) Both modes will undergo loops in phase space which should return to the origin after the interaction.

Recently many modulated gates schemes have been developed which are robust against motional mode drift [LLF$^{+}18$, KLZ$^{+}21$, MEH$^{+}20$]. These gate schemes take advantage of the ability to modulate the control pulses in an optimized fashion such that the range of acceptable drift is larger. Here we experimentally test a diverse population of robust-FM gates which are composed of discrete pulse sequences which are optimized to minimize the time average displacement error [LLF$^{+}18$]. In Figure 6.19(a) we see an example of an FM gate time. The gate time is divided into 20 discrete segments whose position in frequency space are free parameters for the optimization. In Figure 6.19(b) we see the path in phase space that each motional mode will take during the gate interaction. Ideally, after the gate time elapses both modes should return to the origin. For long chains of ions, the segmented pulse provides many degrees of freedom to optimize gates which will not leave residual population in the other detuned modes.

We also test the ‘f-robust’ FM gate scheme. The f-robust FM gate has been optimized using the filter function (FF) formalism [GSUB13] to increase robustness
to time-varying mode frequency drift [KWF+22]. Kang et al. notes that much work in this area [GB15, MEH+20] focuses only on the displacement FF \( F_\alpha \) for robustness optimization.

The time-dependent motional frequency drift can be expressed as \( \nu_k(t) = \nu_k + \delta(t) \) where the power spectrum of frequency noise is \( S_\delta(f) \) [KWF+22]. To first order, in the regime of weak noise, the infidelity can be calculated by evaluating a set of FF integrals [KWF+22]:

\[
E = E_\alpha + E_\Theta = \int_{-\infty}^{\infty} \frac{S_\delta(f)}{f^2} F_\alpha(f) df + \int_{-\infty}^{\infty} \frac{S_\delta(f)}{f^2} F_\Theta(f) df \tag{6.1}
\]

The filter function for angle \( F_\Theta \) can dominate at frequencies below \( 1/\tau \) [KWF+22]. Upon inspection, the typical \( F_\Theta(f) \) multiplied by the factor \( \frac{1}{f^2} \) can be seen to have a flat response until 10 kHz, where there is a deep valley, followed by a more complicated landscape trailing down in significance. In this frequency range \( F_\alpha(f)/f^2 \) begins to dominate. Here we take advantage of an environment dominated by strong coherent low frequency noise to explore the physical meaning of the low frequency parts of these FF bands, and how accurately they can be used to describe experimental data in this regime.

The robust-FM gate optimizer is based on the scheme of Leung et al. [LLF+18, Leu20] implemented originally to demonstrate high fidelity gates in ion chains [WCF+20]. The specific implementation was created by Bichen Zhang and Shilin Huang [WCF+20]. Our robust-FM gate optimizer requires the specification of two free parameters: the gate time and the Rabi frequency limit. A limit is imposed on the upper allowable Rabi frequency that a solution should require. Along with the Rabi frequency limit the gate time is given and then a solution will be generated from the inputted motional mode frequencies. The solver will then optimize the gate according to the scheme in
[LLF+18] and then return a discrete frequency modulated sequence with minimized gate error. It will also prescribe a motional sideband Rabi π-time for which the gate is designed. Our solver analytically calculates the trajectory in phase space [Leu20] and returns a predicted gate error rate based on the ending displacement in phase space. To first order this displacement should be equal to the $E_\alpha \approx p_{01} + p_{10}$ when calculated from the optimizer. Analytically we call it ‘gate closure error’ and it should be a good predictor of the odd population error per gate observed experimentally. Our optimizer allows the user to set shorter gate times at the expense of an imperfect phase space loop closure by changing the Rabi frequency limit. In Figure 6.20 we utilize the simulation package of Mingyu et al. [KWF+22] to generate filter functions for particular gates in our survey. We see how the shapes of the FF can vary for particular robust-FM gates, and the relative behavior of different bands of the angle and displacement filter functions.

Figure 6.20: Examples of filter functions calculated from robust-FM gates used in this study (a) We plot the $F_\alpha$ for gates with very different analytically predicted displacement errors ($|\alpha^2|$); these displacement errors are reflected in the low frequency behavior of the filter functions (b) These gates have fairly similar low frequency behavior in $F_\Theta$ despite the large differences in displacement error; they are all similar gate times (120 µs, 135 µs, 150 µs). (c) Here we look at the low frequency behavior of $F_\Theta/f^2$ for gates with a larger difference in gate times but the same imperfect displacement error (120µs,215µs,300 µs).

Examining Figure 6.20(a) more closely we see the effects of optimizing a gate with imperfect closure error by setting a Rabi frequency limit which is sub-optimal for the
chosen gate time. The gate times used here are very similar (120\,\mu s, 135\,\mu s, 150 \,\mu s) but the Rabi frequency limit can create a very different FF. The low frequency portion of a $F_\alpha(f)$ with imperfect closure error will flatten out to a particular value at DC. In Figure 6.20(b) we see that plotted on this scale, the behaviors of the $F_\Theta(f)$ appear much less different. At the $10^4$ Hz level (where we draw a dotted line) in both of the filter functions we see that there is a change in structure, where sharp resonances begin to appear; above this line $F_\alpha(f)$ will dominate by comparison.

In 6.20(c) we see an example of how the $F_\Theta(f)$ behaves when multiplied by $1/f^2$ as in Equation 6.1. $F_\Theta(f)/f^2$ is always flat in the $f < 10^4$ band, and the level at which it flattens can be quite different. Do not be fooled by the low suppression of these frequencies; this function only translates to an error when integrated over a noise spectrum as in Equation 6.1. If you have a frequency noise PSD with peaks on the order of $10^4\, rad^2/Hz$ then you will have values on the order of $10^{-4}$ integrated over a band. With a lot of noise power in the band these errors add up. When looking at these plots we may expect that in this low frequency band $F_\alpha(f)$ would still dominate over $F_\Theta(f)$ when integrated over the same noise PSD. However in this work we provide compelling experimental and simulated evidence that in the regime where coherent low frequency noise dominates over the $T_2$ white noise motional dephasing time, the displacement filter function is unimportant and our gates are still robust to population errors, even with imperfect predicted gate closure error.

### 6.3.4 Power Spectrum of Frequency Deviations Obtained by Electrical Measurement

The motional frequencies for radial motion in a Paul trap is proportional to $\nu_{x,z} \propto V_{RF}$ the RF voltage amplitude [SSWH12]. For the axial confinement, it goes approximately as $\nu_y \propto \sqrt{V_{DC}}$ [DSW06]. Exactly calculating this for surface traps is com-
plex [HLC+16]. The proportionality still holds experimentally. We can write the time dependant RF voltage as:

$$V_{RF} = V_{0RF} + \Delta V_{RF}$$  \hspace{1cm} (6.2)$$

where $\Delta V_{RF}$ is a small noise deviation in $V_{RF}$. We can write for the radial motional frequencies then $\nu_{x,z} + \Delta \nu_{x,z} = \xi(V_{0RF} + \Delta V_{RF})$, where $\xi$ is the constant of proportionality such that $\nu_{x,z} = \xi V_{RF}$. We can solve for $\Delta V_{RF}$ and will see that up to a constant $\xi$, the deviations in radial trap frequencies go as $\Delta \nu_{x,z} = \xi \Delta V_{RF}$. We define the time dependent behavior of $\Delta V_{RF}$ as $V_{\text{noise,RF}}(t)$, and we define $\delta_{x,z}(t)$ as the continuous time dependent functional form of many aggregated $\Delta \nu_{x,z}$ variations. We then see that the time dependent frequency deviations will go as $\delta_{x,z}(t) = \xi V_{\text{noise,RF}}(t)$. The power spectral density (PSD) of frequency deviations for the radial modes are then $S_{\delta_{x,z}}(f) \propto S_{V_{\text{noise,RF}}}(f)$ because the source of their deviations in time is this voltage noise.

For the DC axial confinement we have the relationship $V_{DC} = V_{0DC} + \Delta V_{DC}$ implying $\nu_{y} + \Delta \nu_{y} = \lambda(\sqrt{V_{0DC} + \Delta V_{DC}})$. We can solve this relationship for $\Delta \nu_{y}$ and Taylor expand for small values of $\Delta V_{DC}$ to obtain:

$$\Delta \nu_{y} \approx \frac{\lambda \Delta V_{DC}}{2 \sqrt{V_{0DC}}} + O[\Delta V_{DC}]^2$$  \hspace{1cm} (6.3)$$

We define $\delta_{y}(t)$ as the continuous time dependent functional form of many aggregated $\Delta \nu_{y}$ variations. Similarly to the radial modes, for the axial mode this implies that:

$$\delta_{y}(t) = \frac{\lambda}{2 \sqrt{V_{0DC}}} V_{\text{noise,DC}}(t)$$  \hspace{1cm} (6.4)$$

We should also then see that the power spectrum of axial frequency noise goes as:
\( S_{\delta_y}(f) \propto S_{V_{\text{noise,DC}}}(f) \). A spectrum analyzer can measure the power spectrum \( S_V(f) \) giving units of electrical power \( P_{AC} = |V|^2/R \) over a particular measurement bandwidth. If the noise is wide-sense stationary over the measurement time, this can serve as a good estimate for \( S_V(f) \). Thus we claim that if the electrical power spectrum as seen by the ion could be obtained directly, it follows that, up to some constants \( \Xi \) and \( \Lambda \) the PSD of frequency deviations for the radial and axial modes go as:

\[
S_{\delta_x,x}(f) = \Xi S_V(f) \quad S_{\delta_y}(f) = \Lambda S_V(f) \tag{6.5}
\]

Chains of ions will have higher order motional modes which will differ up to a constant as well as having some dependence to both the axial and radial modes. We cannot use a spectrum analyzer to directly measure the noise which the ion sees. If we suspect that the dominant electrical noise source is from currents flowing through system grounds, we can make a measurement of the power spectrum driven by these ground currents over a measurement resistance. The PSD obtained from the ground currents can serve as a model for the PSD of frequency deviations if we fit the constant of proportionality. This spectrum can be used to generate the motional phase noise which the ion will see.

Later in Section 6.3.6 we will fit the strength of this PSD model using quantum Monte Carlo simulation of the motional Ramsey experiment. Using this method we can estimate the strength of the noise. Directly modelling this from first principle calculations would require an electrical model of the entire device network, as well as finite element modelling of the surface ion trap. If this simple method can be validated for certain cases, it could prove to be a valuable diagnostic tool. For this method to be valid, one must have good reasons to suspect that the white noise motional dephasing contribution is much smaller than the noise contribution of the
measured spectrum. In the next section we will review core parts of our system electronics and then we will explore low frequency noise coupling mechanisms.

6.3.5 The Device Network and Noise Coupling Mechanisms

A complex system such as an ion trap quantum computer can be thought of as a device network. Many individual electrical components must be all connected together and work in harmony towards a single goal. If we isolate each device and perform noise characterization, we may miss important effects which are the result of the network connections. Low frequency electronic noise is particularly susceptible to being enhanced by interconnection. It is widely known in audio-visual (AV) engineering circles that unbalanced interconnect systems are subject to an inherent injected ‘hum’ (in the neighborhood of 60 Hz) and ‘buzz’ (their higher harmonics)[Whi96]. When sensitive signals share a common ground with AC power devices, the hum and buzz noise can be injected onto the signal by the shield grounds which connect devices. These problems are particularly important in AV systems because hum/buzz at -80 dB levels irritate most listeners [Whi08]. Connecting a coaxial cable (from residential cable TV systems) into an AV device network could create as much at -6 dB hum problems without ground isolation [Whi08].

In ion trapping the noise level of devices like digital-to-analog converter (DAC) boxes are typically characterized independently. As noted earlier, measuring the exact electrical noise seen by an ion is difficult. In this section we outline our device network briefly and then expand more on various mechanisms in which this type of noise can couple onto our surface trap. One issue presented may have significance to anyone operating a Gifford-McMahon cryostat to trap ions.
The Device Network

Here we provide more detail about the electrical interconnections in our system. A simplified schematic diagram of the system is presented in Figure 6.21. We use a 100 channel DAC box discussed in Mount et al. [MGV+16].

Our 100 DAC channels are delivered to the Montana system by 4 shielded 25 pin D-sub cables which are each wrapped tightly in a ground strap shield along their entire length. These shields are connected to the DAC box and chamber grounds via copper tape which aims to enclose them entirely at the fringes. The addition of these extra ground shields was correlated with a reduction in heating rate from 90 quanta/s to the 10 quanta/s regime. We have not since removed these shields to confirm this.

The 4 D-sub cables enter a filter box at the Montana system input which both serves as an adapter from 4 25-pin D-subs to 2 50-pin Mini-D Ribbon (MDR) connectors which are the DC feedthroughs for the Montana system. The filter box features RC filters on each channel with a 2 kHz cut-off frequency. Internally the Montana DC feedthrough adapts the 50 pin MDR connectors to two Omnetics Nano-D connectorized flex PCB ribbon cables. These cables are unshielded and terminate onto the trap PCB which sits rigidly attached to the 5K cold finger sample mount. The trap PCB has a second set of RC filters for each channel with a 2 kHz cutoff frequency.

There is a direct digital synthesizer (DDS) and Digital proportional–integral–derivative (PID) lock box which drives our trap RF. The DDS output goes into a high power RF amplifier and then into a four port RF coupler, which is mounted to the aluminum structure depicted in Figure 4.9. The coupler output proceeds through a semirigid coaxial cable to to the Montana RF feedthrough.

Inside the cryostat sample chamber semirigid 50 Ω coaxial cables are used along a thermal gradient which connects to the helical resonator sitting at 90 K. The helical
resonator output has a ground shield made of large diameter copper pipe for much of its length but has a 1.5 inch discontinuity where a bare copper wire makes the connection to the trap PCB. A ground connection between the helical resonator and the trap PCB ground plane is made using a 2 inch long stainless steel wire with a 10 mil (0.254 µm) diameter.

The Sumitomo cryostat has a two stage heat pump. Stage 1 nominally sits at 30 K and is connected strongly to the experimental platform which sits at 90 K (with our heatloads). Stage 2 nominally sits at 5 K and is connected strongly to the cold finger sample mounts which sits at 8 K with RF applied. Inside the Montana system there is a strong electrical ground connection from stage 1 of the cryocooler and the room temperature chassis, the 90K platform, radiation shield, and all cryocooler electronics. The 5K cold finger is thermally isolated from the rest of the system as built by Montana instruments and no explicit electrical connection exists except ones inside the Sumitomo cryocooler module itself. Thus explicit ground connections made between the 90K platform and the 5K sample mount (and ion trap package) are added by our experimental sample chamber design. Some of these connections include the steel wire ground connection, as well as a semi-rigid coaxial cable which routes a capacitive pick-off from the trap PCB to the exterior. This is only a small fraction of the entire device network sharing a ground with our ion trap; however these are the core components necessary to understand a large part of the noise coupling mechanism which we delve into below.

**Noise Coupling Mechanisms and Electrical Noise Measurements**

Here we highlight some mechanisms which allow low frequency noise to couple to the ion. Electrical noise in this frequency band can couple into the experiment by a variety of mechanisms. Many shielding techniques which provide good shielding
for radiated EM noise at high frequencies can exacerbate noise problems at low frequencies [Ott11]. Many myths exist related to grounding; chief among these is that connecting to an earth ground is necessary for low noise operation [Ott11]. A corollary to this is that currents which flow to earth ground are ‘gone’. Currents flow in loops and as such currents that flow into the building/earth grounds must eventually flow elsewhere [Ott11]. Many mundane devices in common use today will inject noise back into the power grid and this noise will vary in time [VTGO84].

Of particular interest in our system is two mechanisms for noise injection: common impedance coupling, and induced ground currents. Inverter driven AC motors power the mechanical action behind the Gifford-McMahon cryostat’s heat pump cycle. When a device network is connected with a shared ground, especially when involving AC power circuitry, a low frequency current will circulate between the two systems through this shared ground [CBBZ16]. It is widely known that inverter powered motors drive a common mode voltage between the motor windings and ground [SM04]. These inherent common mode currents drive noise currents into the system ground and can be a source of electromagnetic interference (EMI) [ZH13].

This current noise from the AC motor can be injected onto signal line by the interchassis currents. These currents can be also caused by the charging and discharging of real and stray capacitance between AC power circuitry and the device grounds [Whi96]. These hum/buzz currents can be exacerbated by the routing of interconnects in ways that form loops which can pick up AC magnetic fields, especially if these loops are located near-field to AC motors, transformers, and AC power wiring [Whi96]. Additionally the ‘pin 1 problem’ can occur. The pin 1 problem is a name for a device design flaw in which shield grounds are terminated onto the ground planes of signal circuitry PCBs [Bro05]. These ground currents can be injected then into the signal circuitry. They will cause voltage drops to develop across their path.
to the equipment and power grounds [Bro05]. It is actually better to terminate shield grounds into the system chassis than into the signal circuitry ground plane.
Figure 6.21: A schematic diagram of our experimental device network. (a) Inter-chassis current can couple into the signal by common impedance coupling. AC current noise from other parts of the experiment will be injected in the the DAC box ground (b) Underneath the optical table 60 Hz noise is detected by a loop antenna. Loops between the experimental structure and the DAC box itself can increase the noise amplitude. (c) A strong 70 Hz noise and higher harmonics can couple into our trap PCB and trap ground. Ground lines from the 90K platform to the cold finger will create a loop with respect the cryocooler itself.
We highlight two possible mechanisms for our particular system in which these ground currents can effect the ions. The first is the interchassis current mechanism where these noise currents will be injected through the DAC box signal channels. These signal injections will be filtered below 2 kHz by our RC low-pass filters. The second mechanism is the injection of ground currents through ground loops which pass through the trap PCB ground plane. Low frequency ground loops exist where the trap ground can experience noise with respect to the DC and RF signals.

The first major ground loop being that we connect our capacitive pick-off circuit to the exterior via semirigid coaxial cable. Its shield ground terminates on the trap PCB, and is thermally lagged both at 90K and 5K inches away from the sample mount forming a loop in which current can be injected into the trap ground. Another instance of this is that we connect our helical resonator to the trap PCB ground plane with stainless steel wire of 0.254 mm diameter which has a low resistance at low frequencies but > 6 Ω resistive impedance at RF frequencies.

For thermal reasons, the ground shield of the helical resonator RF output cannot be strongly tied to the trap PCB. There is no explicit ground connection between the DAC box and the trap PCB. The DC cable ground shield terminates on the cryostat exterior. Thus the RF and DC signals area are referenced strongly to the room temperature chassis and the 90K platform as well as their own enclosures. The 90K and 5K platforms are only connected electrically inside the Sumitomo cryocooler. Thus any additional low frequency ground connections to the trap PCB from the 90K platform will create a large loop in which currents will flow through the trap ground. While noise will not be injected directly on the trap RF signal line because the helical resonator is a narrow band-pass filter, there is a potential difference at RF frequencies between the helical resonator ground (on the 90K platform) and the traps ground which will have low frequency currents passing through it. Additionally
because the DAC box ground terminates on the room temperature chassis, it will also have a potential difference with respect to the trap ground.

We think of the trap ground as the electrical reference frame which the ion will see. The voltage drops created due to these noise currents will then make the trap ground oscillate with respect to the RF and DC signals (referenced to more distant grounds). In Figure 6.21 we see this principle illustrated as resistors added on the path between the system chassis (of the Sumotomo cryocooler which is connected closely to the AC motor) and the ground points where the DAC box and RF are brought in. As the AC motor drives current into the ground, some of this current will travel through the cold-finger, trap, PCB, and out of the system through the DAC system and PID/DDS system equipment grounds.

Additionally we used a loop antenna connected to a DC-coupled spectrum analyzer to inspect the magnetic field noise around the experiment and sample chamber. Placed on top of the cryocooler of the Montana system (Figure 6.22(a)), there is a predominance of 70 Hz peaks and their higher harmonics. Underneath the optical table, we place many power supplies and two large power strips, but also our DAC and DDS systems which drive the trap DC and RF. When the loop antenna is held under the table (Figure 6.22(b)), it picks up a prominent 60 Hz peak and its higher harmonics. When the loop antenna is held directly over the sample chamber (Figure 6.22(c)), we can view both series of 60 Hz and 70 Hz peaks. The cryocooler motors emit 70 Hz (and all higher harmonics) magnetic field noise which could be picked up on our near-field ground loops and sent into the trap ground plane.
Figure 6.22: We use a loop antenna connected to a DC coupled spectrum analyzer to inspect the spectrum of magnetic field noise in the experiment. (a) We place the loop antenna on top of the cryocooler module of the Montana system and observe a very strong 70 Hz peak and a series of its higher harmonics. (b) We place the loop antenna underneath the optical table, where the DAC system, the DDS system driving the trap RF, and many power supplies are placed. We observe a 60 Hz peak and its higher harmonics. (c) We place the antenna atop the cryostat sample chamber and we observe both sets of peaks clearly.

In addition to the inductive pick-up, there will be a contribution directly driven into the trap ground due to the current noise injected by the inverter driven motor,
transformers, and other AC circuitry of the cryocooler system. Some of this noise will also be injected into the DC electrodes directly via the shared impedance coupling between the DAC box and the sample chamber. This DAC box hum and buzz noise will be low-pass filtered at 2 kHz.

Directly calculating which of these contributions will be most important and exactly what voltage the ion will see would require a complex circuit model of the entire device network. To understand exactly how this noise would effect the ion would require finite element modelling of the surface trap potential under these dynamic conditions. Instead we will measure the spectrum of the noise referenced between two key locations in order to model the spectral distribution of this noise. We will scale this noise model, as justified earlier, to generate a best guess of the power spectrum of frequency deviations which the ion will see. We can then adjust the prefactors Ξ or Λ in a quantum Monte Carlo simulation of the motional Ramsey experiment to match our observed dephasing. We can then check this hypothesis against our experimentally observed MS gate behavior to see if our errors are well modelled by the hypothesis.
Measuring the Noise

Figure 6.23: (a) We plot the $F_\Theta(f)$ filter function scaled by $1/f^2$ until the first steep dip which occurs near 10 kHz and the response is flat. We are interested in examining this region for our work. This feature is unique to $F_\Theta(f)$ as opposed to $F_\alpha(f)$ and so it can be integrated (times a noise PSD) to obtain a constant which doesn’t depend on the choice of lower bound. (b) Our hypothesized noise PSD which is responsible for the majority of dephasing, obtained by a measurement of the electronic power spectrum at two grounded points. We have scaled the spectrum by the parameter $\Xi$ obtained by our motional coherence simulation discussed later in this section to make this the hypothesised frequency noise PSD. Later we test this hypothesis by seeing how well it can explain our parity errors.

In order to measure the low frequency electric noise spectrum, we directly measure currents flowing in our system ground. We perform a two point measurement with a DC coupled spectrum analyzer and a 1 Hz measurement bandwidth. We adapt a BNC coaxial cable into a two probe measurement tool and attach the probes between the Sumitomo cryocooler chassis ground and the system RF feedthrough ground. The measurement will include a component caused by the voltage drop between these two points as well as an inductively coupled component onto the loop created by the measurement probes. We are not interested in precisely measuring the amplitude of this ground noise, and because we are also interested in the spectral contribution from magnetic pick-up in the system, this measurement-loop induced component is not a problem. The results of these measurements are shown in Figure 6.23(b), where we have scaled the amplitudes based on our simulation to the expected frequency noise.
power spectral peaks. The maximum power of the tallest peak is only in the -60 to 
-50 dBm range. Depending upon the locations chosen for measurement, there are 
at least 200 peaks with greater than 30 dB prominence over the noise floor of the 
instrument. In similar measurements conducted on UHV chambers in our lab which 
do not share a strong ground with the cryostat motor, there are no 70 Hz peaks, (and 
thus about half the noise power) and the overall level of the measurements are lower 
when analogous measurement positions are chosen.

Additionally an AC current meter can be used to measure the current driven to 
the building ground during operation; for our system we see an average of 2 mA with 
peaks of 5 mA. This current can be measured with respect to the cryostat chassis or 
with respect to our DAC box at a similar level. This level of current driven to earth 
ground is not seen in UHV chambers in our lab, or in another Gifford-McMahon 
cryostat which is turned off. If an AC voltmeter is applied between the building 
ground and our cryocooler, we measure a 40 mV AC voltage, while the same probe 
applied between the DAC box ground and the building ground gives a 70 mV AC 
voltage. Devices with a strong connection to building safety ground will act as voltage 
sources, while devices with a weak connection to building ground will act as current 
sources [Whi96].

In general it is a difficult problem to estimate or measure the exact noise amplitude 
that the ion will see. This noise will show up on all DC electrodes and on the RF 
signal in some complex fashion which depends upon the details of the device network, 
and the DC impedance to ground seen by the differing trap electrodes. A similar 
spectrum can be seen by directly measuring the DC feedthrough pins referenced to 
the DC shield ground on the cryostat and this signal has a 0.7 mV maximum peak 
measured over a 50 Ω load. In the real situation with the DC filter box installed and 
the full DC cabling attached to the system ground, it will support yet more ground
loops and inter-device common impedance currents which will change the amplitude seen by the ion.

The salient features of this spectrum are a series of 60 Hz and 70 Hz peaks. A long series of their higher harmonics reaching through 10 kHz, as well as a series of other peaks which are not harmonics of 60 Hz and 70 Hz whose origin are unknown. We hypothesize that the spectrum of the frequency drift seen by the ion will have these frequency components with their relative strengths, and we further claim that this is the dominant motional dephasing mechanism for this system. By using this noise profile, we can reproduce the unique shape of the motional Ramsey experiment. We can also include these noise peaks as coherent phase noise sources in a Monte Carlo master equation simulation of our FM gates. All that is left is to use these two tools to fit the amplitude of the noise by simulation.

We have tried using the method of [WUZ+17, KWF+22] to obtain the spectrum via CPMG sequences. We have found that for system, with a potentially very strong coherent noise profile, it was difficult to get obtain a scan with good contrast as the number of pules is increased. Furthermore at these low frequencies the bandwidth of the measurement FF obtainable is on the order of 70-100 Hz. For these reasons this method was not suited to providing a noise PSD with the detailed spectral information we desire at low frequencies. Some examples of these CPMG scans for n=6 are shown in Figure 6.24. The collapse and return of these scans is an indicator of a strong and broad coherent noise spectrum in the frequency range we have been interested in from Figure 6.23. We have normalized the contrast of these scans in order make them amenable to analysis.
Figure 6.24: Examples of the CPMG sequences with $n=6$ in our system. Because of the low contrast of the scan we have had to normalize the datasets to make them more amenable to spectral analysis. Because of this the amplitude is shown in relative units. They are provided as some evidence for coherent noise and as an example of why our method of fitting the spectrum may be useful for electric hum. The bandwidths for the CPMG FF’s are in the order of 70-100 Hz.

6.3.6 Motional Coherence Simulation

We have measured a motional coherence at the 3.4 ms level on both tilt and common radial modes. Experimentally we sideband cool both motional modes with the tilt mode cooled last before measurement. We perform a $\pi/2$ rotation followed by a Ramsey wait time and another $\pi/2$ rotation scanning the phase to trace out a parity curve. We fit the result to obtain a contrast for each wait time point. These results can be seen in Figure 6.25 where the errorbars are obtained by the standard error
of the sine curve fitting. We measure the tilt mode because we design our FM gate solutions to couple strongest to this mode. The decay is not a pure exponential decay as would be expected from a white noise model of motional decoherence. Instead we see a slow decay for short delay times followed by a faster decay regime for times greater than 1 ms.

We use the measured spectrum as a proposed model for the spectral components of the frequency deviations and fit the amplitude. Here we justify this for the tilt mode which is related to the common radial and axial modes by \( \nu_t = \sqrt{\nu_c - \nu_y} \).

Small deviations in the tilt mode if we assumed only axial deviations, we can write \( \Delta \nu_t = \sqrt{\nu_c - (\nu_y + \Delta \nu_y)^2} - \nu_t \), if we Taylor expand this for small deviations:

\[
\Delta \nu_t \approx -\frac{\nu_y \Delta \nu_y}{\nu_t} - \frac{\nu_y^2 \Delta \nu_y^2}{2 \nu_t^3} + O[\Delta \nu_y]^3
\]  

Similarly the expansion for small changes in \( \delta \nu_c \) is written:

\[
\Delta \nu_t \approx \frac{\nu_c \Delta \nu_c}{\nu_t} - \frac{\nu_c^2 \Delta \nu_c^2}{2 \nu_t^3} + O[\Delta \nu_c]^3
\]

Thus to first order this implies that small deviations in either the radial common modes or the axial common mode should result in a combined square deviation going as:

\[
(\Delta \nu_t)^2 \approx \left(\frac{\nu_c \Delta \nu_c}{\nu_y} - \Delta \nu_y\right)^2
\]

So we see that the contribution from the common mode spectral noise is much stronger. If the noise was only experienced as a clean axial motional frequency noise, then we should expect a substantially different behavior from the tilt modes and common modes; however, we experience a similar dephasing time and functional form for both the tilt and common modes. Thus we apply the logic of Section 6.3.4
and use the power spectral density inferred from the measurement of ground noise.

**Figure 6.25:** A quantum Monte Carlo simulation of our motional Ramsey experiment. We generate 100 realizations of frequency noise according to the proposed noise PSD, which was turned into Fourier amplitudes, then multiplied by a random vector of phases. Averaged over many realizations, this noise PSD produces a frequency noise with a standard deviation of 73 Hz when we obtain a fit for a prefactor of $\Xi = 9000 \frac{\text{rad/s}^2}{\text{Hz}}$.

We take our spectrum analyzer trace, convert it from dBm to linear units of power. Then we normalize the spectrum and multiply it by a prefactor we call $\Xi$. We interpolate the spectrum from 5 Hz to 10 kHz with an approximately 2.5 Hz pitch, chosen because it divides evenly with our dominant 60 Hz and 70 Hz peaks, resulting in 4000 samples. Using the method presented in [KWF+22] we use this proposed power spectrum to generate Fourier amplitudes. We create a uniform random vector of phases, and then perform the inverse Fourier transform to create a time-trace of random frequency noise. We run a simulation of the motional Ramsey experiment using QuTiP, a python-based quantum simulation package, and inject the generated frequency noise during the free evolution. We do this for 100 random realizations for frequency noise.

From the simulations we generate a contrast decay curve. We then repeat this
experiment until we fit the overall prefactor $\Xi$ to an amplitude which fits to our experimental contrast decay. We observe a good fit for $\Xi \approx 9000 \frac{(\text{rad/s})^2}{\text{Hz}}$. Averaged over many realizations, this noise PSD, scaled by this factor of $\Xi$, results in real frequency noise with a standard deviation of 73 Hz. This frequency noise amplitude is surprisingly small, but we can test the hypothesis that such a frequency noise PSD can explain our dephasing by seeing how well it explains our MS gate errors.

### 6.3.7 A Survey of FM Mølmer–Sørensen Gates

We sought to test behavior of our gate errors over a range of system parameters. We wanted to verify that the observed dephasing time of 3.4 ms was indeed the main source of gate error. In a typical white noise model, this would be a very low motional coherence and our gate fidelity would be a strong function of the ratio between motional coherence and gate time. Furthermore we sought to examine the true sensitivity of our robust-FM gate odd population error to the calculated displacement error. To test a diverse population of robust-FM gate solutions, we sought a wide distribution in calculated gate closure error, blue sideband $\pi$-time, over a range of gate times.

In order to create this population of gates, we moved through an even distribution calibrated Rabi blue side-band $\pi$-times between the shortest achievable without MEMs power limitation and 450 $\mu$s. Within each neighborhood of sideband $\pi$-time we selected a few gate times to try, one which was too short to achieve good predicted displacement error within our Rabi frequency limit, one which was slower than necessary (so that the optimizer easily finds a perfect solution), and one in between. We found that if the predicted displacement error was $> 1$-$2\%$ the gate would be clearly ill-formed and amplitude calibration could not be performed. This would be
evident in the detuning scan whose $p_{01} + p_{10}$ minimum could never overlap with the point of $p_{00}$ and $p_{11}$ crossover. We limited ourselves to gates which produced an ideal overlap between the odd population minima and the crossover of $p_{00}$ with $p_{11}$. This behavior did seem like a threshold, in that there was an abrupt change in the nature of the detuning scan after this approximate predicted error level. The tolerance to imperfections in the predicted displacement error was surprising.
Figure 6.26: (a) An experimental survey of a diverse series of robust-FM gates was conducted. Each point is a series of odd-numbered concatenated gates (b) We plot these unique gates against their predicted odd population errors and find that there is not strong correlation between the odd population errors and the predicted displacement errors. All of the gates are sufficiently robust against these types of errors for our noise model.

In Figure 6.26 we can see the results of the gate survey. Each unique FM solution was characterized by odd numbers of successive MS gates we fit linearly to find the infidelity, odd population, and parity error per gate (as in Figure 6.18). All numbers quoted on scatter plots are these fitted slope errors. We have attempted to concatenate up to 21 gates when possible but some series were stopped at 13 or 17
due to calibration drift or large parity error at that point. We conducted all of these
gate sequences in a 5 day period back-to-back with a calibration routine similar to
Wang et al. [WCF+20]. The main analyzed population is a set of 32 robust-FM
gates which can be seen in Figure 6.26. We also tested a set of 5 f-robust gates as
described in Kang et al. [KWF+22] which we only include as a comparison in Figure
6.28. Each FM gate used was composed of 20 discrete segments. The main features
of the robust-FM gate survey in Figure 6.26 are a generally low odd population error
per gate with a mean of 0.3% per gate and a high parity error per gate with a mean
of 3.8% per gate. The mean infidelity per gate was 2% which means we have an
average gate fidelity of 98%. On Figure 6.26(a) we see a weak trend of improvement
with gate speed in the 100-250µs time-span with a flattening of the response above
this time. There is a large variation in parity error among the gates with a standard
deviation of 0.8%.

In Figure 6.26(b) we can see a surprising insensitivity to the analytically calcu-
lated displacement error. As a sanity check we had a gate with a 2.5% predicted
displacement error, and at this point there is a response. We don’t include this
dataset in other plots. It would seem our system is very insensitive to these errors
and imperfect FM solutions are well tolerated. In Figure 6.27 we see our population
of gates plotted as a function of the blue sideband π-times at which we calibrated
the gates. Again we find a surprising lack of correlation directly between the blue
sideband π-times and the quality of gate by any of the three metrics.
Figure 6.27: The gate errors plotted as a function of blue sideband $\pi$-time for gate survey, we notice a lack of correlation.

6.3.8 Filter-function Transmitted Noise Profile and f-Robust Gates

Finding no strong correlation to gate time, sideband $\pi$-time, or calculated displacement error, we sought to sort the solutions by the properties of their filter functions. We noticed that the only strong correlation in the data was to the low frequency (< 10 kHz) portion of the $F_{\phi}$. To test our hypothesized frequency noise PSD we then solved the integral for $E_{\phi}$ and used the relation $1 - c = 2E_{\phi} + p_{01} + p_{10}$, where $1 - c$ is the parity contrast error. We used the experimentally observed odd population and this FF calculation to predict the parity error contribution from this low frequency portion of the filter function and our PSD. Fitting a line to this parity scatter plot we see a slope of 0.66 with an intercept of 0.027. The Pearson’s correlation coefficient we calculate for the fit is $r=0.7$ (with $r=1$ being a perfect linear model) with a p-value of $9.4 \times 10^{-7}$. Thus we find that this low frequency portion of the $F_{\phi}$ underestimates, but predicts a large fraction of, the observed parity error. This also indicates that our prediction of $S_3(f)$ is a good model for the noise. Using the calculated $E_{\phi}$ directly models the infidelity with a slope of 0.7 and an intercept of 0.016. The Pearson’s
correlation coefficient is \( r=0.56 \) with a p-value of 0.00027. With a full master equation simulation we can make sense of the additional errors by using a more complete model.

![Graph](image)

**Figure 6.28:** We integrate over the PSD by evaluating \( S_\delta(f) F_\Theta / f^2 \) for each experimental filter function (x-axis) to predict a parity error contribution. We plot this against our experimental errors and fit to find correlation. We see that this proposed noise PSD is a good model of our observed parity error. We also include our observed odd population error using the relation \( 1 - c = 2\epsilon_\Theta + p_{01} + p_{10} \). The Pearson’s r-value is 0.7 with a p-value of \( 9.4 \times 10^{-7} \) indicating this correlation is extremely unlikely to have occurred by chance. While this hypothesized PSD and the \( F_\Theta \) alone cannot predict all of the parity, it is good evidence that this type of noise is responsible for the majority of our errors. We also include 5 unique f-robust optimized gates into this population showing that, as expected, these gates are consistently robust to this unique error source.

### 6.3.9 Monte Carlo Simulation of the Gate Survey

Our gate survey provided a unique opportunity to test the hypothesized noise PSD with a real population of gates. We were able to make sense of the gate survey only by plotting it with respect to the \( F_\Theta(f) \) integrated over our hypothesized noise
spectrum. However this does not capture all of the errors, our PSD stops at 5 Hz and
and it cannot account for stochastic DC drifts. We also exclude other errors in our
gate such as optical and motional dephasing, as well as the motional heating rate.
There can also be a stochastic error in our Rabi frequency (due to optical intensity
drift experiment-to-experiment) and trap frequency drift. We include these errors
along with our spectrum of coherent noise in order to fully model the gate survey.
In order to properly include coherent noise, we must use a Monte Carlo method to
randomize the phases. This Monte Carlo will also make it easy to include stochastic
random errors and the coherent noise.

**Simulation Details**

We use the same full master equation (ME) FM gate simulator as [KWF+22] to
model the system with our particular error input. Similarly to [WCF+20] we use
the ME in Lindbladian form to simulate our decoherence processes. The relevant
decoherence rates simulated use the heating rate and laser dephasing as seen in
Figure 5.11. We find that aside from a few outliers, the $\tau_m = 27$ ms gives a reasonable
reproduction of our odd population over the gate population. To simulate the low
frequency coherent aspect of the dephasing, we utilize the PSD fitted in Section 6.3.6
to generate a frequency noise trace from a uniform random vector of phases by inverse
Fourier transform. This phase noise is generated uniquely for the time spanned by
our simulated concatenated gates. We can see an example of one realization of
frequency and phase noise for the time spanned by 21 of our longest gates in Figure
6.29. While the level of the phase noise seems large over the course of 21 gates, the
phase drift during the timescales of a single gate is small. We use a phase-insensitive
geometry for our gates [IVH+14] and so phase noise of the addressing beams relative
to the motional modes do no effect the spin phase and are only effect the motional
phase [IVH+14]. The tilt mode which we use primarily for our gates will move collectively for the two ions. Drifts in motional phase between gates can be tolerated because the motional phase has no effect on the final state [LBD+05].

Figure 6.29: Frequency and phase noise generated from our spectrum over 21 450\(\mu\)s gates. We have completed the inverse Fourier transform on a Fourier spectrum obtained by our hypothesized PSD, multiplied by a randomly generated vector of phases. While the phase drift is large over the timescales of many gates, it is nearly static during the duration of a single gate. We can tolerate phase drift gate-to-gate because the motional phase has no effect on the final state [LBD+05, IVH+14]

For each number of concatenated gates, the simulation results are averaged over 100 realizations of phase noise. Just as in the experiment, we fit to the infidelity, odd population error, and parity error per gate and we simulate a series of odd numbered gates up to 21. We can see an example of the simulated concatenated gates under Monte Carlo in Figure 6.30. This noise spectrum generates noise which on the timescale of a single gate is almost constant, but over the timescale of a circuit, it will have an interesting drift. We find that only two other randomized
errors are required to fit the data to our experiment which are not captured by the FF integral of Section 6.3.8. One aspect is that at the beginning of each gate sequence, we generate a Gaussian distributed random mode frequency offset with standard deviation of 220 Hz. The magnitude of this error is reasonable given long term motional mode frequency scans, and analysis of Section 6.3.1. This randomized DC motional frequency offset arises from a sampling of the frequency drift profile experiment-to-experiment by the control system which has a Gaussian distributed temporal jitter of the experimental start time. We also add a Rabi frequency noise with a standard deviation of 1.5% per experimental realization.

**Figure 6.30:** We simulate each particular FF solution as a series of concatenated gates up to 21; we then fit the linear regime to a slope for each error type. We average over 100 realizations of our stochastic and phase noise.
Simulation Results

We simulate with just our measured error sources. We use an optical coherence of 500 ms as measured in 6.2. The heating rate use is 0.44 quanta/s on the tilt mode and 8 quanta/s on the common mode as measured in Figure 5.10. For the motional decoherence error term we find that using a 3.4 ms motional coherence greatly overestimates the odd population error. This result can be seen in Figure 6.31. This is shown here as more indication that our decoherence process does not represent a typical $T_2$ time of $\tau = 3.4$ ms.

![Figure 6.31](image)

**Figure 6.31:** Simulation of the gate survey population with $\tau = 3$ ms for the motional dephasing time.

The results of a fitting simulation can be seen in Figure 6.32. In Figure 6.32(a) we see a trace analogous to Figure 6.26(a) where each FM solution simulation is scattered over the gate time. To directly compare it to the experimental data, we have included 6.32(b) and (c) which have the simulated data on the x-axis predicting the experimental points on the y-axis. The simulation predicts the trend of the parity with a Pearson’s correlation coefficient of 0.5 with a p-value of 0.003. The fidelity fit has an $r = 0.46$ with a p-value of 0.006.
Figure 6.32: A Monte Carlo master equation simulation of the gate survey. (a) The immediate simulation results for each FM solution explored in our experimental study. These points were simulated at odd numbered concatenated gates and the errors per gate fitted. Each point results from 100 realizations of the randomly generated time trace of phase noise generated from our hypothesized PSD, then added to our motional modes during the evolution. (b) A comparison of the experimentally (y-axis) and simulated (x-axis) parity errors for each unique FM solution. The parity fit has a Pearson’s r-value of 0.5 with a p-value of 0.003 such that the correlation is extremely unlikely to have occurred by chance. (c) The fidelity fit has an r=0.46 with a p-value of 0.006. (d) Rather than a linear fitting correlation, we just compare these values by eye where the majority of the odd population errors track at the level of our simulation. The reason for the outliers is unknown. A motional coherence of $\tau = 3.4$ ms overestimates the odd population errors under a Lindbladian white noise model.
In order to test the power of the model we have fitted lines with a slopes of 1 and the y-intercept as a free parameter (lines shown in Figure 6.32(b) and (c)). The y-intercepts are fitted to approximately 0.003 for both the parity and fidelity curves. We calculate the residuals with respect to this model and run statistical tests of normality in which the null hypothesis is that the residuals are normally distributed about the line. We conduct the Shapiro-Wilk test [SW65] and calculate a p-value of 0.58. We then cannot reject the null hypothesis of normal distribution at an $\alpha = 0.05$. Additionally we conduct the D’Agostino’s $K^2$ test for normality [d’A71, DP73]. This test is specialized at detecting departures from normality due to skewness or kurtosis and so it is an interesting metric to quantify our simulation’s fit to the $m=1$ model. A good simulation of this noisy data set would not have skewed residuals. We find a $p=0.9$ for this test indicating that it is likely sampled from a normal distribution.

There is a wide variation in residuals, but there is a good reason for this. When calibrating the FM gates experimentally, a detuning scan must be performed. The scans for gates with this type of noise have a band with many crossings, especially for larger numbers of concatenated gates. These crossings will change scan-to-scan, and the landscape becomes very complicated when scanned more finely. Thus in practice under this strong low frequency coherent noise, simple algorithms used to choose the detuning point for a gate in $[WCF+20]$ are unstable and cannot be used. The detunings are chosen manually. Choosing the correct detuning point is often difficult and this is a large source of fixed parity variance dataset to dataset.

Simulating this fixed calibration error which manifests itself in solutions uniquely in a way which exactly matches the data is impossible. This exercise is like throwing dice, however the normal distribution of residuals about the $m=1$ line provides some confidence about the predictive power of this model. This is good evidence that the hypothesized and scaled noise PSD derived from measurement can explain the
majority of our gate errors. Confidence is added when we consider the goodness-of-fit of Section 6.3.8.

We sought to understand the relative importance of the low frequency noise during a single gate experiment versus the sampling of the coherent noise profile at random places over the course of 100 experiments. This sampling experiment-to-experiment, we have modelled as an initial Gaussian distributed motional frequency offset as discussed earlier. Furthermore we wanted to test the repeatability of the Monte Carlo simulation. We have repeated the entire simulation for each FM gate 30 times with all error terms the same. We have repeated these 30 simulations twice; once with the phase noise generated from our PSD turned enabled and once with the phase noise turned off. In the case of no phase noise the only stochastic error terms are the Gaussian distributed DC offsets in Rabi frequency and motional frequency. The results of these simulations can be seen in Figure 6.33. The result is that the effect of the noise PSD (during a single gate) is small compared to the Gaussian distributed DC offsets. This does make the variance of the data larger, fanning out at earlier gate times. However the general behavior is quite similar and it may be sufficient to model this type of noise as simple stochastic offsets compared to the gate time.

If we compute the mean parity error of these two datasets, with and without the phase noise PSD, we obtain 3.49% and 3.39% parity error per gate respectively. This means that without the inclusion of our PSD generated phase noise the guassian DC offsets account for 97.3% of our parity error. The standard deviation of parity errors, with and without the noise PSD, are 0.79% and 0.76% parity error per gate respectively. This confirms that the noise PSD has some effect of increasing the variance of the parity errors over these simulations, but this effect is still small, without the phase noise 95.79% of the parity variance is still accounted for. Under full simulation the most important factor is the DC initial frequency noise offset.
This is not necessarily surprising as the dominant noise peaks are 60/70 Hz whose periods are around 15 ms, while the gate times are 0.1-0.5 ms. The largest drifts will be seen as randomized due to experimental jitter on the timescales of multiple experiments.

The errors which account for 97.3% of our gate error include both the Gaussian distributed motional frequency drifts and the Gaussian distributed Rabi frequency drifts. Both of these low-frequency effects will cause parity error and they can be seen as equivalent as a constant calibration error experiment-to-experiment. The motional frequency drifts account for > 90% of the actual parity error, however to achieve good fitting to the particular shape of the data they are both required. The Rabi frequency drifts will act to increase parity error in a way which is independent of gate time (simulated alone it will uniformly increase parity error as a function of gate time). While the motional frequency drifts will provide the unique, gate time dependent, shape of the data as well as most of the parity errors, the inclusion of Rabi frequency error at this level will allow for the most accurate model. Simulations were run for a grid of values around what was physically reasonable to find the best fit.
Figure 6.33: A repetition of the simulation under the same conditions as Figure 6.32 30 times with and without the noise PSD. This plot is useful for visualizing the overall range that this sort of low frequency noise can take. It also shows that the most important contribution to errors is the stochastic drift experiment-to-experiment when the low frequency coherent noise is sampled.

It is interesting that we can use the noise PSD to model a large fraction of the parity error without considering any DC offset using the FF formalism (as in Figure 6.28). The slope of the parity error predicted by this FF model is not one, however the correlation is extremely strong. This tool can be reliably used to predict the performance of gates subject this this type of low frequency noise. In the $f < 1/\tau$ regime the spectrum of $F(\theta)(f)$ is flat and our motional coherence simulation of Figure 6.25 allows us to fit a noise amplitude and then the integration of this spectrum in the FF formalism will give a reasonable lower bound estimate of parity error. We do have some independent indication that our modes do experience coherent noise from our CPMG scans of Figure 6.24, similar collapse and returns indicative of coherent
noise are also seen on motional spin-echo experiments.

It is clear though that from the full simulation that the phase noise drawn from the noise PSD in reality is responsible for only 2.7% fraction of the total parity error per gate. This phase noise as simulated will never include an initial frequency detuning, and so we have needed to add the initial DC offsets, when then capture most of the problem when averaged over many experiments. The system is insensitive to the effects of motional mode phase noise in this The combination of the gate survey dataset and the Monte Carlo simulations provide a lot of food for thought when trying to understand the implications of low frequency coherent noise on ions. The relative insensitivity to the predicted displacement error experimentally is quite surprising, in this regime the behavior of the low frequency parts of $F_{\Theta}(f)$ is most predictive of gate quality. Furthermore to fully model our errors it we must consider both the coherent action of the noise within the gate and the sampling of this noise between experiments. This sampling loses phase coherence and we are left with an essentially stochastic process again. This stochastic Gaussian distributed motional frequency drift experiment-to-experiment is responsible for the majority of our error.
Chapter 7

Conclusions

Ion trap quantum computers based on surface traps provide a qubit platform with ultimate flexibility. Optically addressed multi-ion gates allow for arbitrary qubit network topologies. The ion trap qubits can be reconfigured as the application demands. Furthermore the ion’s place as a natural qubit means that the qubits are not subject to manufacturing irregularities. The ion however is fundamentally a charged particle, and so its motion will be highly sensitive to electric field noise. Surface electric field noise is thought to be involved in the anomalous heating experienced by surface traps. These heating rates have been shown to decrease sometimes dramatically at cryogenic temperatures\cite{LGA+08}. Silicon-based ion traps have had mixed results when cooled and seen only modest improvements. Here we have demonstrated exceptional common mode heating rates \(< 10\) quanta/s in silicon Sandia traps cooled to \(8\) K. This heating level is much suppressed from the average room temperate examples which can see levels ranging from \(100\)-\(1000\) quanta/s heating rates. There are indications that our heating rate may still be limited by technical noise rather than intrinsic surface noise.

A leading error source for two-qubit gates is motional decoherence\cite{WCF+20, SYBH22}. The motional decoherence process has been subject to intense study since the ability to manipulate motional states coherently has arisen\cite{TKL+00, MKT+00, TGD+16}. Here we report on the behavior of gates within a rare regime of motional errors. In this regime coherent and stochastic drift at the \(Hz\) level dominates over the effects of heating or the exponential dephasing time \(T_2\). There is some evidence
that the source of this coherent noise is well modelled by a low level electric hum and buzz. Ion trap experiments could benefit by applying the principles of audio engineering to mitigate this type of noise.

Cryogenic trapping provides many benefits over room temperature operation; however, it also presents unique challenges. Many researchers struggle with the vibrations inherent in cryogenic operation. We have been able to overcome the vibration problem by using a compact low-vibration optical cryostat integrated into an underlying structure which supports modular optical blocks. Having overcome the mechanical problems, we have described some unique electrical challenges presented by the Gifford-McMahon closed-cycle cryostat.

Even in the face of this slow, adiabatic, motional mode instability, we have been able to demonstrate frequency modulated Mølmer–Sørensen gates with an average of 98% fidelity over a population. This result was obtained without any triggering schemes to the mechanical pump. Furthermore we have demonstrated that the majority of our two-qubit gate error is not intrinsic but rather a matter of applying the principles of electromagnetic compatibility engineering at low frequencies. This means that given the other errors of our system, we should be able to achieve 99.5% or better gate fidelities once our grounding scheme is improved and RF amplitude stabilized.

We have demonstrated a compact cryogenic package approach to ion trapping. Our compact package is manufatureable and can be handled outside of a clean-room environment once packaged. All sample chamber modifications can be done on the optical table without disturbing optical alignment. The pressures in our package during operation allow for single ion dark lifetimes of more than 72 hours, and unlimited Doppler cooled chain lifetimes. Additionally the collision energies of cryogenic systems are below the threshold for chain reordering. Our optical design
approach allows for a modular system which can go years without fine adjustment. We were able to completely disassemble our experiment and in a single morning reassemble every optical block module in the experiment and be trapping ions by the afternoon.

All of these features mean that this platform is uniquely situated to host an error corrected logical qubit. The indefinite chain lifetimes and lack of reordering are crucial for logical qubits to be practical at scale. The manufacturability of our design means it is uniquely situated to being integrated into a modular approach to scale the quantum computer [MK13]. In the modular system-level approach many ion traps are duplicated and networked together by photonic interconnect. This approach to design could help forge the way to an era where researchers can focus more on quantum software and applications, and less on the hardware details. Abstraction and system integration were the keys that unlocked Moore’s law and enabled the classical computer revolution.

7.1 Future Prospects

7.1.1 Electrical Noise

In order to reduce the coherent electrical noise which couples onto our qubits we must rethink the immediate electrical design. Many design principles for reducing this sort of noise exist [Ott11]. It is unfortunate that we have an unavoidable grounded point very close to the AC motor which drives the cryostat. This grounded point may be difficult or impossible to remove. It may help to give the AC motor chassis a very low impedance path to the building ground, which should reduce the tendency of injected currents to travel through the trapping equipment; doing this may create
ground loop but it may not pass through the ion trap ground.

There is a desire to use this unavoidable ground inside the cryocooler as the ground star for trapping electronics and then slightly float the DC and RF inputs. In principle it does not matter if there is a mV level offset between the DC box output voltage and the DC voltage which the trap sees, or between the the RF signal (with respect to the helical resonator ground) and trap RF (with respect to trap ground). Some reduction in hum injected by the DAC box may come by grounding the shields strongly to the DAC box chassis but capacitively coupling them to the Montana system at the other end. This makes the shields a DC open with respect to the Montana chassis but a good shield at higher frequencies, preventing interchassis currents [Ott11]. This scheme will add a DC offset between the DAC box and the Montana system, but such an offset may be preferable to the hum noise injection. How stable this offset will be is an open question. Full isolation transformers can be used where possible and appropriate. For hum problems often the only practical solution is to attempt sensible remediations until the noise level becomes acceptable[Whi96]. If the method presented in Chapter 6 is valid then we should see some effects on the motional Ramsey scan.

Many of these interventions may not be necessary if we simply cut the ground loops which travel through the trap PCB. In UHV chambers the ground shields often terminate into the chamber itself at the feedthrough without problems. The injected hum/buzz from the DAC box may be at an acceptable level if we fix this ‘pin 1 problem’ [Bro05] which allows noise currents to propagate through the trap PCB ground plane. Attempting this should be the first step. There is no reason why the pick-off ground shield must terminate on the trap PCB at low frequencies.

Some benefit could come by attaching the helical resonator directly to the cold-finger. We can also then isolate the RF input ground shield by capacitive coupling.
This could make both the thermal connection and DC ground connection between very weak while preserving a good 50 Ω transmission line into the helical resonator at RF. This which would give the trap a very strong RF ground, even though there may be an offset between the input ground and the trap ground, this offset may be small. In this case, it is still not desirable to complete the helical resonator shield connection to the trap PCB, as this creates another loop. If the cold finger sample mount and the stage 2 cryocooler could be galvanically isolated from the Sumitomo cryocooler AC motor this approach could be even more powerful.

A more unconventional idea is to devise a method of differentially signalling the DC electrode signals, and route them to the trap in twisted-pairs. Twisted pairs are robust to magnetic field noise and differential signaling removes the hum/buzz problems, which probably exists to some degree in all ion trap systems wherein the DC lines are driven by unbalanced signals. How exactly a DC differential signaling scheme would work in an ion trap would while maintaining a coherent RF/DC ground remains to be seen, but the idea is worth considering. This would requires a major redesign of the DAC system, sample chamber, and trap PCB. Trap solutions under a differential signaling scheme may be able to be designed with symmetries in mind. Additionally some benefit may come from applying mu-metal magnetic shielding around the cryopump motor to reduce the magnetic noise which the experiment will see.

An even more ambitious system design change which may greatly reduce the 60 Hz noise contribution could be totally isolating the ion trap system from the building grounds. A DC power distribution system could be made, wherein 120 V power is only supplied to a bank of linear power supplies through isolation transformers. The DC power bank could be physically isolated form the system and the DC power delivered in a bundle to the setup. Core trapping electronics such as the DAC box, DDS box,
RF amplifier, would only draw their power from this DC bank. Floating equipment grounds in this way has hazards and must be done with the utmost caution.

7.1.2 Ion Imaging and SPAM Error

We have acceptable SPAM performance for ion ion at the 0.5% level. During but our large MS gate survey it became clear that drift in detection and pumping calibration is still a major problem for this system. While the fiber couplings of our CW upstream plate remain stable for years without adjustment we must open the box to adjust the waveplates frequently and also the polarization for our free-space 370 nm beam also must be adjusted throughout the day. The source of this drift must be tracked down. A polarization filter at the output of our free-space fibers could bring benefits. Another contribution to SPAM errors is our imperfect lens alignment. Ideally having full 6 degree control of the lens in-situ would help, but in practice accomplishing this in such a small space is difficult. Furthermore adding two more degrees of freedom just adds more equipment which could break. The best approach would be to redesign the tapped-hole matrix on the sample mount to feature dowel pins for alignment. In this way the lens stage could be removed and relocated reliably during trap installation. The lens stage could then be designed to be very close to the optimal position, and any time adjusting the X/Y micrometers for fine optimization would not be wasted upon the next trap installation.

7.1.3 Longer Ion Chains

In this work we have trapped and worked with long chains for many hours in Chapter 5, but we have not shown results of gates on longer chains of ions. We have spent some time working on gates in 5 ion chains but decided to scale back and study
the nature of our gate errors first. Working with ion chains was cumbersome with only two individual addressing beams, we had to invest much more time to sideband cooling in sequence, this made us even more susceptible to trap frequency drift. The system has been ungraded to feature electromagnetically-induced-transparency (EIT) cooling [FTD+20] which will allow us to prepare an entire chain to close to the ground state, only needing to make fine adjustments with sideband cooling afterwards on the modes nearest to our gate. If we can reduce our electrical noise problem noise problem by 20 dB this EIT cooling will allow for rapid progress to work with long chains and quantum algorithms. Furthermore as demonstrated in Chapter 6 the use of f-robust FM gates allows us to tolerate the level of coherent noise currently in our system. If we can reduce the stochastic DC drift by a factor of 2 by improving the RF amplitude stabilization lock then with the f-robust FM gate we would be comfortably into the 99% gate regime, according to simulation. There are other modulated gate schemes currently in development that could be even more robust against gate area errors. Modulated control schemes can allow us to tolerate this electrical noise if we reduce the slower drifts.
Appendices
Appendix A

Trap Cryopackage Assembly Recipe

Figure A.1: The packaged Sandia Peregrine trap depicted in (a) is assembled with the lid depicted (b) which has been prepared with carbon getter satchel, the ytterbium target, the side and imaging windows, and indium for thermal and mechanical bond. These two pieces are simply joined together by hand in a clean room to form the trap cryopackage shown in (c). The processes to make the lid are the bulk of the procedure and the actual packaging together is very simple.

List of Parts and Tools

Here we list in more detail the assembly process for a Peregrine cryopackage. A mechanical drawing for the lid is included in Appendix B.

1. Machined Copper lid, we fabricate the lid out of Copper 101 which is 99.99% copper and has low oxygen content. Oxygen impurities can impede thermal conductance at low temperatures.

2. 100-pin CPGA packaged peregrine trap

3. Indium wire of diameter 0.03” or indium preform fitted to the lid groove
4. Indium sheets of 5 mil thickness

5. 1 1/4-20 × 3/16” set screw, this is for plugging the carbon getter chamber

6. 1 8-32 × 1/8” set screw, this is for plugging the ablation oven chamber

7. Finest pitched copper mesh which can be obtained an example being McMaster-Carr item #9224T8 with pore size: 0.006”, open area: 30%, wire dia: 0.0045”

8. A fine aluminum mesh, preferably even more delicate than the copper mesh.

9. Carbon getter packages pre-activated, often sold in paper satchels.

10. Ytterbium metal source (usually sourced as a 1 mm square shaped oblong log) cut into a piece which can fit snugly into a 0.0315” diameter hole without slipping, ideally force is required to jam the piece into the hole. Natural ytterbium is preferable as the isotope selective loading gives you variety. Some protocols may require multiple isotopes.

11. Epotek T7110 epoxy (epoxy and activator)

12. Electronic scale with 0.01 mg accuracy is required for measuring the epoxy and activator ratio by mass

13. Wooden toothpicks

14. Plastic mixing dishes for epoxy

15. Disposable syringes for weighing epoxy

16. (5 for HOA cryopackage 6 for Peregrine cryopackge) N-BK7 AR coated windows fitted to lid side windows loosely. We use 5 mm diameter 2 mm thick windows which can be obtained with a few different anti-reflectionv (AR) coatings from Edmund Optics part no. 13-286
17. 15 mm diameter by 2 mm thick imaging window in Fused Silica or N-BK7, we have used Edmund Optics part no. 13-334 successfully which can come with a variety of AR coatings

18. Acetone and IPA

19. Clean room towels

20. Clean aluminum foil

**Cryogenic Package Assembly**

1. Prepare the ytterbium source

   - Ytterbium is commonly sold in logs with rectangular cross section of approximately 1 mm on a side.
   - A freshly cleaned razor blade is sufficient to cut the ytterbium log.
   - Place the ytterbium log onto a metal optical table or other such metallic working surface.
   - By holding the razor blade at a 45° angle with respect to the log axis a small triangular shaped piece can be cut.
   - Cutting the piece as small as possible, at this angle, ensures that there will be some segment of the ytterbium metal sample with a small enough dimension to fit into the 0.0315” diameter hole snugly.

2. Clean the Metal Parts

   - Collect the metal parts including lid, set screw, copper mesh, and ytterbium source.
   - Sonicate these parts in acetone for 10 minutes, dispose of acetone.
• Sonicate these parts in IPA for 10 minutes, dispose of IPA.

• Leave parts to air dry in clean room on top of clean room towel.

• The parts can be baked in a vacuum oven at approximately 200-300°C. A small aluminium foil dish can be formed to keep the smaller parts from blowing around during vacuum pumping and venting.

• Backflow nitrogen into the oven and allow the parts to return to room temperature before exposing to lab air. Care must be taken to minimize the oxide which will build up onto the machined copper part. This process will be enhanced at high temperatures in oxygen.

3. Lid window attachment

• Once the lid and its parts have been cleaned you can attach the lid windows and cure the epoxy.

• This step is best completed inside a fume hood as the epoxy emits toxic fumes before fully cured.

• Open the 5 side windows and the imaging window and place them onto a clean-room towel for easy access along with the cleaned lid, opening the optics only before you are ready for attachment.

• Prepare the working space with epoxy (T7110), activator, syringes, toothpicks, plastic dishes for weighing epoxy, a milligram (jeweler's scale) or micro-gram lab scale. You need to wear gloves and be as clean as possible.

• In the plastic dish you can use syringes to weigh the epoxy and activator in a 10:1 mass ratio epoxy:activator. The epoxy is a viscous opaque liquid and the activator is a thin clear liquid. The best method is to first weigh the epoxy into the dish (bulk of the mass), then tare the scale, and use
a syringe to drop the activator into the dish until 0.1*(epoxy mass) has been added.

- Remove the dish from the scale and mix with toothpicks until completely homogeneous.

- The side windows can be pushed into their holes and will tend to stay as long as not jostled around too vigorously. The windows are again, 5 mm diam. by 2 mm thick sliding into an approximately 5.1-5.2 mm diameter 2.03 mm deep hole and they easily remain in place due to the aspect ratio. Furthermore the epoxy can be applied more generously to the side windows and the surface tension and adhesion will tend to keep the windows in place.

- It is possible and even advantageous to cure all of the windows simultaneously. In order make sure the side windows stay in place during transit we can cut 3 small aluminum foil strips to wrap around the entire lid and keep opposing windows in place. They will rest underneath the lid and can all be joined at top to keep the windows in place while moving it towards the oven. Prepare some aluminum foil strips wider than the windows and long enough to wrap around the entire lid.

- For each side window you can dip a toothpick into the epoxy to obtain a single drop. Then apply it to every window indentation by swiping once around the corners and perimeter. Take another drop and spread along the flat surfaces which the windows will press up against. Take care that the sides and the flat have a thin film of epoxy. Do not apply so much epoxy that there will be a risk of dripping or flow onto the window surface. Be conservative especially on that flats. For the small side windows we
depend more on the corners and edges to bond. Apply epoxy to every window opening. For application you may want to pick up the entire lid and examine your work closely in light.

- Once the epoxy has been applied to every window opening then place the lid on the work surface on top of the aluminum foil strips which have been arranged into an X-shape to cover the windows at opposite angles, with one strip down the middle to cover the two center side windows.

- Shove the windows into each window opening firmly and without much fidgeting. Make sure the windows are seated down entirely.

- Lastly we leave the lid down and use a toothpick to apply epoxy for the top imaging window. Be careful to not get any epoxy into the corners or side-walls. Paint the epoxy conservatively on that flat surfaces making sure that aside from the pumping trench it forms a seal around the perimeter. Here it is important not to use too much, then simply drop the imaging window into place and center it with respect to the opening. Take care not to block the pumping trench with epoxy.

- Wrap the aluminum foil side protectors inward onto each other and crumple together. Try not to touch the top window.

- Bring the assembly to the vacuum oven and place inside.

- It is important to control the cure time precisely. We opt to use an 80C cure for 2 hours as the recommended cure schedule. For this step it is best to pump out, backfill with nitrogen, then bring the oven up to temperature. Once up to temperature you can then pull vacuum. If you immediately pull vacuum it can take a long time for it to come to 80C and it will be hard to control the baking temperature and time precisely. It is important
to supervise the oven at this step. It has been found difficult to get a steady 80C cure for only two hours without any manual supervision of the temperature loop in the vacuum oven we have available. Sometimes it overshoots and the heater must be manually turned off. It is suspected that curing at too high of a temperature, or for too long of a time can change the glass transition temperature, or the sheer hardness in ways that are unfavorable for the window integrity. Others have had trouble getting fused silica windows to survive cooldown cycles without breaking. There is a limited dataset but the three things suspected to contribute to window breakage are: the use of T7109 rather than T7110, which has a larger and uncharacterized sheer hardness, the use of an excess of epoxy, especially on the sidewalls, and lastly setting the oven temperature and walking away. This may not give the proper temperature curve over the duration of the cure. Perhaps overshooting for a significant fraction of time, after taking too long to heat up. We have observed two cases of a fused silica window attached with T7110, and conservative amounts of epoxy, cured at 80C for two hours, to survive many cooldown cycles without any cracking.

4. Prepare the carbon getter satchel

- It is best to prepare the carbon getter only after the rest of the lid has been prepared such that it can be taken directly from the vacuum oven to the lid for assembly and installation. It is best to minimize its exposure to lab air after baking.

- Activated carbon getter can be purchased in large canisters filled with smaller pouches of pre-activated carbon.

- Remove one carbon pouch and bring to clean room.
• Use a razor blade to cut several strips of copper mesh which is to be no more than 1 cm wide by approximately 6 cm long, the exact dimensions can be varied a little bit for each strip in width and length.

• It is best to make 3-6 of these at one time and then pick the best one. There is some trial and error involved in getting the fit correct.

• The strip of copper mesh is to be folded on a line going down the long dimension, and approximately 0.3 cm (or 1/3) of the short dimension to make a V-shape.

• Once you have formed a few of these long V-folded copper mesh strips you can cut into the carbon getter pouch and obtain a small pile of the carbon granules.

• Use a tweezer to select the largest, most robust carbon granules. Use care in selecting them to not crush them with the tweezer.

• Only 2-4 large carbon granules are needed. It is better to have a few robust large ones than a big pile of small flakes. The reason it that it is very easy for a crushed carbon flake to contaminate the trapping chamber with carbon dust, care should be taken during the rolling to minimize crushing of your 2-4 carbon grains.

• You do not need much, you just need it to be well contained and well thermalized by the copper mesh matrix. Even just a few carbon grains have vastly more surface area than the entire interior of the machined copper lid. No need to put huge amounts of carbon in there. It just becomes harder to contain the grains, and harder to cool the mass of insulator. The aim is for a small well-contained satchel with approximately 20% carbon by volume surrounded by a a few layers of copper mesh on all
sides. Once installed the set screw will inevitably crush the carbon a bit when securing the satchel into the lid. You want to contain the broken pieces. An additional benefit of only working with a few large grains of carbon is that they are easier to handle during rolling.

- Place the grains of carbon into the bottom of the V-shape at one end of the copper strip, and then fold over that end (so that you are folding along a line orthogonal to the first fold). Take this fold and begin to roll down the length of the strip until you have a roll of carbon, neatly contained on the bottom by the two folds, with a diameter which will fit into the getter chamber of the lid. The height of this roll (set by the short width of the copper strip) should be such that there’s not a lot of extra material after rolling, and the total height can fit into the getter chamber with room for the set-screw to engage with the threads.

- This roll has an open end. To close the end get another small strip of copper mesh which is wider than the opening and can wrap one time around the entire copper satchel. The wrapping direction is around an axis orthogonal to the previous rolling. The goal is to fully seal the open end. Squeeze the entire thing into the most neat package possible.

- Size test this against the copper lid, and continue making them adjusting the size of the copper strips until a good fit is achieved. This can be difficult and take 5-6 tries.

- Post-select only the best few of them, place in an aluminum foil basket inside the vacuum oven.

- Bake the getter satchels inside the vacuum oven set to 350°C overnight.

- The getter can be left inside the oven until we are ready for the final
assembly.

- Backfill with nitrogen and let come to room temp before opening the vacuum oven.

5. Final lid assembly

- For this step a circular cut aluminum mesh which fits at the bottom of the getter chamber is required. This is another layer of protection from the carbon debris. Place this at the bottom of the getter chamber first.

- Remove the getter satchels from the oven, place the best and most well-fitted example into the getter chamber of the lid. Screw down with the 1/4-20 setscrew until the setscrew is flush with the outer surface.

- Place the ytterbium chunk into the ablation oven hole. The ideal situation is that it is triangularly shaped such that a large fraction of it can fit into the hole, requiring pressure to securely fit. Some fraction of it is wider than the hole so that it cannot fall through. In this way you can guarantee that some of it will have a line-of-sight to the ablation laser. It will remain securely placed, but it will not fall out into the main chamber of the cryopackage during handling. The worst case scenario is that either it falls out during construction and you have no ytterbium source.

- It is good to take up space behind the ytterbium target so that the setscrew applies pressure to the back of it. A good way to do this is to make a small roll or ball out of the copper mesh which goes in behind the ytterbium target. It must be large enough such that it takes force to fully screw the setscrew in behind it, which then will fill the voids behind the target.
- Screw in the 8-32 set screw until flush with the exterior.

- Past designs have used epoxy to seal these set screws. The performance of these systems were not the best. The problem with this approach is that you will have to do an epoxy cure in the presence of the carbon getter. This can contaminate the getter. There is no need for a hermetic seal at the set screw threads because the package is already leaky. The threads have no line-of-sight thus gas particles are unlikely to make it inside. Even if they did the particles will get lost in either the getter or the copper meshes. Fantastic and chain lifetimes have been achieved with no epoxy used to seal the set screws.

- At this point the lid assembly is complete.

6. Creating the trap cryopackage

- A surface trap must be packaged with compatible ringframe CPGA for the lid you have created. An example is seen in Figure A.1(a)

- A clean room is needed to perform this step safely. This step requires indium sheets (or ringframe preforms) which have been cleaned by IPA and allowed to dry.

- Cut four indium sheets to fit at the four sides between the CPGA and the lid. Optionally have available a length of indium wire sized for the ringframe, or an indium preform which fits into the groove.

- With the lid upside down place the indium sheets into the lid. An example of this is shown in Figure A.1(b).

- Open the trap package and remove from its carrier.
• Ensure you are orienting the lid in space correctly. The relationship between the RF pin and the ablation oven depends upon your system PCB and laser geometry. Get clear on this before the next step.

• Lift the lid and trap up to eye height keeping the lid upside down with respect to gravity, carefully mate the two pieces taking care to not damage the wire-bonds with the lid, or to knock the indium sheets out of place. It is best to practice this step with a dummy trap/package if available and become confident and quick at this step.

• Once the lid is seated firmly on the package to create the trap cryopackge it can be flipped over such that the lid is on top and the pins are on bottom.

• Apply light pressure to make sure the lid and trap are fully engaged.

• It is possible here to use a room temperature curing epoxy to add an extra mechanical bond between the two pieces. If the ringframe groove has been properly tolerated it should be a tight slip fit and friction should help it stay in place during normal delicate handling. The design of the trap-to-cold finger sample mount, in concert with the trap PCB puts pressure between the lid and the package, especially after cooling, and so once installed the system is stable. The preference of this researcher is to avoid elevating the cryopackage temperature again at all or introducing any epoxy cure steps after the lid has been placed onto the trap package. Also it is important to minimize the time exposed to lab air between opening the trap and pulling vacuum on the cryostat.

• At no point is this system baked with the trap. The temperature is never elevated once received from Sandia.
• Cryopackage assembly, before installation into the PCB. This is a time where the windows can be inspected, at a glancing angle, in good lighting to look for smears and blemishes. The windows can be cleaned with lens tissue and spectroscopically clean IPA/Methanol.

• The completed cryopackage is seen in Figure A.1(c)

7. Final PCB installation

• In this system the trap is installed into the PCB in the clean room immediately after the lid has been applied.

• We make sure that the RF pin on the trap cryopackage is aligned with the RF pin on the trap PCB.

• We place the trap into the CPGA socket.

• Two soft plastic handled screwdrivers are needed to apply pressure here. The handles are cleaned thoroughly before this step. Any other firm but forgiving polymer material which is narrow enough to engage with CPGA corners, but with a wide area for gripping is fine.

• This step, to be done safely, requires two workers, one whose job is to hold the lid down while also holding the trap PCB up off the table slightly so that stress is not transmitted to the PCB components and connectors during the application of pressure. A well designed jig could replace this person.

• One person makes sure to keep the lid in place, and also to hold the entire PCB thus they become a safe, soft, intelligent conduit for the applied forces.

• The other person uses the two screwdrivers (the plastic handle ends) to
apply force to the CPGA corners. They apply force on opposite corners, then the same corners, then the corners individually, then back to opposite corners. The other person watches to gauge the progress and direct the pressure application events. We want to make sure that the package is completely seated into the socket up to the ‘hats’ of the pins.

- Once the cryopackage is seated entirely into the CPGA socket we then switch to applying pressure to the lid itself rather than the CPGA. We apply an entire 120 lbs+ body-weight of force onto the top of the lid, with the second person still holding their fingers underneath the PCB to absorb this force onto a soft surface rather than between the PCB components and the metal table. We are trying to enhance the adhesion between the indium sheets and the copper/ceramic surfaces using pressure. Too much pressure could break the ceramic. Adding heat to melt the indium has resulted in poor vacuum performance anecdotally.

- It is still important to handle carefully as the lid could get dislodged. As much as possible it is good practice to keep the entire lid cryopackage flat with respect to gravity until the last moment before installation where it is turned 90° about the center with the getter and ytterbium targets moving down with respect to gravity into the final position. Avoiding jostling the assembly is good practice. As constructed it should be robust, but in the rare case that the ytterbium has come dislodged, or some carbon debris is loose, avoiding shaking and turning the entire assembly in all directions will mitigate the impacts of these events.
Appendix B

Design Drawings

B.1 Peregrine Cryopackage Lid Design

Here we present the Peregrine cryopackage lid design. The tolerance of the ractrack-shaped ringframe groove need to be tightened slightly, this design was too loose.

Figure B.1: Peregrine Cryopackage Design Drawing View 1
Figure B.2: Peregrine Cryopackage Design Drawing View 2

Figure B.3: Peregrine Cryopackage Design Drawing View 3
B.2 Cold Finger Sample Mount Design

Here we present the design of the cold finger sample mount for reference.

Figure B.4: Cold Finger Sample Mount Design Drawing View 1
Figure B.5: Cold Finger Sample Mount Design Drawing View 2
Appendix C

First Generation Compact Cryogenic Ion Trap Platform

C.1 The Trap Cryopackage

In this Appendix we explore the motivations, design philosophy, design, and experiments performed on a first generation prototype of the compact cryogenic ion trap platform, although ions were never trapped in this system it was used to thoroughly study isotope selective ablation loading, to prototype our experimental design philosophy, and to test a directly driven lumped element RF tank circuit which used a low output impedance amplifier. While ions were never trapped in this system we learned from it a lot about what it would take to make the concept work and what elements needed to be paired back for the future design.

C.1.1 Trap Cryopackage Design and Assembly

The ions require UHV pressures, the surface ion trap with its DC and RF electrodes, and a ytterbium source. Nothing else is needed inside the UHV space provided that DC/RF can be fed through to the trap, and that there is adequate optical access to the ions. We sought to design an ion trap package which allows for handling and only includes these minimally required functions. This initial design was conceived as a fully vacuum sealed trap cryopackage.

In Figure C.1(b) we see an image of the fully assembled trap cryopackage and in Figure C.2 we see a schematic representation, in Figure C.1(a) we see the modified ceramic pin-grid array (CPGA) package with trap attached. The cryopackage is
made by bonding the top ‘lid’ to the modified CPGA package. We use the CPGA package as the DC/RF vacuum feed-through once bonded to the lid. A standard 100-pin CPGA was sent to HI-REL to be modified for cryogenic packaging. During modification a square shaped (with rounded edges) gold-coated Kovar ‘ringframe’ was located and then brazed to the CPGA’s ground plane. This ringframe is $0.096 \pm 0.002$ in tall and $0.03 \pm 0.002$ inch in width and will be used to mate with a trench shaped groove inside the lid to the trap cryopackage. A gasket made of indium wire was used between the lid and the ringframe extrusion on the CPGA. By mating the package lid to the modified CPGA package and applying combination of heating and force a good vacuum seal can be created. Indium has a long history of use in cryogenic vacuum seals[108].

![Figure C.1](image)

**Figure C.1:** (a) a view of the modified CPGA package with ringframe attached and the EPICS trap installed. The EPICS trap sits atop a spacer which raises the trap height above the ringframe and uses vias to route the bone pads to a taller height. (b) A picture of the complete assembled package we can see the copper mesh encapsulated carbon getter as well as the epoxy attached viewports.

In Figure C.2 we can see that the lid is machined with all previously mentioned features which are required in the UHV trapping volume. An ablation oven for
the storage of our ytterbium source is fashioned as a simple cylindrical hole with a 0.1272 in (3.23 mm) depth and 0.05 in (1.27 mm) diameter. While optimizing the ablation oven in this first iteration a cover was added serving as a Faraday cage which hoped to reduce the number of charged species exiting the ablation oven. An upper chamber containing activated carbon getter (wrapped tightly in copper mesh) is included. This chamber for getter is positioned within a cone-shaped geometric feature meant as an electromagnetic shield between the dielectric window and the ion location. Each window is affixed to the titanium housing using EPOTEK 7110 cryogenic epoxy. The assembly procedure is included in this chapter. This cryogenic lid originally was entirely fabricated out of gold coated Kovar but this material as abandoned for Titanium for its non-magnetic properties.

In this first generation of experiments we used two simple single layer gold on fused silica surface traps. Both were designed originally for cavity QED experiments
and thus are fabricated on top of a dielectric mirror. They were chosen to test this concept because of their simple integration to our height requirements and low trap capacitance (compared to silicon-based Sandia traps) which makes the use of lumped element RF circuits more forgiving. These traps have three round mirror trapping locations and the trap was designed for cavity QED experiments. For the early stages of this project we used a trap fabricated at Duke by Van Rynbach et al. [VRMK16]. In later stages we began using the EPICS trap fabricated at Sandia National Labs and thus the entire CPGA stack (including filter capacitors and a trap spacer (for raising the height of the trap to accommodate our work) was assembled and verified there on top of a modified CPGA package.

The general procedure for sealing the package was to first prepare the titanium lid, by gluing and curing the windows, baking out the activated carbon getter and encapsulating it in a fine copper mesh. The carbon getter is a messy material which can create particulate contamination and the copper mesh serves to keep it contained. Care must be taken that the copper mesh package is well sealed. Small particles of carbon can migrate to the trap during handling and damage the delicate electrodes, and disturb the potential. The ablation material is then loaded into the ablation oven cavity. The entire assembly is assured to be clean with a solvent cleaning and low temperature bake.

Once the lid is prepared it is loaded into a ‘puck’. A puck has grooves for handling inside vacuum using manipulation forks. It has a cavity for containing either the lid or the CPGA in a stable way. These two pucks as seen in Figures C.3(a) and C.3(b). In the case of the lid it is retained, and aligned to the CPGA using four screws which hold it in from the backside. The CPGA is held into the puck via two small aluminium pieces pinning it down, which are in turn screwed into the puck itself. The CPGA is stabilized first and then the lid is aligned to it so that the ringframe
Figure C.3: (a) Image of the surface ion trap (EPICS trap) fastened into the indium crushing puck. Screws and small aluminium bracer hold it down at two edges, male dowel pins protrude for alignment purposes. (b) Image of the bottom puck which contains the lid assembly, as well as female dowel pin receptacles for alignment purposes and male stands for holding the trap aloft during the baking stage. The lid is secured into this puck with four 4-40 screws (c) A view of both pucks assembled in the final orientation for crushing. Both dowel pins are engaged and within the vacuum chamber a vertical crushing force will be applied pressing the two pucks and thus the CPGA and lid together with an indium gasket causing the seal.

mates with the indentation on the lid (between which the indium is crushed). Before the final crushing (in the final design) first the trap is held aloft above the lid for the baking step, so as to provide venting space, with an additional set of pins which act as stands.

The indium seal is performed inside a large multi-purpose vacuum chamber with a trolley system which was first built as a wafer bonding system in [McK11]. In Figure C.4 (a) and (b) we see an image and schematic of the entire chamber. The trolley system is used to lead the puck stack into a chamber at the load lock. We then ferry the parts to the wafer bonding chamber for heating and crushing. The water bonding chamber where the seal if performed is shown in Figure C.4. Heating to nearly a temperature of 150° C is required to soften the indium enough for sealing. This chamber also has a vertically adjustable puck holder which can be moved into strong contact with a preheated crushing surface. By screwing the bottom puck
Figure C.4: (a) An image of the trap cryopackage vacuum sealing station. We utilized a wafer bonding chamber first built by [McK11] on our campus. (b) We load the trap cryopackage pucks into the load lock and ferry them to the wafer bonding station. (c) Within the wafer bonding station there is a heater (for approaching the indium softening point, and a pressing mechanism for completing the seal. This figure was reproduced from the thesis of Kyle McKay © 2011

holder one can move the pucks until they make contact with the top, then one can continue to add force between the top heated crushing element and the rest of the stack. One wants to avoid completely melting the indium ideally, and the internal temperature of the package itself is difficult to measure well in this particular setup, as
the temperature sensor is tied to the heating element. So one gradually increases the current through the heating element until the crushing is accomplished successfully. The package is then removed from the system after letting it cool.

This entire manufacturing procedure (baking and crushing) can be completed in two directions: with the lid on bottom and trap on top (normal to the trap surface points down with gravity), or with the trap on bottom and lid on top. Doing it with the lid on the bottom is somewhat more mechanically difficult with our vacuum feed-through manipulators however is lessens the risk that indium (if it melts) will spill into the trap bond pads and cause electrical shorts. It also lessens the likelihood of debris landing on the trap during the process.

C.2 Lid Assembly Procedure

C.2.1 List of Parts and Tools

1. Titanium lid
2. 100-pin CPGA packaged trap
3. Indium wire of diameter 0.03”
4. Puck system (for crushing)
5. Ytterbium metal source
6. Epotek 7110 epoxy (epoxy and activator)
7. 4 N-BK7 AR coated windows
C.2.2 Cryogenic Package Assembly Outline

1. Clean and bake constituent parts
   - Wash in Acetone
   - Sonicate in 2-Propanol
   - Load into oven (vacuum lab) and high temperature bake 300-400°C

2. Prepare carbon getter material in copper satchel which involves rolling a few grains of carbon tightly in fine carbon mesh. Ensure that at least two layers of the mesh surround the carbon on all sides.
   - Load onto molybdenum puck and bake in UHV assembly chamber at 800°C
   - Move carbon getter satchel to the parking stage for storage

3. Load Ytterbium into lid securely inside clean room
   - Solvent clean assembly (Acetone wash, sonicate in 2-Propanol)
   - Load into vacuum oven and to high temperature bake overnight (greater than epoxy cure temp)
   - Remove lid from vacuum furnace

4. Epoxy all windows, excluding the cryo pumping window.
   - First thoroughly clean windows and lid with solvent wash and baking in sample prep
   - Using Epotex T7110 mixed by mass in correct 1:10 by mass ratio (activator to epoxy), apply epoxy using syringes to a dish on a miligram scale
• Cure each window separately: Applying all windows at once makes windows sag creates a weak seal due to gravity acting on them. So we apply the window (with window facing up) at 150 C in sample prep oven for 15-20 minutes (after activation 200 C is new temperature limit)

• Solvent wash the lid assembly in 2-Propanol (acetone degrades T7110)

5. Load indium gasket

• Clean indium wire with acetone of diameter 0.03”

• Push into the sealing groove with small overlap of indium after shaping into groove

• Use the oven in the range 150-170 C until flow to increase attachment to the lid

• Clean assembly with 2-Propanol

• Warning: if you do not attach the indium to the lid with heat then the gasket could fall out during manipulation inside the vacuum sealing chamber.

6. Load getter satchel into the lid, epoxy last window 1

• Partially cure for 15-30 minutes in dry argon-filled (or Nitrogen filled) bag until tacky

• Fit screws of two-puck assembly for indium crushing, fixing lid assembly into top puck

1Its tricky epoxying the last window while minimizing the contamination of the getter material in principle we could let the epoxy vacuum cure with the lid resting window-side up (like the other windows were) to avoid the window falling out of alignment, however this could let chunks of carbon fall down into the trapping volume if the satchel has leaks. This would also lead to another step of loading and unloading the chamber before finally fitting the lid to the crushing-puck assembly
• Cure lid assembly at 150 C for 15 minutes in vacuum furnace

• Bake at 120 C for 1 hour

• Leave to out-gas overnight

7. Load the ion trap/CPGA-puck assembly into transfer chamber

8. Flip the ion trap/CPGA-puck assembly and carefully lower onto lid-puck assembly

9. Move the entire assembly up into heating chamber

10. Bake at 120 C for an hour

11. Let assembly out-gas overnight

12. Heat vacuum furnace to 170 C

13. Apply pressure for full vacuum sealing of package

14. Remove assembly from chamber system

15. Remove cryo package from puck assembly
C.3 Cryostat Design

We chose to utilize a low-vibration cryostat specially designed for optical access by Montana Instruments, this system is called the Montana Cryostation (CR-197). It has a quoted temperature range of 3-350 K with less than 10 mK temperature stability. Our system had a large number of electrical feedthroughs and in practice had a base temperature of 4.71 K. This system has a quoted 5 nm pk-pk vibration...
amplitude. The cooling power at 4K is [GET NUMBER]. Unlike most conventional closed cycle cryostats which must hang from above with the sample chamber dangling, this system is able create a thermal transfer joint which isolates the sample space from vibrations as well as makes the sample chamber available directly on the optical table. In Figure C.5(a) we see a labelled model of the system we used with the vacuum chamber portion open where the 40K shield can be seen, the system features 100 pins of DC feed-through made available through two 50 pin DC Mini-D Ribbon (MDR) connectors. Four RF feedthroughs are provided using SMA connectors on the outside and MMCX connectors in the inside. In Figure C.5b we see the 40K shield is lifted to reveal the central sample chamber which has a tall cold finger sample mount pedestal, at the top of this pedestal is our sample thermometer which sits below the trap PCB. The trap cryopackage can be installed into a 100-pin CPGA socket into the sample chamber. The viewports in the vacuum chamber and 40K shield are made of AR coated fused silica.

In Figure C.6 we see a real image of the sample chamber with labels, the sample chamber was co-designed with Montana Instruments. There is a bottom 40K PCB and an upper 4K trap PCB. The 40K PCB and the routing beneath is was done by Montana instruments, the 4K PCB and the routing above that was build at Duke and subject to design iterations throughout the trapping experiments. The MDR connector vacuum feedthroughs are soldered to a PCB on the other side which is mounted inside a small box shaped module which has o-ring seals to rest of the sample vacuum chamber, underneath the 40K PCB these wires are routed to the 40K PCB and exposed as header connections. Inside the main sample chamber all of the DC cables are brought between these 40K and 4K PCBs with varnished wires twisted into bundles and terminated on male header connectors. They are clamped to the 4K sample mount as a thermal anchor and then mated with female header

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Figure C.6: The real sample chamber is shown: varnished DC wires are routed in twisted bundles to the trap PCB they are clamped to the cold finger sample mount to temperature lag them at 4K simple header connectors are used for connecting to the PCB. The trap PCB is made of G10 epoxy glass and is screwed down at four corners into the cold finger. The semirigid RF cables can be seen going from the bottom 40K PCB to the top 4K PCB, these can also be temperature lagged to the cold finger using a scheme explained in the text. The trapped cryopackage can be seen as installed into the 100-pin CPGA socket. A thick ground wire can be run from one corner of the PCB (under the screw) to the top of the lid. A toroidal RF inductor can be seen behind the cryopackage.

connectors on the trap PCB.

The RF feedthroughs are SMA on the outside which have immediate MMCX connectors in the inside. Between the feedthrough and the 40K PCB there are semirigid Pasternack RF cables with stainless steel outer conductors for lower thermal conductance. Between the 40K PCB and the 4K PCB we have used Pasternack PE-020SR
which have copper outer conductors. These cables were temperature lagged (not shown here in any figure) before the connecting with the PCB by clamping a thick copper wire with a thin electrically insulating coating to the cold finger (under the DC cable clamps). This wire acted as a thermal bridge which was brought as close to the cable termination at the 4K PCB as possible (maximizing the temperature gradient between 40K and 4K) and then the thick wire was bonded to the semirigid cable by weaving a knot of very thin coated copper wire, which was then varnished lightly. Making sure the RF input and pick-off cables were properly heatsunk was essential to our lumped element RF circuit operation.

C.4 Optomechanical Design Philosophy: Optical Blocks

Most atomic physics experiments are complex assemblies of off-the-shelf parts such as optical posts, clamps, cages assemblies, mirror mounts, and modulator mounts. These parts must be placed by hand and aligned to the central experimental region. The optical posts and camps serve solely to be able to locate in space a variety of optical components (whose sizes may differ) in such a way that the optical axis can be aligned between these components, thus this is a highly flexible platform for creating complex experiments. However there are shortcomings to this optical design paradigm both practical and technical.

Thermal and mechanical noise in the lab environment over long periods of time tends to make all of the components creep, and furthermore stacks of such components have a more severe tendency to creep. In practice the experimentalist must expend large amounts of time both assembling the setup and maintaining alignment, this becomes a long-term maintenance problem and becomes a large fraction of the man-hours expended. Provided one has settled on a design for the experiment it makes
sense to begin removing degrees of freedom in favor of ease-of-assembly and structural stability.

A technical concern for ion trapping experiments is that individually addressed Raman operations are sensitive to mechanical noise which reduce interferometric stability of the beams and will lead to lower fidelity gate operations. The effect of mechanical noise on beam stability is worsened by optical components being located upon tall posts. Air turbulence also affects the mechanical stability of the beams and this effect gets worse as a function of optical path length, and the degree to which the optical path is open to thermally active devices. In practice it is difficult to reduce the path length greatly when using the conventional approach to optical design. As you shrink an optical system built conventionally (and shrink the optics) and the path length becomes both dominated and limited by the posts, camps, and mounts.

The ultimate in stability is a monolithically designed metal piece that underlies the experiment and contains direct mounts for all system optics, however this is often impractical, undesirable, and unmanufacturable. A compromise that balances both of these approaches is an underlying structure that can be dowel pin referenced between the center of the experiment (in this case ions) and the optical components. Special optical breadboards can be made which a matrix of tightly tolerated dowel pins as well as tapped holes for fastening. If these breadboards can then be located accurately with respect to the chamber then and each unit of optical functionality can be modularized and and separated into a single smaller blocks dubbed ’optical blocks’ then we have a setup which can be freely modified but machining smaller customized plates while having all the stability benefits of a monolithic construction. Each optical block has dowel pin alignment on the bottom, to mate it the the breadboard, and then all optics are aligned on top with dowel pins as well. Today there are a wide variety of 1/2 inch kinematic mounts available with dowel pin alignment fea-
tures included for industrial integration. Certain specialized static mounts for gluing optical components can also be designed and placed with the dowel pin alignment matrix. More implementation details for designing using this approach are given in Chapter 4 where the second generation system is discussed.

C.4.1 Optomechanical and Optical Design

Figure C.7: An image of the PIVOT ion trapping system with optomechanics. The Montana cryostat sits atom a base plate located by custom dowel pins which mate with 1/4-20 counterbores. Two custom bread boards can be mounted at movable heights.

The entire cold head/sample chamber module was then bolted to a large aluminium platform, which serves as the basis for the construction of all final optical components. These optical components are built into a tower of aluminium plates, supported by four steel shafts. Each aluminium plate is custom machined to sup-
port aluminium block sub-components. Each aluminium plate is held aloft by shaft clamps for vertical support. For horizontal placement and locking wide ellipsoidal pieces with rounded edges and a counter-sunk thru hole were used as clamps. They mate with the aluminium plates through large side clearance holes (which have 8/32 tapped holes) are used to screw aluminium locking pieces into contact with the shafts. The aluminium locking pieces are rounded with the radius of curvature of the shafts at the interior. The entire coldhead/sample chamber is located to the optical setup via dowel pins to which the shafts are referenced in the bottom aluminium plate. The cryostat sample chamber module has through holes in the back designed for 1/4-20 screws, but which we commandeer for dowel pins so that we can be sure of the cryostat locations in the design. This entire cold head setup, with all final stage optics is secured to a large aluminum breadboard, which is then placed atop vibration dampening foam, to isolate the cryopumping vibrations from other experiments on the optical table which may be negatively effected, such as the 369.5 nm locked laser, and a local cavity experiment. This micro-cavity experiment was able to operate less than 6 feet away while the cryostat was running and yet still lock the micro-cavity \cite{VRSS+17}.

The entire assembly as designed in CAD is seen in Figure C.7, the helium compressor which attaches to the cold head/sample space via a 10 foot long hose is not pictured. This hose must be placed such that it undergoes an 180° turn between the sample chamber and the compressor. From this point all light sources would have already been prepared and fiber coupled. Thus the general function of this set-up is to combine beams in free space and direct them through the final focusing lens. On the optical level there was a lot of extra room provided for possible future use and the capability of bringing in counterpropagating beams was included.

In Figure C.8 we see a section view of the pivot system optics. We have de-
Figure C.8: A section view of the PIVOT System with the second half of the imaging optics excluded, the lens was designed to have an \( \text{NA} = 0.35 \) with low aberration. The axes of the CW/ablation beams are shown in relation to the trap axes.

signed a 0.35 NA imaging lens using 2 inch off-the-shelf Thorlabs optics to image ion fluorescence with a 5.5x magnification. Two positive meniscus lenses are used
(power split) to begin to bring the rays back to center and then the rays are shaped until a plano-concave lens is reached which creates an infinite conjugate image. The plano-concave lens has the most stringent positional tolerance requirements. Z-focus Adjustable 2 inch lens mounts which would fit our space requirements could not be found so a series of drop-in spacers were custom machined which could touch each lens at a point allowing for tight positioning. The geometry could be set up in CAD and then the thickness of the spacers adjusted until the edge-to-edge distances match our Zemax design. All of the lenses and spacers are dropped into a tube which was tightly slip-fitted and has a lip at the bottom. At the top of the stack the lens tube becomes threaded and a lid can screw into the assembly loading down the stack. At the top of the lens stack is a turning mirror which does not aperture the beam, this is screwed down into the top lens cover. The entire lens is held in a slip-fitted cradle which mounts to a three axis stage (shown also in Figure C.10). The lens can be smoothly rotated in its cradle to aim the ion fluorescence but there is enough friction that gravity hold it in place once positioned.

The axis of the trap as well as the relevant beams (ablation and CW) are shown in Figure C.8. The view of this sectioned CAD model is the direction that the ablation laser would point, the ablation ovens can be seen in the cryopackage. An f=100 mm focal length achromatic doublet is used to focus the CW beams and its position is adjust with an Optosigma TSDS-255SL stainless steel XYZ-stage which was and is the most compact high stability steel three axis stage available and we have used it in all of our compact designs.

In Figure C.9 we see the central optical breadboard and an the CW optical block, there is lots of space on this breadboard for future expansions, but is is already apparent how compact the free-space optics can become. We originally chose to combine all of our blue light (369.5 nm / 399 nm / 935 nm ) lasers into a single
Figure C.9: We bring in all blue beams combined in a single fiber which can be accomplished by dichorics upstream. The 935 nm light was combined in free-space in the setup. There is a 3 inch by 3 inch pitched matrix of dowel holes which can mate with the optical blocks, there is also a 1 inch by 1 inch pitch of 1/4-20 counterbores which allow the optical blocks to be screwed in from the bottom (the blocks have 1/4-20 tapped holes on the bottom). A f=100 mm achromatic doublet focuses our beams to the ions, there is 3-axis platform stage on the other side which can be used for bringing in counterpropaagaing beams.

optical fiber. This can be done with off-the-shelf Semrock dichorics, however later on it was discovered that the 369.5 nm and 399 nm light had a very different collimation condition in the Microlaser collimators that use used. This threw the 399 nm beam off in z we later added some 1/2 inch off-the-shelf optics to bring in a separate 399 nm / 391 nm beam, this is not dipicted in this document. This issue of collimation can be fixed by using reflective parabolic collimators which are now available in compact form. We have included a separate 3-axis platform stage in the designs for
the eventual integration of a counterpropagating beam should the need arise.

The optical block is mounted upon a custom breadboard which has a 1 inch by 1 inch pitched matrix of 1/4-20 counterbores which are on the bottom of the plate looking up, and interspersed with this matrix is a 3 inch by 3 inch pitched dowel pin matrix for repeatable mounting of blocks. Each bock has locating dowel holes in which pins are inserted. There are also some 1/4-20 tapped holes on the bottom of the blocks. The blocks can then be screwed in from the bottom of the plate. The advantage of this scheme compared to traditional optical breadboards (which are matrices of tapped holes) is that on top of the optical block no space needs to be reserved for counter-bores used in mounting, so the entire space of the block is free for optomechanics.

The imaging system is in the top level and can be seen in Figure C.10 the entire imaging system can fit on top of a single Newport 3-axis brick stage. The entire imaging system can be translated together in space. It includes a magnetic mounted pellicle which can be mounted to direct light into either of two opposing directions. One side directs fluorescence to another steering mirror and then through a reimaging lens with 2x magnification, this throws image far enough away from the setup to mounts a large Andor EMCCD camera for more sensitive ion imaging. The other leg of the imaging optics can send light to a Guppie CMOS camera with a small pixel size (2.2 µm) but less sensitivity. The pellicle allows about half the light to directly enter a PMT state detection, the pellicle can be removed when no camera imaging is needed. As previously discusses the lens is mounted on a 3-axis stage which holds a cradle for the lens to rest in.
Figure C.10: (Left) A schematic diagram of the imaging system, the ion fluorescence exit the lens and is turned 90° towards a block mounted on top of a Newport brick stage. A 1/2 inch 45/55 pellicle can be mounted to direct the fluorescence either to another steering mirror and then to a reimaging lens on a z-stage, this sends the image with 2x magnification to a large off-setup Andor EMCCD. With the pellicle directed the other way we can send the ion image to a Guppie CMOS camera with small pixel size (2.2 µm. The pellicle will allow half the counts into the PMT, which has a 369.5 nm notch filter and an adjustable iris. (Right) another view of the imaging system from a different perspective where we can see how the imaging lens is mounted to a 3-axis stage. The lens sits in a slip-fit cradle with enough friction that it is firmly placed due to gravity, yet can be rotated by hand.

C.5 Lumped Element Directly Driven RF Tank Circuit

Trapping ytterbium ions stably in a surface trap requires at least 150-200 V amplitude of RF voltage at frequencies of 30-50MHz with higher frequencies in that range generally leading to more trap stability. Traditionally most ion trap setups use the helical resonator technology detailed in [SSWH12]. This is an impedance matching
transformer which can be matched to the ion trap load via the adjustment of the relative positions of the two coils. The main drawbacks of this technology are its large size and mechanical instability.

We worked towards a more integrated solution a simple lumped element RLC resonator. A superconducting inductor can be used and lumped elements can be used for the impedance match, high Q-factors have been obtained\cite{GNK+12}. This can work well however the impedance matching elements must be placed close to the circuit and so be located inside the cryostat. Then the impedance matching procedure becomes tedious because the system must be temperature cycled iteratively. Once a good matching condition is found the actual voltage gain obtained from this more complicated circuit can be less than the Q \cite{GNK+12}.

We attempted a design where the circuit was driven with a low impedance source. Rather we employ amplifiers with a low output impedance and high output current to drive our circuit and do not attempt to impedance match the system. A schematic of this circuit can be seen in Figure C.11 where we use the OPA2674 a DSL line driver for this application. The OPA2674 amplifier has less than an ohm of output impedance at 45 MHz. A great improvement from the first amplifier which was tested, the ADA4870, which had a 3 Ω output impedance at 30 MHz, and an unworkable 7 Ω output impedance at 45 MHz. Higher frequencies between 40 and 50 MHz are needed to ensure trap stability. At these higher frequencies the ADA4870 was unworkable.

The operation of the circuit depends critically on the reduction of resistive impedance between the source and the capacitive ion trap load. The compact incarnation of the circuit included an amplifier board external to the cryostat with a 6 inch long coaxial cable leading into the cryostat sample space. The sample space had a short semi-rigid coaxial cable going from room temp to the 40K shield. Another longer cable leads from the 40K shield to the 4K circuit board which can be seen in Figure C.12.
Figure C.11: RLC resonator circuit driven by a low output impedance operational amplifier.

Originally all of these semi-rigid cables had stainless steel cores to reduce the thermal conductance from room temperature. We opted to replace them with copper cores. This however minimized the series resistance in the circuit and lead to larger voltage gains.

Figure C.12: The in-cryo printed circuit board which both routes the 100 DC lines as well as contains the RF circuit and capacitive pick-off for measurement purposes. The inductor was made out of NbTi superconductor from SUPERCON (model 56S53) with a diameter of 0.33 mm when insulated.
The inductor used as a NbTi wire from SUPERCON (model 56S53). It was wound on a Teflon core into a toroidal shape to be about 1 $\mu$H and held in place using VGE Varnish. The advantage of the toroidal shape is the reduction of fringe fields which can interfere with other parts of the experiment.

In order to understand and analyze the circuit certain analytical expressions must be solved for. Using Figure C.11 as a reference the circuit, first we reduce the measurement circuit. The AC resistance from larger pick-off capacitor must be small compared to the measurement resistance or the act of measurement will load the circuit. Having met this condition one can neglect the measurement resistor in our analysis and the voltage at the trap will be given by:

$$V_{trap} = \frac{1}{jR_s\omega C_T - \omega^2 LC_T + 1} V_{amp}$$

We then define two parameters, the resonant frequency $\omega_0 \equiv \frac{1}{\sqrt{LC_T}}$ and the Q-factor, $Q \equiv \frac{\omega_0 L}{R_s}$, we can then substitute these values giving:

$$V_{trap} = \frac{1}{1 - \frac{\omega^2}{\omega_0^2} + j\frac{\omega}{\omega_0 Q}} V_{amp}$$

When the driving frequency approaches this resonant frequency $\omega = \omega_0$ then we have:

$$V_{trap} = -jQV_{amp}$$

Thus the voltage gain factor present across trap capacitance is simply the Q. The Q can be calculated from this resonance curve as: $\omega_0/(\Delta\omega)$ where $\Delta\omega$ is the full width at half maximum of the power or the spectral width between the max($\frac{V_{trap}}{\sqrt{2}}$)
C.5.1 Testing the Compact RF circuit

One can see the general dynamics of the amplifier with driving by looking at Figure C.13(a) which shows the decay of the quality factor as a function of driving amplitude applied to the amplifier. The Q of our circuit was calculated by finding the spectral width point on the resonance curve where the voltage drops to $\frac{1}{\sqrt{2}}$ of its maximum and dividing the resonant frequency by that spectral width. We measure the trap voltage by using the capacitive pick-off shown in Figure C.11. In these figures we call the voltage level driven to the amplifier the pump amplitude.

![Figure C.13](attachment:figure_c_13.png)

**Figure C.13:** (a) The Q as a function of amplifier input voltage which we call the pump amplitude. One can see a high Q values for low pump amplitudes which degrade as a function of drive amplitude. The base temperature of the cryostat also increases as a function of the driving amplitude from 4.79K to 4.97K. This temperature is measured at the cold finger and thus does not represent the full temperature swing of the coil itself, which must be greater due to the large change in Q. (b) We see the voltage estimated by capacitive pick-off measurement as a function of driving voltage.

There are two regions of Q-factor decay with a plateau in the middle. The initial quality factor was almost 100 but it rapidly drops to 75 by the time there is a 50 mV pump amplitude. Once the pump amplitude is above 150 mV there is another linear regime of Q-factor decrease. This seems to be the effect of the amplifier saturation.
for large driving amplitudes. The saturated amplifier heats up the cryostat increasing the series resistance by heating up the conductors on the inside. The cryostat temperature increases as a function of the driving amplitude from 4.79K to 4.97K. This temperature is measured at the cold finger and thus does not represent the full temperature swing of the coil itself which is greater. The saturated amplifier may also experience increased output impedance itself as its hits its current limit.

**C.6 Isotope Selective Ablation Loading and Time-of-flight Spectroscopy**

In most room temperature ion trapping systems an atomic beam is created via resistive heating of a metal tube. The metal tube is filled with the species to be loaded, and due to the increase in the rate of sublimation, a diffuse beam of neutral atoms are produced. These atoms can be ionized with lasers and then loaded into the ion trap. This method, albeit reliable and time-tested is difficult to adapt to the trap cryopackage. This is due to the need to pass several amps of current through the tube thus dumping heat into the 5K environment. This limitation can be mitigated by placing the oven behind the 40K shield and using clever methods to send the neutral atomic flux through a hole in the shield such as was designed for[VWB⁺13a]. However this solution is not very compact and is not compatible with the space limitations present in the Montana Instruments Cryostation.

We opted to use laser ablation to create our atomic beam. Metal ablation with pulsed lasers has a history of use in laser etching systems, film deposition systems, and general study [CPH90, TCG⁺00, TTS⁺02, VWR⁺98, KRN⁺89, ACFH04, TPA⁺04]. In the past ion traps have been loaded using Nd:YAG and \( N_2 \) pulsed lasers[LCL⁺07, ZOHSP12, HGH⁺07]. Usually these systems have directly loaded ions (lacking photoionization laser) into deep traps. The problem with this method is that there is no
isotope selectivity unless one uses a isotopically pure source. On our system we seek
to load ytterbium into a shallow surface trap while doing so in an isotopically selective
way. We also want to create the most efficient loading process possible so that the
amount of ytterbium sputtered onto trap electrodes is negligible. As such the direct
loading of ions from an ablated plasma is incompatible with the trap cryopackage
and we must optimize ablation in the weak regime and increase our photo-ionization
efficiency as much as possible.

The ablation process is characterized by a threshold [TPA+04] of onset which is
measured in fluence. The fluence being the energy per unit area in the pulse. The
slope of the ablation yield (in mass per pulse) and threshold (fluence intercept of this
yield curve) is uniquely determined by the material to be ablated and the wavelength
of incident light.

The ablation process can be messy and generate nano-particles [BBC+98], neutral
atoms, ions of various charge states, and small clusters [LJK+03]. It is of interest
to shield out these unwanted charged species, and to load ablated neutral atoms
to allow for isotope selectivity. Additionally high ablation powers yield significantly
more flux than is needed for trapping as well as macroscopically large chunks of metal
to the trap surface. This is not a problem when using macroscopic ion traps, however
the delicate nature of the surface ion trap necessitates a careful consideration of the
appropriate ablation conditions. The plume dynamics of ablated metals have been
well studied for common manufacturing metals such as gold, aluminium, and steel.
There is a specific lack of a study for ytterbium ablation, including measurements
of its ablation threshold and the collection of information on the dynamics of the
ablation plume in vacuum. Due to the widespread use of ytterbium in atomic clocks
and ion trap quantum computers, as well as the system integration advantages of
laser ablation loading, it would serve the community well to have, on record, careful
measurements of the ytterbium ablation process.

The BigSky ULTRA Nd:YAG laser system with 1064 nm output was used for ablation. A half wave plate and polarizing beam splitter is used to adjust the power level. A generic red laser is co-propagated with the original 1064 nm beam in order to serve as a targeting laser. The 1064 nm beam has a collimated output waist of 2 mm, it is focused through a 30 cm lens to produce a beam waist which (along the shorter axis) is 50 µm and along the long axis is closer to 200 µm in diameter. Laser on it was found out that the mode shape of this 1064 nm was severely distorted due to laser degradation, when filtered 80% of the power was removed, a beam profiler not not available to diagnose the exact problems with the focused beam. In a close-by experiment they used this same laser, but with a filtered and thus Gaussian mode to measure the ablation threshold [VAS+19], the fluences in this data have been re-scaled by our ablation thresholds, specifically the energy where the ablation enters the linear regime. Exactly determining the correct factor from a beam profile or the amount of power filtered out via pin-hole is extremely complicated, but the ablation thresholds provide a metric by which the incorrect fluences (in the original data) values can be re-calibrated, and the data is worth presenting because a large range of powers is sampled, and a more fine scan of the low-power regime is performed in time-of-flight spectroscopy. This beam is brought into the experiment in a direction mutually orthogonal to the cooling laser and imaging path (as seen in Figure C.2). A video camera with a long working distance objective is used to check for the alignment of the laser to the target.

A micro-controller (Arduino) and differential amplifier based PID locking system was created which grabs data from the network output of the wave-meter and uses it to feed back the piezo position on the Toptica 399 nm laser. The lock bandwidth is limited only by the rate at which the wave-meter updates, which is in turn power
dependent. No inherent arduino bandwidth limits have been reached because this wave-meter update is, itself, much slower and on the order of 15 Hz when optimized. This is sufficient to keep the frequency within the range of the isotope line-widths (40 MHz), and has fluctuations below 5 MHz.

C.6.1 Natural Ytterbium Isotope Spectroscopy

The most basic goal is to identify of the center of the $^{174}$Yb peak. Despite the fact that we use Yb 171 as the hyperfine qubit, the trapping scheme for Yb 174 is more easily implemented and serves to test the basic functions of the ion trap. Using the PID lock we can scan the 399 nm laser frequency through the entire span of natural ytterbium isotopes we can compare this spectrum to the literature values. This enables us to determine the extent of Doppler shift due to deviations from orthogonality between the photoionization beam and the atomic beam. From experiment to experiment this is subject to change and in fact in our experiment a change of angle of this photo-ionization beam of $2.3^\circ$ produces an approximate 100 MHz shift. Figure C.14 shows the observed spectrum of Natural ytterbium.

The data in Figure C.14 was taken by triggering the Andor camera to begin a 0.5 second exposure before triggering a series of 6 pulse sequences. We measure the total counts in a region of interest near the trapping location. In each pulse sequence the 1064 nm laser was triggered, the control system waited a specific time (92 $\mu$s) measured to be the delay between laser trigger and the arrival of the optical pulse. The 399 nm light was triggered on for 200$\mu$s to capture fluorescence from the entire velocity distribution. All six pulse sequences occur within a single frame of the Andor to aggregate the signal. The lock point of the 399 nm laser was moved between discrete intervals for each iteration of the experiment. The scan intervals
Figure C.14: The PID lock was used at varying frequencies to aggregate data over a wide range. This reveals the many different (power and velocity broadened) isotope peaks.

are more closely spaced (10 MHz) near to peaks, and more far spaced (max 50 MHz) in large flat periods, with an average spacing between collection points of 25 MHz.

These spectroscopy lines are broadened by laser power and the velocity distribution. More narrow spectroscopy lines could be obtained if the power of the beam was adjusted to be below saturation, as well as the region of interest in the Andor beam reduced to exclude fluorescence from off-axis atoms.

Ablation Oven Geometry and Plasma Shielding

Ablation flux will travel normal to the surface of the ablated material independent of the angle that the ablation laser makes with the surface. A beam with a large angle to the surface will be impacted insofar as the fluence is reduced due to the beam being spread over a larger area, however the angular flux distribution will not be impacted. Ablation creates a broad angular flux distribution without using the geometry of the ablation oven to limit it.
Figure C.15: (a) An image of the ablation spot test package, we have taken the Duke cavity trap as installed and wire-bonded (a broken example) and installed a glass slide in order behind it in order to aggregate a ytterbium film to both ensure flux (and an upper bound of ablation intensity) and to get a rough measurement of the angular flux distribution from our oven. The 1064 nm ablation laser is sent through the glass slide and over the trap towards the ablation target and then the ablation flux is allowed to travel back over the trap and deposit onto the glass slide. (b) An image of the sample from behind after the test, the contrast has been artificially enhanced to make the ytterbium film more easily visible and a dotted line at the interface of the film has been added. The two salient features noted from this experiment was that our 1.27 mm diameter ablation oven does a poor job of collimating the flux, and that the ablation flux has a large angular distribution, depositing everywhere not directly shadowed by the trap substrate. Furthermore no deposition is allowed to aggregate in the film where the 1064 incident and reflected beams are hitting the glass slide and apparent holes in the film are seen.

In order to help shield out charged species, the large ablation oven (2mm diameter opening) opening was covered with a titanium mask with a smaller 250 µm opening. The grounded metal being nearby the plume of charged species would hopefully act as a Faraday cage screening the electric fields produced by the plasma discharge, and also attract the charged ions to the grounded metal shielding most of them out. Furthermore by providing a smaller aperture between the ablation site and the
trapping location are can improve the collation of the flux.

**Figure C.16:** An image of the cover for shielding charged particles. The backside has an extruded feature which mates with the larger ablation oven hole.

**Figure C.17:** The angle between the ablated flux and the 399 nm spectroscopy probe beam determines the extent to which the flux is subject to the doppler effect.

One can investigate the presence of detectable spectroscopy signal at both 369.5 nm and 399 nm in the ablation plume to determine their relative abundance as a function of ablation power. This was accomplished by measuring the energy with the
ablation pulses over a wide range of powers and then using this information, along with the known beam area, and the fluorescence counts to calibrate the intensity of spectroscopy. This allows the measurement of both the ablation threshold for our system as well as the onset of detectable 369.5 nm counts indicating the presence of ions.

![Graph showing spectroscopy counts as a function of power for 370 nm and 399 nm light.](image)

**Figure C.18**: Spectroscopy counts as a function of power for 370 nm and 399 nm light, this reveals the neutral ablation threshold as well as the onset of general plasma discharge at the trapping location.

Shown in Figure C.18 are the results of this experiment. One can see that there are highly insignificant counts at 369.5 nm until very large ablation fluences are reached, and at this point can can see that the received counts from 399 nm begin to become noisy and go down on average. The most interesting thing to note about these 369.5 nm counts is that they are not even dependant upon the presence of the 369.5 nm beam and do not show up on the screen as an isolated beam of counts but rather
a global increase in brightness with a uniform distribution throughout the frame. These counts also only begin to appear at very large ablation powers such that the visible plume on the camera seems to fill the chamber. This is indicative to such a strong ablation that there is a plasma discharge all the way at the trapping location rather than simple a plume of travelling ions/neutrals. This level of ablation power is very destructive and not desired, none-the-less it is quite interesting. No above-background 369.5 nm counts are observed for a very wide region of powers. This indicates that the small opening used to shield out charged species may be effective. Repeating a systematic investigation of this wide power range scan for both spectral components in the ablation oven lacking a cover may yield an interesting plot to compare it with. This has not been performed yet because warming up the cryostat is done rarely to preserve the integrity of the package and performing trapping attempts in this new package was the main concern.

One can see in Figure C.19(a) the effect of over-ablation on a trap surface, in this picture of the after-effects of this study. Large amounts of fluorescence spots accumulate on the trap surface. They light up when sending the 369.5 nm laser through the system. This sort of deposition also contributes to the electronic destruction of the trap. Electrodes get shorted together or to the ground and this makes trapping impossible. Just for reference Figure C.19(a) shows the trap with the 369.5 nm laser but no ablation laser. Figure C.19(b) shows the trap under the same 369.5 nm conditions with the ablation laser present at large enough powers to display fluorescence under the 369 nm filter. Comparing Figures C.19(b) and C.19(c) we can see that the latter (without the 369.5 nm laser shining through) still displays this fluorescence around that frequency. This is only the case for high ablation powers.
Figure C.19: (a) An example of the trap (with no illumination) after lots of ablation spectroscopy at a mix of powers. What should appear mostly black is now illuminated by the presence of the 370 beam. In this image there is no ablation laser being shot. One can clearly see the presence of deposition on the trap. This is especially apparent on one side of the trapping location. The EPICS trap contains 11 µm tall wall (compared to an ion height of 50 um) for shielding ablated flux from the mirror surface at the trapping location. This wall has taken much of the deposited flux and it appears as a splattering of scattered 370 photons when viewing the trap. (b) This is the same situation but with the ablation laser running at these higher powers in which we observe 396.5 nm fluorescence as in Figure C.18. The 369.5 nm laser is still on in this image. (c) This is the same situation but with no 369.5 nm laser running at all. One can clearly see that still we get 369.5 nm photons through out filtering system from the ablation plume alone and no excitation laser. These ablation powers are clearly destructive.

C.6.2 Time-of-flight Neutral Ytterbium Spectroscopy

Surface ion traps have a low trap depth compared to the macroscopic traps which have been used in ablation loaded experiments in the past. The ablation by nanosecond pulse lasers causes what has been called a 'phase-explosion' [GTS09]. And there is a sharp threshold in fluence which, when crossed results in a rapid onset of strong ablation which causes the creation of a high-temperature plasma [CITE]. The larger the energy of the ablation event, then the higher the temperature of the initial explosive cloud becomes, which makes the average ion and atom velocity higher. This causes a smaller population of trappable atoms and ions. We sought to study the
dynamics this ablation event using pulsed spectroscopy from our neutral ytterbium laser to answer the following questions: at what fluence is the ablation threshold, how does the balance between neutral and ionized ytterbium scale as a function of fluence, and what is the velocity distribution of the ablated cloud. Once understanding these three questions we could then estimate the probability of trapping an ion per ablation pulse which we seek to be on the order of unity.

An acousto-optic modulator (AOM) was employed to act as a fast switch. We collect fluorescence with an Andor EMCCD camera. The camera has a low data collection bandwidth, however the fast fast switching of the 399 nm pulses from the AOM can be used to gain fine time resolution. Time-of-flight spectroscopy was used to probe the velocity distribution of the ablated flux. Additionally Doppler spectroscopy was performed, as was a full frequency scan of the neutral ytterbium lines in order identify the exact frequency of the desired isotope, ytterbium 174.

Using the timed 399 nm and 1064 nm pulse delivery system in sync with the Andor trigger the velocity distribution could be extracted from the ablation plume. This is important information if one is to calculate the number of neutral ytterbium atoms which travel through the trapping region. Surface ion traps have a lower maximum to the amplitude of voltage that can be applied to the RF rails. This is because of the small feature size inherent in chip fabrication. This lower applied feature size means that in practice the depth of the pseudo-potential well is low. This is particularly the case in our system due to the limits of our RF circuit.

From an investigation of the trapping solution the trap depth dictates that we should only be able to trap ions with a velocity of 310 m/s or less. This corresponds to a trap depth of a slightly less than 90 meV. The distance from the ablation source to the trapping location is about 1 cm and as a result the approximate time it takes for an atom at this velocity to traverse the source-trap distance is about 32 µs.
The first step was to see how this velocity distribution scaled with ablation pulse energy.

**Figure C.20**: Measured velocity distributions using the time-of-flight analysis for a variety of fluences.

In Figure C.20 we see a plot of this time of flight distribution for a variety of powers. The data was collected by capturing 6 repetitions of the same pulse sequence all contained within one 0.5 second exposure of the Andor camera. A single 1064 nm pulse was followed by a delay time and then a 2 µs 399 nm pulse. The delay time included 92 µs delay between the electronic 1064 nm trigger and the actual optical pulse plus an additional delay time to probe the atomic flux with different arrival times at the trapping location. These 2 µs long pulses were scanned in 2µs increments until there was no discernible difference from the background counts.

One notable feature of interest in Figure C.20 is that the peak speed of the distri-
bution increases with increasing power before seeming to saturate both in fluorescence counts and in speed. The behavior of this peak in arrival time is plotted in Figure C.21 where a linear relation is found until the last point where the maximum saturates. If we calculate the centroid in arrival time of each of these time-of-flight distributions we can arrive at an average ablated atomic velocity. This value scales linearly with fluence in this regime as seen in Figure C.22, even though the peaks seems to saturate the mean velocity of the ablated atomic cloud continues to increase with fluence. Thus the exact nature of the velocity distribution depends upon pulse energy. Also noted is that there are no counts for the lowest pulse energy in Figure C.20 as this energy is also below the measured threshold for our system from Figure C.20 while in the next higher energy (right around threshold) we see the turn-on of significant counts. This lowest energy trace is excluded from Figures C.21 and
C.22. We estimate the number of trappable atoms per pulse using a method similar to [VAS+19] which has been adapted to use a small area of interest on the Andor camera rather than a PMT. We estimate for the trappable atoms per pulse near the ablation threshold (measured at 0.049 J/cm$^2$) to be 0.35 atoms/pulse (with our low RF voltage potential well).

C.7 Thermal Design Flaws and Postmortem

This system helped improve our understanding of the RF circuit and of the dynamics of ytterbium ablation, but we never demonstrated a trapped ion in this system. In this section we will explore some reasons why the experiment failed and what lessons were applied to our next generation design.
Vacuum seal problems

Anecdotally we had noticed some problems with the vacuum integrity of the packages. One ion trap package that we were using was tested by conducting some temperature dependant 399 nm spectroscopy. If our vacuum sealing process worked, thus preserving a rough vacuum within the package, then we should still be able to observe neutral spectroscopy, even in a warm cryostat. This was tested by watching the 399 nm fluorescence as we warmed up. The spectroscopy signal is lost between 150 and 225 K before being completely absent at room temperature (Shown in Figure C.23). If one looks at the phase diagram of water, it seems likely that water at low pressures would melt in this temperature range, most of the common gasses would have already evaporated, furthermore the cryostat was still under vacuum (although not cooling) at this point and so the presence of these gasses, given a leak, would be suppressed. It is unknown when the vacuum seal broke, this could have come about through repeated temperature cycles, or be a fundamental problem in the sealing process. The two points of failure could be the epoxy seal used to attach the windows. Water trapped in the carbon getter, or possibly some failure in our method to make indium seals.

This problem was further confirmed by attempting warm spectroscopy in a newly assembled package. This package was assembled and cooled down only once. Trapping attempts and low intensity spectroscopy experiments were performed over the course of 2 Months. Upon warming it was found that the spectroscopy signal went away. In order to check the severity of the leak the system was pumped out with the roughing pump to a vacuum. After 10 minutes the spectroscopy signal returned with just rough pumping. Upon venting the spectroscopy signal went away once more. This is indicative of a gross leak in the cryopackage. This leak could very well have
Figure C.23: Spectroscopy signal at 399 nm was a function of cryostat base temperature for a particular package.

existed during the transit of the package from the sealing chamber to the cryostat. Such a leak would have contaminated the system with water. This is shows that our current assembly procedure has some deep flaw, either in the indium seal, or in the epoxy window seals.

C.7.1 Thermal Design

The vacuum seal problems alone, even with a gross leak, should not preclude trapping, cryopumping should still make the pressures in the package low, and a small leak may cause a differential pumping effect which still creates UHV pressures. The biggest central oversight in the design of this system is the thermal design. The cooling of the package, due to the mechanical design, must be done through the CPGA pins which are temperature lagged onto the cold-finger. Ideally the coldest part of the package should be the titanium lid, and critically, the activated carbon getter capsule. This means that ideally the package should be thermally grounded to the cold finger.
with large surface area. The DC lines and PCB should also be tied there, as they necessarily attach to the lid. However the current design ensures that the RF heating dumps into the CPGA, which is itself tied to the cold finger, however the lid, which we depend on for cryopumping is cut off between this heating element and the cold finger. A temperature sensor was added to the top of the lid, without RF applied temperature measurements that its temperature was 8K. The base temperature of the cold finger sample mount by contrast was 5.10 K (for this measurement). Measurements of the temperature increase of the lid during the RF operation at high voltage showed large spikes in the lid temperature, all the way up to 25K (when the base temp had also risen to 7.4K). This result was from a more primitive version of the circuit, which had poor thermal grounding, nevertheless it demonstrates a large temperature difference between the lid and sample mount during RF driving. This result makes sense because the main heat load during RF driving is the dissipation in the trap itself.

C.7.2 RF Voltage

Lots of optimization has led to RF amplitudes on the low end of what is needed to trap. The RF voltage amplitude at 45 MHz being around 120 V. The deeper our potential well the larger population of ions available to trap from the velocity distribution. With this low voltage the trap depth is limited to 90 meV but depths of 200-300 meV are achievable in surface electrode ion traps. There are many possible improvements to the RF circuit, however most of them require more space within the chamber. In the next system a compact helical resonator will be used to rule out the RF circuit as a variable.
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Biography

Robert F. Spivey III goes by ‘Tripp’, his given nickname. After graduating high school Tripp did summer internships at Applied Quantum Technologies (AQT) in 2010 and 2011. While at AQT he worked on programming an automated micro-lens array characterization tool. Tripp attended Rensselaer Polytechnic institute (RPI) where he majored in Physics. While at RPI he did undergraduate research primarily with Professor Toh-Ming Lu, whose lab is primarily an applied solid state physics lab which works on creating novel materials for advanced electronics and energy applications. Prof. Toh-Ming Lu has also long had an interest in the dynamic scaling of rough surfaces and the Monte Carlo simulation of thin film growth. Tripp developed a discrete Monte Carlo simulation of pit-type defect evolution during EUV multilayer deposition and published on it in 2013 (RF Spivey, R Teki, TM Lu, Thin Solid Films, 2013). Tripp then developed a new model of atmospheric and high pressure deposition processes. Tripp graduated before this work bore fruit but trained his successor on his simulation who pushed the project forward and published (T Merkh, RF Spivey, TM Lu, Scientific Reports, 2016). Tripp also worked with Professor Gyorgy Korniss who studies the physics of social networks (T Jia, RF Spivey, B Szymanski, G Korniss, PloS one, 2015).

Tripp graduated Summa Cum Laude from RPI in 2014 and decided to matriculate to Duke University to join Professor Jungsang Kim’s MIST lab and pursue his doctorate. While at Duke Tripp has mainly worked on advancing the cryogenic package project. Tripp is excited by the unique challenges presented by the field of quantum control, engineering, and computing. Tripp is particularly interested in advancing the quality, manufacturability, and scalability of the hardware platform.