Chapter 1

Introduction

Chapter 2

Ion Trap Apparatus

A vast effort is spent on the initial build-up of the an ion trap system, but throughout the life of the experiment, a greater effort is spent on its daily maintenance. I hope that this chapter will serve as a resource for future members of the FastGates team, as well as provide a useful recipe for anyone building a similar system.

Due to the size and complexity of the system, we introduce an inital overview of the design, motivated by the desired functions. As the name suggests, an ion trap experiment aims to confine arrays of single ions, this is achieved by static and dynamic electric fields which, due to the ions possesing non-zero electric charge, can provide trapping potentials, section 2.2. Due to the fragility of the internal states of the ion (these are state of the art sensors after all), we must take great care in isolating the ion from any noisy environment. This neccesitates the use of ultra-high vacuum (UHV) systems, section 2.3, vibration isolation, and magnetic shielding, section 2.4. To manipulate the internal electronic states of the ion, we create local electric and magnetic fields using RF antennae and, in this work, lasers, sections 2.5 and 2.5.3. Finally, to interface with the apparatus we have built, at the time scales set by our interaction strengths, we require a sophisticated and custom control system which is discussed in section 2.6.

2.1 System Design

2.2 The Ion Trap

2.2.1 Trap RF Chain

2.2.2 Trap DC Voltages

2.3 Beam Geometries and Vacuum System

2.3.1 Vacuum System

2.3.2 Atomic Source Oven

2.3.3 Optical Access

Beam Paths

In Vacuum Prisms

Dual High NA Objectives

2.3.4 Imaging System

2.4 Magnetic Field

2.5 Laser systems

2.5.1 Photoionisation

2.5.2 Ca^+ Laser Systems

2.5.3 Narrow Line Width 729 Laser

2.5.4 Single Addressing System

2.6 Sinara Hardware and Artiq

Chapter 3

Experiment Characterisation

Before we can dive into running novel experiments involving the motion and spin of the atoms, we need to characterise our apparatus. This allows us to both benchmark our system against state of the art results, and to reveal any current limitations of the apparatus which we may need to address.

3.1 Available Transitions

3.2 Spin

3.2.1 Rabi and Ramsey Scans

Here we briefly describe the method in which we extract Rabi frequencies and single qubit gate durations. Long Rabi flop (100 us), fit out carrier frequency of 2*pi*0.0716(1) MHz and Decay rate of 0.0107(7) 1 / us. See long flop. This is 1 mW of power on m1_m3 transition before we changed polarization. Fit with decaying cos function

$$\frac{1 + e^{-\lambda t}\cos(2\Omega t)}{2}, \tag{3.1}$$

3.2.2 Extracting Laser Offset and Magnetic Field

3.2.3 Spin Coherence Times

Individual gate fidelities are ultimately limited by loss of coherences of the two qubit states due to either dephasing or by the natural lifetime of the upper level. By our choice of ion and qubit levels, defined between the ground $4S_{1/2}$ state and the metastable $3D_{5/2}$ state, we can expect a lifetime limited coherence time of $\tau = 1.1$ s [?]. In practise, mainly due to imperfect tracking of laser frequency and magnetic field drifts (as mentioned above), we see coherence times dominated by dephasing. To discern between these two noise sources, we may exploit the fact that we have multiple Zeeman levels within our $3D_{5/2}$ state with varying magnetic field sensitivities. We also have the ability to define our qubit on the Zeeman split ground state, which decouples dephasing due to the laser from measured coherence times. We perform Ramsey scans with varying mid-sequence delay durations to extract the coherence times, an example of which can be seen in figure **??**. In characterising the spin coherence times, we hope to explore both the efficacy of the magnetic shielding surrounding the ion trap, as well as the stability of the 729 nm laser.

Figure **??** shows how the magnetic shielding effect coherence times of three transitions, XX, YY and ZZ, with magnetic field sensitivities of XX, YY and ZZ respectively. From this we find that without the shielding, we are strongly limited by external magnetic field noise, and with full sheilding we suppress this noise to where we are dominated by laser phase noise. To find the factor by which the magnetic field noise is attenuated, we can compare the coherence times of the laser phase insensitive transition with and without the box. We find an attenuation factor of XX, which is XXconsisitent with the expected attenuation factor of the mu-metal shielding.

With the shiedling in place, we compare the coherence times of the $4S_{1/2}, m_j = -1/2 \leftrightarrow 3D_{5/2}, m_j = -5/2$ with fibre noise cancellation (see section 2.5.3) and without, figure **??**. We find that the coherence time is improved by a factor of XX, with FNC enabled. Our current spin coherence time of XX ms is limited by the laser phase noise, and we expect to be able to push this to [ref R. Oswald] by improving the laser PDH stability. However, for the immediate planned experiments (see section **??**), these improvements will be a low priority due to other likely dominating error sources in the motion of our ions.

3.2.4 State Preparation and Measurement

To utilise two levels of the ion as a qubit, we need to be able to selectively prepare the ion into one of the Zeeman levels of the ground state. As mentioned in section 2.4, we are operating at a low magnetic field of 5 G, leading to a splitting between the Zeeman levels of less than 21 MHz, the natural linewidth of the 397 nm transition. This means we cannot optically pump using 397 nm frequency selectivity. Further, due to the constraint of beam geometry from the in vacuum optics, we can not use polarisation selectivity of the 397 nm transition. Instead, we use the narrow line width 729 nm laser on resonance with the $4S_{1/2}, m_j = +1/2 \leftrightarrow 3D_{5/2}, m_j = -3/2$ transition, and the 854 nm deshelving laser on resonance, to optically pump into the $m_j = -1/2$ Zeeman level we define as our qubit ground state.

To measure the qubit state of the ion, we apply the 397 nm and 866 nm lasers and count 397 nm photons scattered. From the level diagram shown in figure **??**, we can see that upon turning on the 397 nm laser, if we are in $|0\rangle$, photons will be scattered, and if we are in $|1\rangle$, then no photons will be scattered. To optimise the fidelity of measurement we ensure that the signal is discernible with low error from any background counts on the camera. In general, improving the number of signal counts can be achieved by tuning the 397 nm laser near to the transition resonance, by increasing the readout duration, or by increasing the percentage of scattered photons captured by the imaging system. Practically we desire that the readout step does not heat the motion of the ion and so we red detune the 397 nm laser to a similar setting as for Doppler cooling (see section 3.3.3). The parameters we use for readout are summarised in table **??**, and a typical histogram of readout counts for one and two ions can be seen in figure **??**. We find that the readout fidelity is XXX.

To measure the effect of state preparation and measurement error on longer experimental sequences, we will discuss randomised benchmarking in the following section 3.2.5. Due to the relevance here, we quote the measured state-preparation and measurement error (SPAM) of $\epsilon_{SPAM} = 1.46(6) \times 10^{-3}$.

3.2.5 Randomised Benchmarking

High fidelity unitary operations (gates) are essential for both near intermediate scale quantum computing and for reducing overheads in required physical qubits and operations in fault-tolerant schemes [**?**]. To evaluate the quality of both our state-prep and single qubit rotations, we employ randomized benchmarking (RBM) [**?**, **?**]. RBM consists of applying random combinations of a pre-chosen discrete set of gates to estimate an average error per gate. We chose the single-qubit Clifford group as our set of gates to evaluate. The single-qubit Clifford group is the set of unitaries which map the Pauli matrices to one another through conjugation. This can be thought of as the complete set of rotations of the Bloch sphere such that all valid combinations of the axis $(x \rightarrow \{\pm x,\ \pm y,\ \pm z\})$, $(y \rightarrow \{\pm x,\ \pm y,\ \pm z\})$, $(z \rightarrow \{\pm x,\ \pm y,\ \pm z\})$ are realized. There are 24 unitaries in this set. We followed the RBM protocol described in the Thesis [**?**] to evaluate our single-qubit gates. First the qubit is prepared in some known initial state, i.e. prepared in some chosen basis. A gate sequence is then applied which consists of multiple random Clifford gates followed by a final 'inverting' Clifford, where the 'inverting' Clifford is chosen such that the full sequence performs the Identity operation. The state is then measured in the same basis to find any deviations from the Identity being performed due to gate errors. This is repeated with the same preparation and sequence multiple times to calculate the probability that the Identity was performed - thus giving the sequence fidelity. These steps are repeated for many different random sequences with a range of sequence lengths. The decay model we fit to the fidelity versus number of Clifford gates is given by,

$$F(m) = \frac{1}{2}\left(1 + (1 - 2\epsilon_{SPAM})(1 - 2\epsilon_c)^m\right), \tag{3.2}$$

where $F(m)$ is the fidelity of the sequence of length m, ϵ_{SPAM} is the state-preparation and measurement error, and ϵ_c is the average error per Clifford gate. We use this method to bench mark our qubit transition. The Clifford gates are decomposed into sequences of $\pi/2$ and π pulses about either the $x-$ or $y-$axes. We probe up to $m = 100$ Clifford gates, and fit the decay of the fidelity to the above model. We measure the error per Clifford to be $\epsilon_c = 9.5(3) \times 10^{-4}$, while the SPAM error is $\epsilon_{SPAM} = 1.46(6) \times 10^{-3}$. The decay plot for this RBM sequence can be seen in figure **??**. The error bars are given by the standard deviation of the survival populations. There are on average 3.50 $\pi/2$ pulses per Clifford, with a typical $\pi/2$ duration of 1.3 μs.

3.3 Motion

3.3.1 Finding Motional Mode Frequencies

3.3.2 Motional Mode Stability

3.3.3 Cooling

For any interaction involving the motion of the ion, we require both the ability to prepare the motional state with high fidelity, and to subsequently measure this motional state to verify correct preparation. For entangling gates, and the creation of squeezed states which we are considering in this chapter, we assume that we begin in the motional ground state, or in other words, Fock state zero. Our initially trapped ions will be in some high temperature thermal state, (*given by the oven temperature and the PI laser momenta kicks*). We first doppler cool our ions, and then subsequently sideband cool them. We give a brief description of these two cooling processes here.

Doppler Cooling

Doppler cooling exploits the fact that incident light onto a moving ion will appear frequency shifted in the rest frame of the ion. For Doppler cooling of $^{40}Ca^+$, we apply both the 397 nm and 866 nm lasers. We initially red detune the 397 nm laser by around 100 MHz. This results in the preferential absorption of a quanta of 397 nm light by ions with a velocity vector antiparallel photon k-vector. After this absorption, the ion will be in the excited $4P_{3/2}$ state and spontaneously decay to either the $4S_{1/2}$, or the $3D_{3/2}$ emitting a photon of either 397 nm or of 866 nm respectively into a random direction. These two decay paths have a branching ratio of XX. As we desire many photon kicks to cool our ions, we repump the electron out of this metastable $3D_{3/2}$ level by applying an on resonant 866 nm beam. The absorption and sequential emission of this 397 nm photon will lead to a net reduction in the motional energy of the ion if the photon is emitted at a higher energy than when absorbed. The equilibrium temperature is given by the condition where the doppler cooling rate is equal to photon recoil heating of the ion. Assuming a Lorentzian absorption profile, the minimum temperature is given by,

$$T_{Doppler} \approx \frac{\hbar\gamma}{2k_B}, \tag{3.3}$$

where $\hbar$ is the reduced Planck constant, γ is the natural linewidth of the transition, and k_B is Boltzmann's constant.
For $^{40}Ca^+$, the natural linewidth of the 397 nm transition is $\frac{\gamma}{2\pi} = 21$ MHz, leading to a Doppler temperature of approximately 0.5 mK. Given a radial mode frequency of $\frac{\omega}{2\pi} = 4$ MHz, and the mean occupation number of the oscillator being given by,

$$\bar{n} = \frac{1}{e^{\hbar\omega/k_BT} - 1}, \tag{3.4}$$

we find the final thermal distribution to have an expected Fock state of $\bar{n} = 2.3$. Using parameters summarised in table **??**, we find practically the final temperature after Doppler cooling to be around XXX mK.

Sideband Cooling

To further cool the ions toward their motional ground state, we use resolved sideband cooling. The motion of the ion, described by a harmonic oscillator, modulates the transition frequencies of the ion, leading to sidebands at multiples of the motional frequency. For the $4S_{1/2} \leftrightarrow 3D_{5/2}$ transition, at appropriate laser intensity and motional mode frequencies, these sidebands can be resolved spectroscopically. The pulsed sideband technique we employ consists of red sideband pulses, followed by deshelving, and repumping pulses on the 854 nm and 866 nm transitions respectively. An example pulse sequence can be seen in figure **??**, and experimental parameters we use are summarised in table **??**.
To verify the efficacy of our sideband cooling, we perform thermometry experiments by driving on resonance red sideband (RSB) and pi-RSB-pi pulse sequences. We record the time dynamics of population flopping as we vary RSB pulse length. In the case of Fock state zero, we expect to see a strong signal on the RSB, and no signal on the pi-RSB-pi pulses. We fit a thermal Fock state distribution (with truncation at Fock state = 100) to these signals to extract the mean occupation number, and $\eta\Omega$, the carrier Rabi frequency multiplied by the Lamb-Dicke parameter. A typical thermometry scan after Doppler and sideband cooling can be seen in figure **??**. We find that the mean occupation number after sideband cooling is $\bar{n} = 0.03()$, and

$\eta\Omega = XX$ MHz.

Optimisation of the cooling parameters can be roughly performed by fitting temperature while scanning RSB pi-pulse durations, total number of pulses, repumping and deshelving times. One can optimise for minimum temperature, however it is also important to optimise for total cooling duration. For single ion, single mode experiments, this duration is often a non-issue, however for multi-ion crystals, any interaction involving the motion, may require the sequential sideband cooling of multiple motional modes. This can not be easily parallalised due to the requirement that the RSB pi-pulse is performed near resonance to one of the motional sidebands. This sequential cooling strategy can be either limiting when heating and cooling rates are comparable, or in the best case, painful due to long data collection times.

To mitigate this issue, other sub-Doppler cooling techniques with larger accepted frequency bandwidths may be employed. Examples are dark-resonance cooling[], and electromagnetically induced transparency (EIT) cooling[], and Sisyphus cooling[]. These techniques are not yet implemented in our system, but will be likely additions once we move to larger ion crystals.

3.3.4 Heating Rates

As mentioned, the cooling of our ions is only relevant if we have acceptable heating rates. Heating of the motion is predominantly caused by the ion trap itself. This can be due to imperfections in the surface of exposed dielectric and metals causing stray fields, or can be due to noise on the DC and RF drive voltages[]. Noise due to the surface of the trap can be mitigated by increasing ion-electrode distances, or by using traps with smaller surface area direclty exposed to the ion. In our case, as mentioned in sections 2.2, the NPL trap has an electrode ion distance somewhat larger than most surface traps, but less than that of a macroscope blade or rod style trap. To verify the heating rate of our system, we performed a series of thermometry scans whilst varying some delay time between cooling and thermometry pulses. A typical plot can be seen in figure **??**. We find that the heating rate of our system is approximately 33(3) quanta per second on the upper radial 4 MHz mode on one ion.

It is expected that the heating rate will be larger for lower frequency motional modes if we assume uniform electric field noise. We also verify this by looking at heating rate on the radial mode while varying the axial mode frequency. This is a useful diagnostic to check for unexpected heating at certain frequencies, perhaps due to RF noise in the lab. We find....

3.3.5 Motional Coherence Times

3.4 Experimental Control

3.5 Spin-Dependent Forces

The spin-dependent force (SDF) is (planned to be) heavily used throughout this thesis. We use the Mølmer-Sørensen (MS) scheme [?], to generate the SDF via a bichromatic laser field. Bichromatic refers to the simultaneous application of two tones symmetrically detuned around the qubit carrier frequency, with absolute detuning approximately equal to the motional mode frequency, $\delta \approx \omega_m$. The resulting interaction, when ignoring off resonant and higher order couplings, is given by,

$$\hat{H}_{MS} = \hbar\eta\Omega\ \hat{\sigma}_\phi \cos(\delta t) \left(a e^{-i\omega_m t} + a^\dagger e^{i\omega_m t} \right), \tag{3.5}$$

where η is the Lamb-Dicke parameter, Ω is the carrier Rabi frequency, $a(a^\dagger)$ is the lowering (raising) operator, and σ_ϕ is the Pauli operator with ϕ being in the $x,\ y$ -plane. Applying the rotating wave approximation, and defining $\delta_g = \delta - \omega_m$, we find that the interaction Hamiltonian can be approximated to,

$$\hat{H}_{MS} = \frac{\hbar\eta\Omega}{2}\ \hat{\sigma}_\phi \left(a e^{-i\delta_g t} + a^\dagger e^{i\delta_g t} \right). \tag{3.6}$$

The SDF realises a displacement of the motional state in phase space. We may control the trajectory of this displacement by varying δ_g: on resonance, $\delta_g = 0$, we see linear trajectories, whilst off resonance, $\delta_g \neq 0$, we see cyclic trajectories where after some time $t = 2\pi/\delta_g$, the motion returns to the initial state. We shall exploit this control in both the two-qubit entangling gate experiments, as well as in the creation of squeezed states.

3.5.1 Calibrating the SDF

The MS interaction is widely used in ion trap experiments due to it being robust against varying intial motional states, and to the effects of heating during the pulse sequence. However, in our use case, the SDF is sensitive to miscalibration of the central qubit frequency, the balancing of power in the two tones, and to motional mode drifts. The first two miscalibrations appear as the SDF having.. We use both "detuning" and "duration" scans to calibrate the SDF. The "detuning" scan is performed by varying the detuning, $delta_g$, of the interaction, whilst the keeping the SDF duration constant, whilst the "duration" scan keeps detuning constant and varies the SDF duration. Duration scan fit given by,

$$P_{\downarrow,\mathrm{th}} = \frac{1}{2}\left[1 + e^{-4\left(\bar{n}+\frac{1}{2}\right)|\alpha(t)|^2}\right], \tag{3.7}$$

from Burd thesis, equation 3.30. $|\alpha(t)| = \Omega_{sdf}t/2$, fit find $\Omega_{SDF} = 2\pi \times 6.6(2)$ kHz, fitted with $\bar{n} = 0.03$ Fitting detuning scan using $|\alpha(t)| = \Omega_{sdf}\sin(\delta t/2)/\delta$, find $\Omega_{SDF} = 2\pi \times 6.3(2)$ kHz, fitted with $\bar{n} = 0.03$.

3.6 Two-Qubit Entangling Gates

Parity contrast = 0.90(2) Population 01 10 = 0.05(2)

3.6.1 Collective Motion of Two Ions

3.6.2 Mølmer-Sørensen Gate

Theoretical Background to the MS Gate

Experimental Implementation of the MS Gate

3.7 Creating Squeezed States

3.7.1 Calibrations

Chapter 4

Outlook

4.1 Appendix