

Towards Cryogenic Dual-species Tweezer Arrays

PhD Thesis

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Abstract

This thesis presents work towards a dual-species array of neutral Rb and Cs atoms built within a cryogenic apparatus for quantum computing. The project aims to demonstrate non-destructive readout through the naturally non-resonant imaging wavelengths of Rb and Cs atoms and to show improved trap lifetimes by using cryogenics for improved vacuum.

This thesis presents analysis of interactions between Rb and Cs Rydberg d -orbitals where Förster resonant d -states are identified as candidates for high fidelity qubit operations. We model C_kZ gates up to $k = 4$ and demonstrate competitive qubit fidelities with d -orbitals, with further enhancements attainable through improved gate protocols. We then move on to discuss the design of a next generation experiment for realising dual-species arrays in a cryogenic environment. This new experiment utilises a He cryostat to reach a base temperature of ~ 7 K to take advantage of cryopumping to lower the vacuum pressure for increased trap lifetimes. In-vacuum aspheric lenses allows for the projection of independent tweezer arrays for trapping single Rb and Cs atoms. This thesis then discusses loading of a cryogenic Cs MOT before loading single Cs atoms in holographic optical dipole traps. Several enhancements in vacuum are used to demonstrate increased trap lifetimes from ~ 10 s up to ~ 1 minute where further enhancements will be attainable with active in trap cooling. Increased trap lifetime is vital for the realisation of quantum computing at scale.

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Chapter 1

Introduction

Quantum computing has seen increasing interest over the last few decades with its promise of providing a computational speed up for a small set of mathematical problems. This not only includes the famous Shor’s algorithm [1] for prime number factorisation and Grover’s search [2], but also applications in quantum chemistry [3] and material science [4]. Neutral atoms confined in arrays of optical tweezers have emerged as a promising platform for quantum computing due to their potential for scalability and flexibility [5, 6]. Across the world there are many research groups dedicated to the development of this platform [7–13] as well as industry groups who are looking to commercialise neutral atom hardware [14–17].

Here in the UK, the government recently announced investment of £670 m in quantum computing. The UK has had a history of neutral atom research with groups dedicated to molecules in tweezers [18], Rydberg physics [19–22] and quantum computing [23–25]. The National Quantum Computing Centre has recently been established in part to accelerate the UK’s progress in the field of quantum computing, including with neutral atoms arrays. Here at the University of Strathclyde, the SQuAre project [26] has been a trailblazer for bringing neutral atom quantum computing to the UK, and it is within this context that the next generation of experiment discussed here is being built.

In this thesis, we present work towards the realisation of a cryogenic dual-species tweezer array of Rb and Cs atoms. The platform that is presented in this thesis will

act as a test bed for error correction by utilising the non-destructive readout offered by dual-species arrays. Improved trap lifetimes through the implementation of a cryostat will be demonstrated as a route to scalability. Finally, interactions between Rb and Cs Rydberg states shall be identified for the use of enhanced multi-qubit gate fidelities.

This first chapter will provide a brief overview of quantum computing in general before presenting neutral atom arrays as a platform. We will then discuss the key aims of the QuERy project, which guides the work presented in this thesis.

1.1 Quantum Computing

The fundamental unit of information in quantum computing is the quantum bit (or qubit for short). The qubit is defined as a two-level quantum system and is the analog of the classical bit. It is convenient to represent a qubit in Dirac notation as

$$|\Psi\rangle = \alpha |0\rangle + \beta |1\rangle, \quad (1.1)$$

where $|0\rangle$ and $|1\rangle$ are orthonormal basis vectors, and α and β are complex coefficients whose squared modulus represents the probability of the associated basis vector. The non-discrete nature of quantum systems is known as superposition and is one of the tools utilised in quantum computing. Collecting information from a qubit requires the projection of the system into a single basis state when measured. The qubit state vector must obey the following normalisation condition

$$|\langle\Psi|\Psi\rangle|^2 = |\alpha|^2 + |\beta|^2 = 1, \quad (1.2)$$

where the probability of projecting into $|0\rangle$ and $|1\rangle$ is $|\alpha|^2$ and $|\beta|^2$ respectively.

Rotations of the qubit state vector are controlled through interactions with external fields — the specifics of which depend on the choice of platform. Mathematically, we represent the evolution as a linear transformation where the interaction is represented by the operator matrix

$$\hat{O} = \sum_{ij} k_{ij} |i\rangle \langle j|, \quad (1.3)$$

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where k_{ij} is the value of a given matrix element. We can therefore represent the action of $\hat{O}$ on the state vector $|\psi\rangle$ as

$$\hat{O}|\psi\rangle = \lambda|\psi'\rangle, \quad (1.4)$$

where the λ prefactor is a scalar constant. As a nod to their classical counterparts, it is common to refer to operators which act on qubits as quantum gates and these are the basic building blocks of quantum algorithms. The set of Pauli operators are examples of basic but important quantum gates which perform bit and phase flip operations

$$\begin{aligned} \hat{I} &= \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}, & \hat{\sigma}_x &= \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \\ \hat{\sigma}_y &= \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, & \hat{\sigma}_z &= \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \end{aligned} \quad (1.5)$$

Any arbitrary rotation by an angle θ can be represented by the Pauli operators with

$$\mathcal{R}_\chi(\theta) = e^{-i\theta\hat{\sigma}_\chi/2}, \quad (1.6)$$

where $\chi \in \{x, y, z\}$. A quantum gate which by conjugation normalises to a Pauli operator is considered a Clifford gate which can be efficiently simulated classically [27]. For example, the Hadamard gate

$$\hat{H} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 1 \\ 1 & -1 \end{pmatrix}, \quad (1.7)$$

can be described as a member of the Clifford group as

$$\hat{H}\hat{\sigma}_x\hat{H}^\dagger = \hat{\sigma}_z. \quad (1.8)$$

A set of Clifford gates does not allow for universal quantum computing, however, the inclusion of a single non-Clifford gate does [28].

The true power of quantum computing can be realised when extending systems to more than one qubit. As an example, consider a two-qubit system of qubit A and qubit

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B where we can represent each qubit in isolation in the form of Eq. 1.1. However, it is far more convenient to consider the Hilbert space of the coupled qubit system which can be done by performing a tensor product between each qubit,

$$|\Psi\rangle_A \otimes |\Psi\rangle_B = \alpha |0\rangle_A \otimes |0\rangle_B + \beta |0\rangle_A \otimes |1\rangle_B + \gamma |1\rangle_A \otimes |0\rangle_B + \delta |1\rangle_A \otimes |1\rangle_B, \quad (1.9)$$

where we now have a four dimensional system. Similarly, quantum operators can be represented in the two-qubit Hilbert space

$$\hat{O}_{AB} = \hat{O}_A \otimes \hat{O}_B. \quad (1.10)$$

To simplify notation we define $|\eta\rangle_X \otimes |\chi\rangle_Y = |\eta\chi\rangle_{XY}$.

Considering the form of Eq. 1.9, it is possible to realise quantum states where the two qubits can no longer be considered independently from one another. This is known as entanglement and is a fundamental resource for quantum computation. Consider the entangled state,

$$|\Phi^+\rangle_{AB} = \frac{1}{\sqrt{2}} \left[|00\rangle + |11\rangle \right]_{AB}, \quad (1.11)$$

which by definition we cannot represent in a separable form. In a practical sense, this means that if we were to measure qubit A in its $|0\rangle$ state, we can infer that qubit B has been projected into the state $|0\rangle$.

Entangled states can be created utilising multi-qubit gates. These canonically take the form of controlled gates, which apply a unitary operation on a set of target qubits conditioned on the state a set of control qubits. One of the most common two-qubit gates is the controlled- X gate (more commonly referred to as a CNOT gate) which can be represented in matrix form as

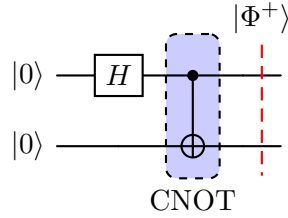
$$\text{CNOT} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{pmatrix}, \quad (1.12)$$

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and has the following action on the computational basis states

$$\begin{aligned}
 |00\rangle &\xrightarrow{\text{CNOT}} |00\rangle, \\
 |01\rangle &\xrightarrow{\text{CNOT}} |01\rangle, \\
 |10\rangle &\xrightarrow{\text{CNOT}} |11\rangle, \\
 |11\rangle &\xrightarrow{\text{CNOT}} |10\rangle.
 \end{aligned} \tag{1.13}$$

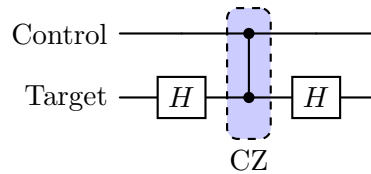
Starting with two qubits in the $|00\rangle$ state, a maximally entangled state can be created by first applying a Hadamard gate to the control qubit, followed by a CNOT gate, as shown in the following circuit diagram



Another common two-qubit gate is the controlled-Z (CZ) gate which is represented in matrix form as

$$\text{CZ} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}. \tag{1.14}$$

The CZ gate is naturally implemented on neutral atom systems using Rydberg blockade (discussed in Part I). A CNOT gate can be constructed by rotating the target qubit with a $\hat{H}$ gate, applying the CZ gate, and finally apply a second $\hat{H}$ gate to the target qubit:



On the surface, entanglement appears to break causality, as information can seemingly transmit faster than the speed of light. This is known as the Einstein-Podolsky-Rosen (EPR) paradox which introduced the concept of hidden variables as a means

to preserved locality in quantum mechanics [29]. In the 1960s John Bell devised a test which could use the correlated observations of an entangled system to determine whether or not hidden variables do govern quantum systems [30]. It was not until the 1980s that this was put to the test and the results indicated that entanglement can not be described by hidden variables, hence, indicating entanglement is fundamental to quantum mechanics [31].

It is often portrayed that the nature of qubits as described above allows for generally more efficient computation on a quantum computer when compared to a classical computer [32]. This is in fact not true, although it can be said that any classical algorithm can be computed on a quantum computer [33], in most instances this will be far less efficient — both computationally and in terms of resources — than if it was performed on a classical computer [34]. There are, however, a set of mathematical operations which can be encoded on a quantum computer and provide results more efficiently than a classical computer [1–4, 35]. Therefore, a fully-fledged quantum computer will supplement classical compute only when necessary. This is in the same vein as a graphical processing unit (GPU) which can efficiently compute floating point arithmetic and matrix operations, but does so without replacing the traditional computer. As such, it may be more apt to call a quantum computer a quantum processing unit (QPU) which will be one of a plethora of compute tools available to a user.

An example of a quantum algorithm which sees a speed up on a quantum computer is Grover’s algorithm, which can search an unsorted list for a single unique index [2]. For a list of N entries, Grover’s performs a search with $\sqrt{N}$ queries whereas the fastest classical algorithm will take N queries. Another quantum algorithm is Shor’s algorithm [1] for finding prime factors which also presents a polynomial speed-up over classical algorithms. Shor’s algorithm is of particular interest as many encryption protocols, including RSA public key encryption, depend on the fact that a classical computer will not realistically factor large prime numbers. These algorithms are often referred to when discussing the advantages of quantum computing, however, quantum chemistry [3] and material science [4] both stand to gain from speed-ups through quantum computing.

1.2 Qubit Platforms

In the previous section we briefly discussed the mathematical representation of quantum computing from which it should be clear that realisation of a qubit requires two components. Firstly, we need a two-level quantum system for encoding information, as per Eq. 1.1; secondly, we require a means of interacting with the qubit. Additionally, any useful quantum computation also requires the quantum system to have long coherence times and a mechanism for high-fidelity state-selective readout [36].

Real physical systems do not naturally fit into these requirements, as interactions with the environment decohere quantum states before any measure of computation can occur. However, in a laboratory environment, it is possible to engineer the conditions which effectively reduce decoherence pathways long enough to allow coherent interactions with quantum systems. Only under these sorts of conditions is it possible to isolate two levels of a real quantum system to encode a qubit.

Below, we discuss some common quantum systems which are used to realise qubits as a means to illustrate the current field of quantum computing. An important take away here is that there are many approaches by which to realise a qubit, each with their own merits. However, each method does have limitations which will need to be overcome to realise a useful quantum computer. In the next section we will see how neutral atom qubits have the potential to overcome some of these hurdles.

Superconducting Qubits

There are a number of ways to utilise superconductors to realise qubits, however, one of the most common methods is known as a transmon which consists of a shunt capacitor in parallel to a Josephson junction [37, 38]. The qubit is formed by using the non-linear inductance provided by the Josephson-junction to break the harmonicity of the underlying resonant microwave circuit. This results in isolation of the two lowest energy levels which can encode $|0\rangle$ and $|1\rangle$ states [39]. To access the lowest energy states of the harmonic oscillator, the superconducting qubit must be cooled to temperatures in the low milliKelvin range.

Single qubit gates can be performed on superconducting qubits using microwave pulses [40], while multi-qubit gates are often performed by implementing additional transmons to mediate the coupling strength between the desired qubits [37, 41]. Single qubit gate fidelities are typically of the order $\mathcal{F} > 0.999$ [40] and two-qubit gate fidelities of $\mathcal{F} > 0.99$ [37, 41, 42]. Qubit operations are also typically performed on the order of 100 – 1000 ns, however, there have been demonstrations for gate times as low as 13 ns [43]. Coupled with coherence times on the order of hundreds of milliseconds [44, 45], these fast gate times allow a large number of sequential gate operations to be applied within the coherence time, as required for implementation of quantum algorithms.

There are some limitations with superconducting qubit platforms with perhaps the biggest issue being scalability. Each superconducting qubit must be fabricated and physically placed in the experimental system, which necessitates the creation of large dilution fridges to house superconducting chips at scale [46]. Fabrication at scale will also prove a challenge as each qubit is inherently non-identical which means each one requires separately optimised gate sequences to operate. Superconducting hardware also has wiring challenges as each qubit needs multiple control lines.

Trapped Ion Qubits

Neutral atom qubits, which are discussed in more detail below, share a lot of similarities with trapped ion qubits. The qubit is encoded on a single ion, with one or more qubits constrained in either Penning traps [47–50] or Paul traps [51]. The former of these traps has seen recent development towards miniaturisation by using a multi-electrode design enabling qubit shuttling [52–55]. The computational basis for trapped ions is typically encoded in the hyperfine ground states of the atom, or between the ground state and a long-lived meta-stable excited state. Two-qubit gate interactions are mediated by a combination of light-matter interactions and the resulting motional changes within the vibrational state of the trap [56]. This form of gate has realised very high two-qubit gate fidelities of $\mathcal{F} > 0.999$ [57], with recent state-of-the-art boasting $\mathcal{F} \approx 0.9999$ for two-qubit gates [58].

These high gate fidelities are impressive, however, there are again challenges with

operating at scale. It is hard to stop trapped ions from interacting, which means gate fidelities have poor scaling as the number of qubits increase. This is due to increases in the density of motional modes which can result in couplings to unwanted states during gate pulses [59]. The development of shuttling architectures [60] — where qubits are moved between storage and entanglement zones — and zoned based architectures [57] — where gates are applied on a site by site basis — has been motivated as a way to circumvent unwanted interactions.

Silicon Qubits

A lesser known candidate for quantum computing is the semiconductor based silicon qubit which presents several advantages when compared to other platforms. The qubit is encoded within the electron or nuclear spin of a silicon quantum dot (QD) [61, 62]. Single qubit operations are performed via single qubit RF microwave pulses which can achieve gate fidelities of $\mathcal{F} > 0.999$ [63]. Multi-qubit gates are performed through spin-dependent tunnelling between adjacent QDs which has shown two-qubit fidelities of $\mathcal{F} > 0.99$ with gate time of ~ 100 ns [64]. Silicon qubits also show superb coherence times¹ [65, 66] which, coupled with the short gate times, makes this platform promising as a means of realising quantum computation.

The largest qubit registers to date contain only ~ 12 qubits [67, 68], so scaling will have to be addressed in order to realise a useful system using silicon spin qubits. However, silicon based semiconductors are a mature industry which will make future development of these platforms easier [68].

Photonic Qubits

Photonic qubits are another system for implementing quantum computing where qubits are made up of single photons [69]. In these systems, computational basis states are typically encoded in some property of the photon such as polarisation and qubit manipulation is implemented using optical elements such as waveplates, beamsplitters and mirrors. Photonic systems have demonstrated single qubit fidelities $\mathcal{F} \approx 0.99$ [70].

¹Single electron lifetimes have been demonstrated in excess of 2 s [65].

Multi-qubit gates prove harder to implement due to the generally non-interacting nature of photons. However, gate fidelities of $\mathcal{F} = 0.88$ within circular Bragg resonators have been demonstrated [71].

The scalability of photonic qubits proves a challenge due to the optical setup requirements for implementation. For one, photon sources are not 100% efficient at producing single photons. Additionally, errors can arise from imperfect readout of photodiodes affecting the ability to perform high-fidelity state detection. The linear nature of photonic qubits makes it a challenge to implement programmable quantum circuits. One approach is to create photonic integrated circuits with programmable interferometers to allow for reconfigurable circuits to be implemented in a similar manner to FPGAs [72]

1.3 Quantum Computing with Neutral Atoms

The platforms briefly discussed in the previous section provide a glimpse at the various means in which qubits can be engineered. Each architecture has its specific strength, but also drawbacks which will limit routes to fault tolerance. Superconducting qubits for example may be fast, but are very susceptible to noise; trapped ion qubits can be confined and scaled with relative ease, but fall short when extending multi-qubit qubit gates. Silicon qubits show some promise, however, they have yet to show true scalability.

Neutral atoms confined within arrays of optical dipole traps (ODTs) — or tweezer arrays — have emerged in the last decade as a candidate for quantum computing [5, 6]. First proposed in the early 2000s, these qubits are typically made up of single alkali atoms such as rubidium (Rb) [7, 73] and caesium (Cs) [8, 9, 13, 26], or single alkaline atoms such as strontium (Sr) [74, 75] and ytterbium (Yb) [76, 77] — Fig. 1.1 shows how each of these atoms can encode qubits. Neutral atomic qubits have several inherent benefits. For one, each atom is identical which means gate sequences can be applied to all atoms without the need for tuning of individual qubits. Secondly, they are highly scalable, as increasing the number of qubits can be done by adding additional

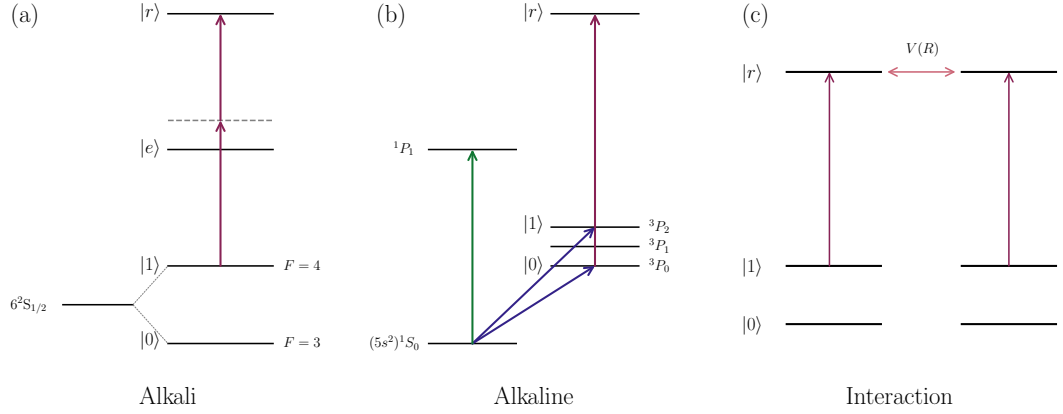


Figure 1.1: (a) An example of alkali (Cs) qubit where the computation basis is typically encoded in the hyperfine splitting of the $nS_{1/2}$ ground state and qubit-qubit interactions are performed via a Rydberg state $|r\rangle$ which has transition couplings to the $|1\rangle$ state. (b) An example of an alkaline (Sr) qubit where the computational basis can be encoded in the ground state, the metastable triplet state, or between the ground state and triplet state. Imaging can be performed on the green transition. (c) Interactions between two qubits is performed by excitation to Rydberg states where the double Rydberg state is forbidden under the Rydberg blockade mechanism.

tweezers. In practice, this requires some additional laser power, and a change in phase pattern on an spatial light modulator (SLM) or additional tones applied to an acousto-optic deflector (AOD). Thirdly, by their nature, neutral atoms do not readily interact which avoids unwanted couplings between qubits. However, multi-qubit gates can be engineered by excitation to high principle quantum numbers known as a Rydberg state where the dipolar nature of the system of neutral atoms becomes apparent.

There are a few methods for making a qubit out of a neutral atom. Rb and Cs qubits are a popular choice as their atomic structure is well understood [78, 79]. In these atomic species, it is common to encode the computation basis $\{|0\rangle, |1\rangle\}$ in the Hyperfine ground states, for which two energy levels exist. This encoding is shown in Fig. 1.1a for a Cs atom. The hyperfine ground states are separated by 9.2 GHz (6.8 GHz in the case of ^{87}Rb) which means single qubit rotations can be performed with microwave pulses [23, 80, 81]. Faster gate pulses can be performed using Raman transitions via an intermediate p -state transition [82].

Multi-qubit gates are mediated by highly excited Rydberg states which can impart an energy shift on a neighbouring atom's energy levels via the dipole-dipole interaction. At the suitably small interatomic distances and with low enough Rabi frequencies, not more than one atom may be excited to the Rydberg state, a phenomenon known as Rydberg blockade. The Rydberg energy shift typically scales with interatomic separation R as $V(R) = C_6/R^6$ where we can define the blocked condition as $V(R) > \Omega$ where Ω is the Rabi frequency. For an atom in a Rydberg state, its blockade radius is defined as $r_B = \sqrt[6]{|C_6|/\Omega}$ which is typically $\sim 5 - 10 \mu\text{m}$ for $50 \leq n \leq 100$ values used in tweezer array experiments. Rydberg blockade is the underlining principal of many multi-qubit gate sequences used for neutral atom qubits [83–85] and is discussed in more detail in Part I.

Early demonstrations of ODTs allowed for the confinement of 1000s of atoms collected from a magneto-optical trap (MOT) [86] which could be held for ~ 200 ms limited by collisions with background gas. The first single atom trapping in ODTs was achieved in the early 2000s [87], where a single trap at 1064 nm was used to confine a Cs atom. The same experiment also demonstrated state preparation and detection in the hyperfine ground state, a vital requirement for engineering qubits. Arrays of single trapped atoms can be made by using an SLM to modulate the phase of a laser [88]. This allows for the generation of reconfigurable arrays of up to thousands of optical tweezers [89] in 1D [90], 2D [88] and even 3D arrays [91] proof sentence. State-of-the-art arrays have demonstrated 6100 atoms trapped in a 12000 sites [8]. SLMs also allow for the generation of arbitrary geometries, allowing for novel designs for a quantum register [7]. These results highlight the relative ease of scalability of tweezer arrays when compared to other qubit platforms. Fig. 1.2 shows a schematic of how SLMs can be used to generate a dual-species array of Rb and Cs atoms.

SLMs are excellent at producing static arrays, however, their refresh rates are limited to $\lesssim 120$ Hz. There have been demonstrations of SLM technologies which boast refresh rates of 10 MHz [92], however, these are not yet viable for use at scale. Additionally, ODTs are typically limited to single atom loading rates of 50%, although techniques have demonstrated loading rates of $> 80\%$ [11, 93]. Multiplexing RF tones

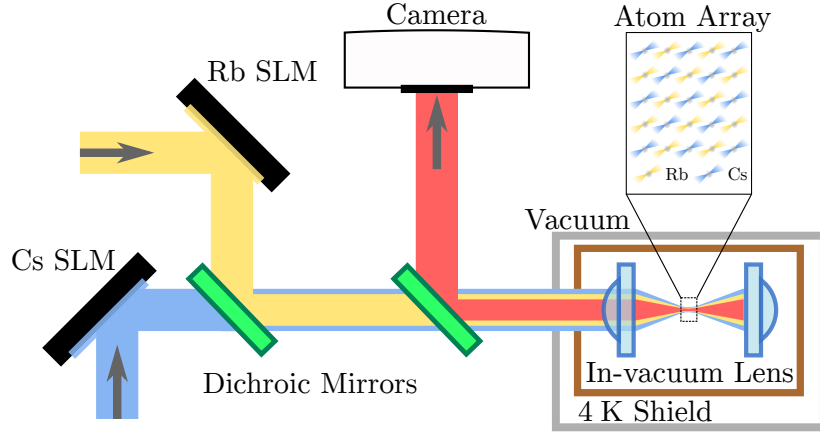


Figure 1.2: Schematic showing tweezer array generation for a cryogenic dual-species array. Each atomic specie’s array is generated by a separate SLM and is projected via an in-vacuum lens. In this example, a single camera is used for fluorescent imaging of the array.

through an AOD can be used to generate spot arrays which can be manipulated with ramped RF tones at the microsecond time scales [94]. Defect free array generation can be achieved with stochastic loading of a static SLM arrays with rearrangement of atoms performed in 2D with a pair of crossed AODs [95].

For scaling to thousands or hundreds of thousands of trapped atoms, SLMs and AODs run into issues with laser power heating the optical elements. To produce even larger arrays, static optical elements called metasurfaces can be used. These optics imparts a phase shift on light which is either reflected or transmitted off the element [96]. Metasurfaces have only so far demonstrated trapping of ~ 100 Sr atoms, however, arrays of 360,000 sites projected onto a CCD camera has been demonstrated [74].

1.3.1 Recent Developments

Currently, there are over a dozen research groups and companies across the world who work with neutral atom systems for quantum computation. Some of these groups are working on improving specific problems, while others are taking a holistic approach for implementing all-in-one systems for quantum registers. Here we highlight some of the recent advancements in this field.

For single qubit operations, microwave gates have demonstrated high gate fidelities

for $\mathcal{F} > 0.9999$ [23], however, these have limited gate times on the tens of microseconds scale and can only be applied globally across the array meaning these gates are not suitable for local control. Faster gate times can be achieved by instead utilising a two-photon Raman transition. Single qubit Raman gates have gate times of $< 10 \mu\text{s}$ and have demonstrated fidelities of $\mathcal{F} > 0.999$ [82]. Raman gates have also demonstrated local addressing using an AOD [7].

Two-qubit gates have also seen recent advancement with improved laser phase noise [97] which has allowed for higher Rydberg driving coherences for improved gate fidelities. Improvements in protocols to reduce gate times with robust phase accumulation across few or even single pulses have allowed for robust control of gate sequences [84, 85, 98, 99]. These gate sequences have spurred the realisation of two-qubit gate fidelities $\mathcal{F} > 0.99$ [7, 14, 85, 100, 101]. Rydberg gates have been performed with AODs to direct light for local control [14].

Reconfigurable mid-circuit operations [102] and zone based architectures [7, 103] have been presented as a means for practical implementation of quantum algorithms. Neutral atom Rydberg arrays have seen advancements beyond digital quantum computing with recent results showing improvements in the field of classical optimisation with analogue quantum computing [16, 104]. Recent results have demonstrated Rydberg arrays for solving maximum independent sets (MIS) [105, 106], weighted MIS [24], quadratic unconstrained binary optimisation (QUBO) [24] and integer factorisation [107]. Programmable analogue simulation have been used to demonstrate 1D systems of many body scars [108] and 2D phase diagrams [109].

1.3.2 Quantum Error Correction

The previous section shows that relatively high gate fidelities can be achieved with neutral atom qubits. However, when considering computation at scale, we require qubits to effectively be error free. In reality, this is impossible to achieve when encoding an algorithm on a set of physical qubits due to naturally occurring decoherence pathways and unwanted interactions that can never be completely suppressed. Consider, for example, a single qubit used in a quantum algorithm that requires 10,000 sequential two-qubit

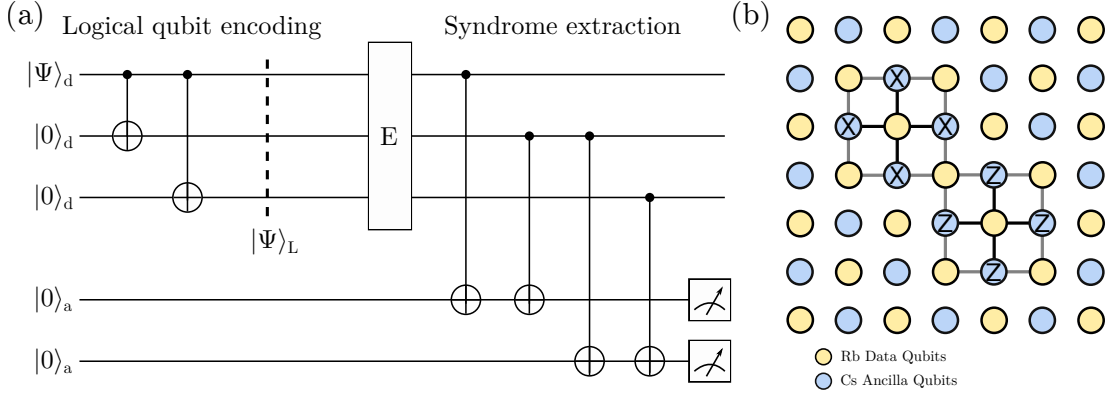


Figure 1.3: (a) Quantum circuit for realising the three-qubit error correction code [110]. The encoding section spreads the quantum information across the three data qubits with two CNOT gates. The unitary operator across the data qubits (labelled E) corresponds to a bit flip error occurring. Syndrome extraction is performed by application of a set of CNOT gates between the data and ancilla qubits, the latter of which are measured to identify errors which may have occurred. (b) Schematic of a surface code with a chequerboard grid of alternating ancilla and data qubit. Stabiliser measurements are performed between adjacent ancilla and data qubits.

gate operations. Even if each gate can be performed with fidelities of 0.9999, the net contribution of all gates at the end of the algorithm results in a fidelity of $\mathcal{F} = 0.3679^2$. A fidelity this low, across many thousands of qubits makes the result of the algorithm useless. Therefore, some form of error correction must be implemented to ensure the total algorithm fidelity is not limited by the gate fidelities [110].

A simple example of error correction is repetition codes, which spreads the quantum information of a logical qubit state across multiple physical qubits. To illustrate how one may implement a repetition code, we shall use the example of the three qubit code [110] where one logical qubit is made of three physical qubits. In this picture, we can encode the logical computational basis states as

$$|0\rangle_L = |000\rangle, \quad (1.15a)$$

$$|1\rangle_L = |111\rangle, \quad (1.15b)$$

where $|\chi\rangle_L$ denotes the logical qubit state and $|\chi\chi\chi\rangle$ represents the physical encoding

²This does not account for the decoherence of the quantum system which will further reduce the value of $\mathcal{F}$.

using three qubits for $\chi \in \{0, 1\}$. The encoding mechanics is shown in Fig. 1.3 which uses a series of CNOT gates to map the first qubit state across three qubits. A crude way of identifying and correcting for errors on $|\chi\rangle_L$ is to measure all three physical qubits, and use a majority vote system to decide what the output state of the logical qubit is. This may identify some errors, however, the method is destructive and not suitable if the logical qubit is in some superposition state which must be preserved.

To instead perform measurement in a non-destructive manner, we introduce the notion of data and ancilla qubits where the data qubits encode the qubit state and ancilla qubits are used to identify errors which have occurred during computation. In the example above, we can initialise the logical basis states as

$$|0\rangle_L = |000\rangle_d \otimes |00\rangle_a, \quad (1.16a)$$

$$|1\rangle_L = |111\rangle_d \otimes |00\rangle_a, \quad (1.16b)$$

where the subscript d denotes data qubits and the subscript a denotes ancilla qubits. After computation, we can identify bit flips of the data qubits by performing a parity check on the ancilla qubits. This is shown in Fig. 1.3 where a series of CNOT gates act to compare data qubits and flip one or both of the ancilla qubits if a bit flip error occurred. Measurement of the ancilla qubits provides us with a syndrome to identify not only that a data qubit has flipped, but which qubit has in fact flipped,

$$|10\rangle_a \implies |000\rangle_d \longrightarrow |001\rangle_d, \quad (1.17a)$$

$$|11\rangle_a \implies |000\rangle_d \longrightarrow |010\rangle_d, \quad (1.17b)$$

$$|10\rangle_a \implies |000\rangle_d \longrightarrow |100\rangle_d. \quad (1.17c)$$

The three qubit code is a nice illustrative example for quantum error correction, however, it is not robust against phase errors. There are a few examples of error correction codes which are robust against bit flip and phase errors, such as the Steane code [111]. However, many repetition codes will be hard to implement on real hardware due to the need for long distance connectivity. With this in mind for neutral atom

arrays, the surface code is an example of an error correction code which fits rather naturally into the implementation restrictions as it can be implemented on a square geometry and requires only nearest neighbour interactions. For surface codes, the ancilla qubits are used for stabiliser measurements of their neighbouring data qubits. Surface codes do however require quite a bit of overhead to realise a single logical qubit. For example, the minimum number of physical qubits to realise a single logical surface code qubit is 17 [110, 112], however, for improved performance, codes with 95 physical qubits have been demonstrated [113].

An alternative error correction code is the quantum low density parity check (qLDPC) which has been shown to require a smaller physical qubit overhead to implement. Analysis suggests the qLDPC codes can be implemented on 12 logical qubits using 300 physical qubits, whereas surface codes would require 4000 physical qubits [114]. However, qLDPC codes require long distance connections between physical qubits to implement, which may be difficult to realise at scale.

Progress towards implementing error correction on neutral atom arrays has seen some recent advancements. Mid-circuit readout of qubits has been demonstrated using microwave pulses to shelve data qubits while ancilla qubits are read out [13]; shielding beams have also been used to shift atoms in storage out of resonance with readout light within zone based architectures [7]. Alkaline atoms have presented a pathway for error detecting by converting Rydberg decay and leakage into erasure errors³ by continually monitoring the ground state [115], as well as mid-circuit readout and reset operations [116, 117]. A novel approach of converting qubit state to position allows for differentiation between a given state and loss [7, 118]. Fault tolerant operations have also been demonstrated on logical qubits [119] as well as below threshold quantum error correction [7, 120]

³Erasure errors are those occurring with a known location, which can be detected during circuit operation and converted to atom loss [115].

1.3.3 Challenges Facing Neutral Atom Arrays

In the last few sections, we have discussed a number of the benefits of neutral atom arrays as a platform for quantum computing. However, as discussed in Sect. 1.3.2 there are a few key challenges which must be overcome to realise systems that could be considered fault tolerant. Below we will discuss three key challenges which the work presented in this thesis will try to address.

1. Scalability

It has already been discussed that tweezer arrays offer a route to scaling from a handful of atoms to thousands of atoms, however, the fundamental limit to array sizes is governed by how long an atom can be retained in a trap. ODTs tend to have shallow trap depths which means most collisions with background gas in the vacuum will eject said atom from the trap. For a vacuum limited loss regime, the trap lifetime is characterised by an exponential decay constant τ . For N trapped atoms, the value of τ scales with $1/N$. Most vacuum systems to date are kept under vacuum with aid from an ion-pump which limits the single atom lifetime to tens of seconds [5]. Scaling to arrays with millions of atoms quickly runs into lifetimes on the scale of gate pulses. Therefore, to truly utilise quantum computing on neutral atoms at scale, low vacuum pressures have to be reached.

2. Destructive Readout

Neutral atom qubits are typically read out using fluorescence imaging. To do this, historically, imaging light will be applied across the array and all atoms will be read out in parallel. Multiple rounds of quantum error correction require a set of qubits to remain unperturbed while a subset of qubits are read out, and therefore globally destructive read out is unacceptable. A method for reading out a subset of qubits has been suggested in Ref. [13] where data qubits can be selectively shelved in a dark state by applying a local Stark shift to the ancilla qubits. After shelving, global read out light can be applied which will only cause fluorescence of the ancilla qubits. Another method used in Ref. [7] involves shielding a subset of atoms with a laser beam which

imparts a Stark shift on its excited state, pushing the read out transition far from resonance. Both the examples presented in Ref. [13] and Ref. [7] struggle to avoid of Stark shifting light leaking to atoms which are not meant to be affected. This leads to imperfect shielding and shelving which will cause unwanted interactions with some qubits. These methods will also have trouble scaling to tens of thousands of qubits.

3. Imperfect Gate Operations

Quantum error correction requires physical qubit gate fidelities to be suitably low before protocols start to realise fault tolerance. One of the biggest limiting factors for higher gate fidelities is imperfect blockade during a gate sequence due to the finite strength of the Rydberg-atom interaction. Interaction regimes known as Förster resonances can resonantly amplify interactions to be stronger than the second-order van der Waals coupling typically obtained in platforms with a single atomic species.

1.4 The QuERy Project

The three challenges laid out in Sect. 1.3.3 are the motivation behind the project presented in this thesis. Internally, this project is referred to as Quantum Error Correction on Rydberg arrays (QuERy) and will take a two pronged approach to address the three challenges.

To address scalability, the QuERy project will implement a cryogenic stage to the vacuum system. This is to take advantage of the effects of cryopumping which occur when background vacuum impurities are frozen-out by cryogenic surfaces at temperatures $\lesssim 30$ K. Recent work has demonstrated the use of a cryostat with tweezer array experiments for increased vacuum lifetimes [9, 121–125] with the longest trap lifetimes recording ~ 6000 s [9]. Additionally, inspired by recent results featuring arrays with > 6000 atoms [8], the QuERy project will implement a titanium sublimation pump (TSP) to act as an additional vacuum pump for longer trap lifetimes.

Due to naturally non-resonant read out wavelengths for different elements, the QuERy project will implement dual-species arrays of Rb and Cs atoms to tackle the

challenge of cross-talk free read out. From an operational point of view, one atom, say Rb, is used as a data qubit while the other ancilla atom, Cs, is read out and reset without decohering the Rb atom's state [5]. Dual species read out does not require selective shelving of qubits [13], or the application of a shielding beam [7] which would add additional sources of error. Dual-species arrays have also been suggested as a means for faster read out [126]. This method involves using one species for data and ancilla qubit computation and, performing read out by transferring the ancilla qubit state onto multiple read out atoms. The collective signal of the read out atoms is used to discern the ancilla qubit states. Experimentally, dual-species arrays of Rb and Cs atoms have been demonstrated [10, 127–130] as well as dual-isotope arrays of ^{87}Rb and ^{85}Rb atoms [131].

To address the final challenge, imperfect gate operation, the QuERy project will take advantage of stronger Rydberg-atom interactions which are more easily addressed in dual-species systems. Specifically, we will address Rydberg states which possess Förster resonances between Rb and Cs atoms [132]. This will be performed by addressing Rydberg d -orbitals which possess enhanced shifts for interspecies interactions while suppressing intraspecies couplings [20]. The identified states display larger energy shifts than dual-species Rydberg s -orbitals previously identified and demonstrated in other dual-species systems [10, 130, 132].

1.5 Thesis Outline

In the previous section, we discussed the overarching aims which guide the QuERy project. This section will outline the structure of this thesis and give an indication of the progress towards realising the project aims. At the time of writing, the QuERy experiment has demonstrated single Cs atoms trapped in arrays of ODTs and has made some steps towards realising single Rb atoms in tweezer traps. Additionally, a set of Rydberg states of particular interest have been identified for use in this system. Work towards these results, as well as research into candidate Rydberg states for gate operations, are spread across the four parts to this thesis, each of which are outlined

below.

In Part I, we will discuss the Rydberg-atom interaction and its utility in realising multi-qubit gates. We will present research we published in Ref. [20] which identifies Förster resonant pair states in Rydberg d -orbitals as candidates for high fidelity multi-qubit gates. Part II discusses the MOT laser setup and vacuum system including the design of the cryogenic stage used for this experiment. In Part III, we discuss the experimental realisation of a Cryogenic Cs MOT and the loading single atoms into tweezer arrays. We go on to perform characterisation of the vacuum system by presenting a series of trap lifetime measurements and show how improved vacuum pressures can be achieved by optimising use of the cryostat and TSP. Finally, in Part IV we discuss the progress towards realisation of a dual-species array of single of Rb and Cs atoms. We also discuss work towards the Rydberg laser system which will be used to investigate Förster resonance between Rb and Cs atoms in d -orbitals.

1.6 Work Derived From This Thesis

- P. M. Ireland, *et al.*, *Interspecies Förster resonances for Rb-Cs Rydberg d -states for enhanced multi-qubit gate fidelities*, Phys. Rev. Research **6**, 013293 (2024)
- P. M. Ireland, *et al.*, *A cryogenic dual-species tweezer array*, [In preparation]

Chapter 2

Theoretical Background

This chapter will set out some of the foundational atomic physics required for understanding neutral atom tweezer arrays. We will begin by discussing how we model the alkali atoms before introducing essential topics that are applicable in the rest of this thesis such as light-matter interactions and optical dipole traps (ODTs).

2.1 Alkali Atoms

The hydrogen atom has analytical solutions to the Schrödinger equation under the approximation that the motion of the nucleus can be ignored due to its mass being three orders of magnitude larger than the electron's. This approximation provides a formula — known as the Bohr formula — which predicts the energy of a given electron orbital by considering only the principal quantum number n [133]. This approach cannot be extended to multi-electron atoms due to the complexities of the electrostatic terms. However, the alkalis can be treated as hydrogenic in structure as they consist of a single valance electron with all other electrons tightly bound to the nucleus. The application of a quantum defect $\delta_{n,\ell,j}$ to the alkali's n value is used to encapsulate the screening of nuclear charge provided by the strongly bound inner electrons. This allows for the formulation of a modified Bohr formula [133, 134]

$$E_{n\ell j} = -hc \frac{R^*}{(n - \delta_{n\ell j})^2}, \quad (2.1)$$

where $E_{n\ell j}$ is the energy of the state with corresponding quantum numbers $\{n, \ell, j\}$, h is plank's constant, c is the speed of light and R^* is the mass-corrected Rydberg constant which is given by

$$R^* = R_\infty \frac{M}{M + m_e}, \quad (2.2)$$

where R_∞ is the Rydberg constant, M is the mass of the nucleus and m_e is the mass of the electron. From this we can define an effective principal quantum number $n^* = n - \delta_{n\ell j}$ where the quantum defect can be calculated with the expansion

$$\delta_{n\ell j} = \delta_{0\ell j} + \sum_{k=2}^{\text{even } k} \frac{\delta_{k\ell j}}{(n - \delta_{k\ell j})^k}, \quad (2.3)$$

where the coefficients are fitted to energies obtained from measurements of the transition wavelengths [135, 136].

2.1.1 The Electric Dipole Transitions

Transitions between different electron energy levels are dependent on the electric dipole operator [137]

$$\mu_q = e\mathbf{r} \cdot \hat{\mathbf{e}}_q, \quad (2.4)$$

where $\mathbf{r}$ is the radial displacement of the electron and $\hat{\mathbf{e}}_q$ is the unit vector which describes the polarisation of the electric field. Eq. 2.4 is presented in the spherical basis for which $q \in \{0, \pm 1\}$ corresponding to $\{\pi, \sigma^\pm\}$ polarisations. The couplings between two atomic energy levels $|n, \ell, j, F, m_f\rangle$ and $|n', \ell', j', F', m'_f\rangle$ can be decomposed by considering the overlap of the radial and spherical wave functions with the dipole operator, with the dipole matrix element $\langle n, \ell, j, F, m_f | \mu_q | n', \ell', j', F', m'_f \rangle$ representing the coupling strength [137].

Table 2.1 provides the transition strength in caesium between a combination of low lying states and high n Rydberg states calculated with Alkali.ne Rydberg Calculator (ARC) [138]. The large overlap of the radial wave functions of high n states results in comparatively strong transition strengths between Rydberg states (evident by comparison of transition 1 and transition 4 in Table 2.1). The opposite is true for transitions

Transition	State 1	State 2	Transition strength (ea_o)	Transition wavelength
1	$6s_{1/2}$	$7p_{1/2}$	0.113	459 nm
2	$6s_{1/2}$	$60p_{3/2}$	0.002	319 nm
3	$7p_{1/2}$	$60s_{1/2}$	-0.008	1041 nm
4	$60s_{1/2}$	$60p_{3/2}$	1680	16.5 nm

Table 2.1: Transition strengths between state 1 and state 2 obtained by calculating the dipole matrix element using ARC assuming $q = 0$. The wavelengths for the optical transitions are also provided.

between low n states and Rydberg states (see transition 2 in Table 2.1) which makes direct Rydberg excitation less attractive. Additionally, direct Rydberg excitation from the ground state requires the use of ultraviolet (UV) wavelengths for which laser power becomes a limiting factor. These issues motivates the use of two-photon excitation schemes which are discussed in Sect. 2.2.4.

2.1.2 Zeeman Effect

In the presence of an external magnetic field, the energy levels of an atom will experience a shift from their unperturbed values born out in the previous section. This interaction is known as the Zeeman effect and is an important tool for manipulating atoms. The Zeeman effect can be studied in the uncoupled, fine-structure and hyperfine structure basis independently, but for the ground and low energy excited states of the alkali atoms, the energy shifts from DC magnetic fields is far smaller than the hyperfine splitting. As such, we will only consider the Zeeman effect in the hyperfine-basis.

The Zeeman Hamiltonian is defined in terms of a static magnetic field $\mathbf{B} = B_0 \hat{\mathbf{z}}$ and the total angular momentum vector $\mathbf{F}$ in the hyperfine-basis [133]

$$\hat{H}_B = g_F \mu_B \mathbf{F} \cdot \mathbf{B} = g_F \mu_B F_z B_0, \quad (2.5)$$

where μ_B is the Bohr magneton and g_F is the fine structure Landé g -factor. In the weak-field regime, energy shifts can be considered via perturbation theory with linear splitting of the first order

$$E = g_F m_f \mu_B B. \quad (2.6)$$

For the strong-field regime, perturbation theory no longer holds and instead interactions must be investigated via direct matrix diagonalisation. The work presented in this thesis will only operate in the weak-field regime, and as such the strong-field regime will not be considered further.

The magnetic sub-level quantum number m_f can take values of $-F \leq m_f \leq F$. The corresponding allowed transitions are $\Delta m_f = \pm 1, 0$ except for the $m_f = 0 \rightarrow m'_f = 0$ transition which is forbidden when $\Delta F = 0$ — this is important to consider for optical pumping as discussed in Sect. 2.2.3

2.2 Semi-Classical Light Matter Interaction

For modelling neutral atom qubits, it is typical to only consider a small number of energy states. Ideally, all interactions can be represented as a two-level system with couplings to other states treated as dispersion terms in the interactions. This is because modelling large systems is computationally expensive so simplified models which encapsulate larger systems are a valuable resource. For certain qubit operations (such as Rydberg excitation and Raman transitions), it is common for interactions to occur across three energy levels, but with population effectively zero in one of them allowing the system to be approximated as a two-level system.

In this section, we begin by discussing a classical light field interaction with a two-level atom. We will go on to introduce the concept of spontaneous emission which is important for operating systems such as magneto-optical traps (MOTs) and for optical pumping which is discussed in Sect. 2.2.3. This section will also discuss three-level light-matter interactions which are used in Rydberg excitation for utilising more favourable wavelengths.

2.2.1 Two-Level Systems

We begin by considering a classical light field interacting with a quantised two-level atom [139], as shown schematically in Fig. 2.1a. This system can be represented by the

Chapter 2. Theoretical Background

Hamiltonian

$$\hat{H} = \hat{H}_A + \hat{H}_I, \quad (2.7)$$

where $\hat{H}_A$ is the atom Hamiltonian and $\hat{H}_I$ is the interaction Hamiltonian. The atom Hamiltonian consists of a ground state $|g\rangle$ and an excited state $|e\rangle$ with eigenenergies

$$\hat{H}_A |g\rangle = E_g |g\rangle, \quad (2.8a)$$

$$\hat{H}_A |e\rangle = E_e |e\rangle, \quad (2.8b)$$

where $E_e - E_g = \hbar\omega_A$ is the transition energy spacing with angular frequency ω_A . We can represent the atom Hamiltonian as

$$\hat{H}_A = \hbar\omega_A |e\rangle \langle e|. \quad (2.9)$$

The interaction Hamiltonian is defined by the interaction between a laser and an electric dipole transition within the atom

$$\hat{H}_I = -\boldsymbol{\mu} \cdot \boldsymbol{\mathcal{E}}(t), \quad (2.10)$$

where $\boldsymbol{\mu}$ is the electric dipole moment and $\boldsymbol{\mathcal{E}}(t)$ is the electric field component of the laser. The oscillating field can be expressed as

$$\boldsymbol{\mathcal{E}}(t) = \frac{\mathcal{E}_0}{2} \left[e^{-i\omega_f t} + \text{c.c.} \right] \hat{\boldsymbol{e}}, \quad (2.11)$$

where $\mathcal{E}_0$ is the amplitude of the electric field amplitude, ω_f is the angular frequency of the light and $\hat{\boldsymbol{e}}$ is the unit vector of the electric field component of the light. The interaction Hamiltonian can be rewritten in the atomic basis where we use the rotating wave approximation [140] to ignore the counter propagating terms of the electric field

$$\hat{H}_I = \hbar \frac{\Omega}{2} \left[e^{-i\omega_f t} |g\rangle \langle e| + \text{h.c.} \right], \quad (2.12)$$

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where the Rabi frequency is defined as

$$\Omega = -\frac{\langle e|\boldsymbol{\mu} \cdot \hat{\mathbf{e}}|g\rangle}{\hbar} \mathcal{E}_0. \quad (2.13)$$

Eq. 2.7 can be solved with the Schrödinger equation

$$i\hbar \frac{\partial}{\partial t} |\psi(t)\rangle = \hat{H} |\Psi(t)\rangle, \quad (2.14)$$

for which the general solution for a two-level atom is

$$|\Psi(t)\rangle = \alpha_g(t) |g\rangle + \alpha_e(t) |e\rangle, \quad (2.15)$$

where the coefficients $\alpha_g(t)$ and $\alpha_e(t)$ complex amplitudes. The total light-matter Hamiltonian (Eq. 2.7), the general solution (Eq. 2.15) and the Schrödinger equation (Eq. 2.14) can be combined to represent the time dependent coefficients $\alpha_g(t)$ and $\alpha_e(t)$ as coupled partial differential equations,

$$\dot{\alpha}_g(t) = i\frac{\Omega}{2} e^{i\delta t} \alpha_e(t), \quad (2.16a)$$

$$\dot{\alpha}_e(t) = i\frac{\Omega}{2} e^{-i\delta t} \alpha_g(t), \quad (2.16b)$$

where $\delta = \omega_f - \omega_A$ is the detuning. The exponential terms can be removed by substituting the following transformations into Eq. 2.16

$$\tilde{\alpha}_g(t) = e^{-i\delta t/2} \alpha_g(t), \quad (2.17a)$$

$$\tilde{\alpha}_e(t) = e^{i\delta t/2} \alpha_e(t), \quad (2.17b)$$

which yields,

$$\dot{\tilde{\alpha}}_g(t) = -\frac{i}{2} \left[\delta \tilde{\alpha}_g(t) - \Omega \tilde{\alpha}_e(t) \right], \quad (2.18a)$$

$$\dot{\tilde{\alpha}}_e(t) = \frac{i}{2} \left[\delta \tilde{\alpha}_e(t) + \Omega \tilde{\alpha}_g(t) \right], \quad (2.18b)$$

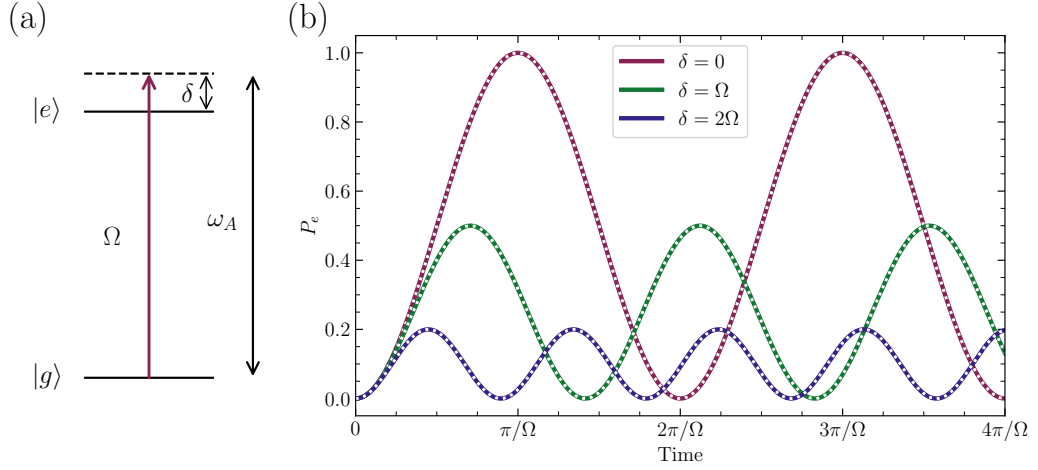


Figure 2.1: (a) Schematic of a two-level atom interacting with a classical light field. (b) Coloured lines indicate Rabi oscillations obtained through numerical QuTiP modelling of the Hamiltonian in Eq. 2.19 for different values of δ ; the white dashes trace the analytical solution to the Hamiltonian presented in Eq. 2.20.

which allows the rewriting of the Hamiltonian from Eq. 2.7 in the convenient form

$$\frac{\hat{H}}{\hbar} = \frac{\Omega}{2} [|g\rangle \langle e| + \text{h.c.}] + \delta |e\rangle \langle e|. \quad (2.19)$$

Solving the Schrödinger equation for the updated Hamiltonian of Eq. 2.19 and assuming an initial condition of $|\Psi(0)\rangle = |g\rangle$ yields the time dependent population of the excited state [139]

$$P_e = |\langle e | \psi(t) \rangle|^2 = |\tilde{\alpha}_e(t)|^2 = \frac{\Omega^2}{\tilde{\Omega}^2} \sin^2 \left(\frac{\tilde{\Omega} t}{2} \right), \quad (2.20)$$

where $\tilde{\Omega} = \sqrt{\Omega^2 + \delta^2}$ is the modified Rabi frequency. Eq. 2.20 demonstrates the realisation of Rabi oscillations in a semi-classical model.

To demonstrate the validity of this solution, Fig. 2.1b shows the results of a numerical QuTiP [141] simulation which solves the Schrödinger equation for the Hamiltonian in Eq. 2.19 where different line colours represent different values of δ . The white marks in the figure trace out the analytical solution presented by Eq. 2.20, which shows good agreement with QuTiP. There are two notable consequences when considering values of $\delta \neq 0$. First is the emergence of dampened Rabi oscillations as increased values of

$|\delta|$ act to reduce the amplitude as presented in Eq. 2.20. Second is the increase in Rabi frequency for larger $|\delta|$, which is also clear from simple inspection of Eq. 2.20.

Above we have demonstrated the emergence of Rabi oscillations due to a classical light field interacting with a quantised two-level atom. This model has a couple of limitations as it does not predict some phenomena which is observed during real light-matter interactions — such as spontaneous decay.

2.2.2 Spontaneous Emission

The light-matter interactions described in the previous section only consider coherent evolution between atomic energy levels without consideration of spontaneous decay due to finite state lifetimes. The effects of spontaneous decay are exploited throughout atomic physics such as for the operation of MOTs and for optical pumping — as discussed in the next section. Here we represent spontaneous decay for a two-level system using the Lindblad formalism as described in Ref. [142].

So far, we have described quantum systems as pure states¹, however, to incorporate decoherence, we now consider the density operator

$$\hat{\rho} = \sum_i p_i |\Psi\rangle \langle \Psi|, \quad (2.21)$$

where $\{|\Psi_i\rangle\}$ is the set of pure states with the associated set of probabilities $\{p_i\}$. The diagonal terms of $\hat{\rho}$ are the basis states populations such that $\text{Tr}[\hat{\rho}] = 1$ and the off-diagonals describe the coherence between basis states.

The time evolution of $\hat{\rho}$ with Hamiltonian $\hat{H}$ can be modelled with by the Lindblad master equation

$$\frac{d}{dt}\hat{\rho} = \frac{i}{\hbar} [\hat{\rho}, \hat{H}] + \mathcal{L}(\hat{\rho}), \quad (2.22)$$

where the Lindblad operator is defined as

$$\mathcal{L}(\hat{\rho}) = \sum_i \left[\hat{L}_i \hat{\rho} \hat{L}_i^\dagger - \frac{1}{2} \left(\hat{L}_i^\dagger \hat{L}_i \hat{\rho} + \hat{\rho} \hat{L}_i^\dagger \hat{L}_i \right) \right], \quad (2.23)$$

¹States which are entirely coherent and can be represented in the form $|\Psi\rangle = \sum_i c_i |i\rangle$.

where $\hat{L}_i$ are the collapse operators. For spontaneous emission in a two-level system, $\hat{L} = \sqrt{\Gamma_e} |g\rangle \langle e|$ where Γ_e is the linewidth of the transition. Solving the Lindblad master equation for spontaneous decay allows for the definition of the optical Bloch Equations (OBEs) [143]

$$\begin{aligned}\dot{\rho}_{gg} &= +\Gamma_e \rho_{ee} + i\frac{\Omega}{2} [\tilde{\rho}_{eg} - \tilde{\rho}_{ge}], \\ \dot{\rho}_{ee} &= -\Gamma_e \rho_{ee} + i\frac{\Omega}{2} [\tilde{\rho}_{ge} - \tilde{\rho}_{eg}], \\ \dot{\tilde{\rho}}_{ge} &= -\left(\frac{\Gamma_e}{2} + i\delta\right) \tilde{\rho}_{ge} + i\frac{\Omega}{2} [\rho_{ee} - \rho_{gg}], \\ \dot{\tilde{\rho}}_{eg} &= -\left(\frac{\Gamma_e}{2} - i\delta\right) \tilde{\rho}_{eg} + i\frac{\Omega}{2} [\rho_{gg} - \rho_{ee}],\end{aligned}\tag{2.24}$$

where $\tilde{\rho}_{eg} \equiv \rho_{eg} e^{+i\delta t}$ and $\tilde{\rho}_{ge} \equiv \rho_{ge} e^{-i\delta t}$ removes the time dependent terms. The steady-state solution can be obtained by setting the time derivative terms of Eq. 2.24 to zero

$$\rho_{ee}^{ss} = \frac{1}{2} \frac{s}{1 + s + (\frac{2\delta}{\Gamma_e})^2},\tag{2.25}$$

where $s = I/I_{\text{sat}}$ is the saturation parameter with laser intensity I and atomic saturation intensity [79]

$$I_{\text{sat}} = \frac{2\pi^2 \hbar \Gamma_e c}{3\lambda^3}.\tag{2.26}$$

The Rabi frequency can be rewritten in terms of the transition linewidth and saturation parameter [79]

$$\Omega = \Gamma \sqrt{\frac{s}{2}}.\tag{2.27}$$

To demonstrate the effect spontaneous decay has on the evolution of a two-level system, Fig. 2.2b shows a QuTiP simulation of the Lindblad master equations for the two-level Hamiltonian of Eq. 2.19 where $\delta = 0$ and $\Gamma_e/(2\pi) = 5.2$ MHz corresponding to a resonant laser driving of the D₂ transition in Cs. The value of Ω was set such that $s = 10$, as per Eq. 2.27. The results demonstrate the emergence of damped Rabi oscillations, occurring due to decoherence caused by spontaneous decay from the excited state to the ground state. This leads to the steady-state solution as predicted by Eq. 2.25. Damped Rabi oscillations are an incoherent scattering process which is the regime in which MOTs and florescent imaging of single atoms operate. For these

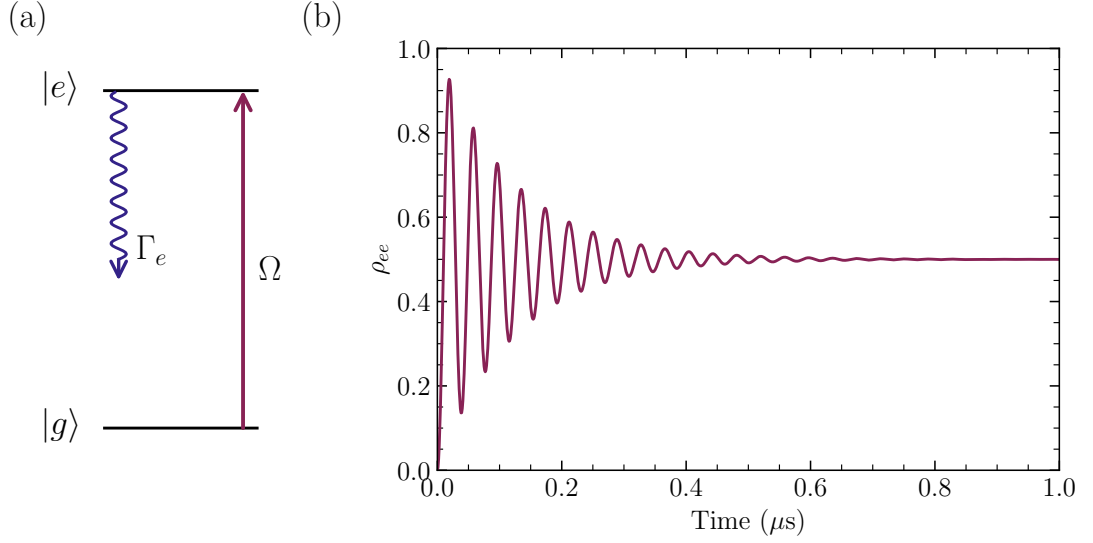


Figure 2.2: (a) Two-level light-matter interaction with spontaneous decay from $|e\rangle$ to $|g\rangle$ with atomic linewidth Γ_e . (b) Resonant driving of the Cs D₂ transition for the modelled with QuTiP demonstrating the emergence of dampened Rabi oscillations.

processes, the scattering rate

$$\Gamma_{\text{sc}} = \Gamma_e \rho_{ee}^{ss} = \frac{\Gamma_e}{2} \frac{s}{1 + s + (\frac{2\delta}{\Gamma_e})^2}, \quad (2.28)$$

is an important quantity, as it can be used to obtain a MOT atom number, as was done in Chapter 7. The decoherence of a two-level system is also important when considering quantum gates as the finite state lifetime ($\tau = 1/\Gamma_e$) puts a limit on how slowly gate sequences can be performed. This motivates the study of strong Förster interactions discussed in Chapter 4.

2.2.3 Qubit Preparation

For quantum computing with neutral alkali atoms it is typical to use the hyperfine clock states² for the computational basis $\{|0\rangle, |1\rangle\}$, as discussed in Chapter 1. As a result of scattering during the cooling stage, all atoms should be in the $|1\rangle$ state by the

²The $F = 3$ and $F = 4$ hyperfine states of the $6s_{1/2}$ level in ^{133}Cs , and the $F = 1$ and $F = 2$ of the $5s_{1/2}$ state of ^{87}Rb .

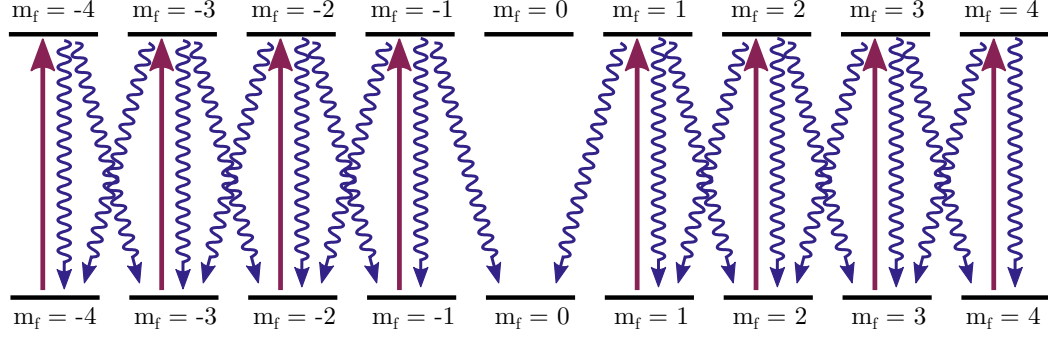


Figure 2.3: Optical pumping between the Cs $F = 4$ ground state and $6p F' = 4$ excited state. A laser (red arrows) with π -polarisation pumps atoms to the excited state where atoms decay back to the ground state (purple arrows). As the $m_f = 0 \rightarrow m'_f = 0$ transition is forbidden (when $\Delta F = 0$), the population builds up in the $m_f = 0$ state.

time they are loaded into tweezers, however, the population will be spread across the m_f magnetic sub-levels. To ensure the qubits are initiated in a well defined state, we must employ optical pumping to bring the atoms to a single m_f level.

Consider the first order Zeeman shifts presented in Eq. 2.6 where the $m_f = 0$ sub-level has two important characteristics. Firstly, the Zeeman shifts vanish under the first order approximation, so consequently, the $m_f = 0$ sub-level is not susceptible to magnetic noise. Secondly, the $m_f = 0 \rightarrow m'_f = 0$ transition is forbidden, which can be exploited when performing optical pumping. As such the $m_f = 0$ magnetic sub-level is an ideal state to use for the computational basis.

To initialise qubits in the $|1\rangle$ state — which is the $F = 4$ hyperfine state in Cs — we exploit the incoherent scattering process described in the previous section to perform optical pumping on one of the D-lines. For illustration, we will consider optical pumping on the $F = 4 \rightarrow F' = 4$ transition which exists on both the D₁ and D₂ lines in Cs [79]. Light with π -polarisation is applied on the Cs qubits which drives $m_f \rightarrow (m'_f = m_f)$ transition except for atoms in the $m_f = 0$ sub-level which remain unperturbed during optical pumping. The excited m_f sub-levels will scatter on $m'_f \rightarrow (m_f = m'_f \pm 1, 0)$ transitions as shown in Fig. 2.3. Over repeated scattering, the atom population will build up in the $m_f = 0$ state where excitation to the F' state no longer takes place. Optical pumping like this requires pure π -polarisation which means it must be performed in the presence of a magnetic basis field.

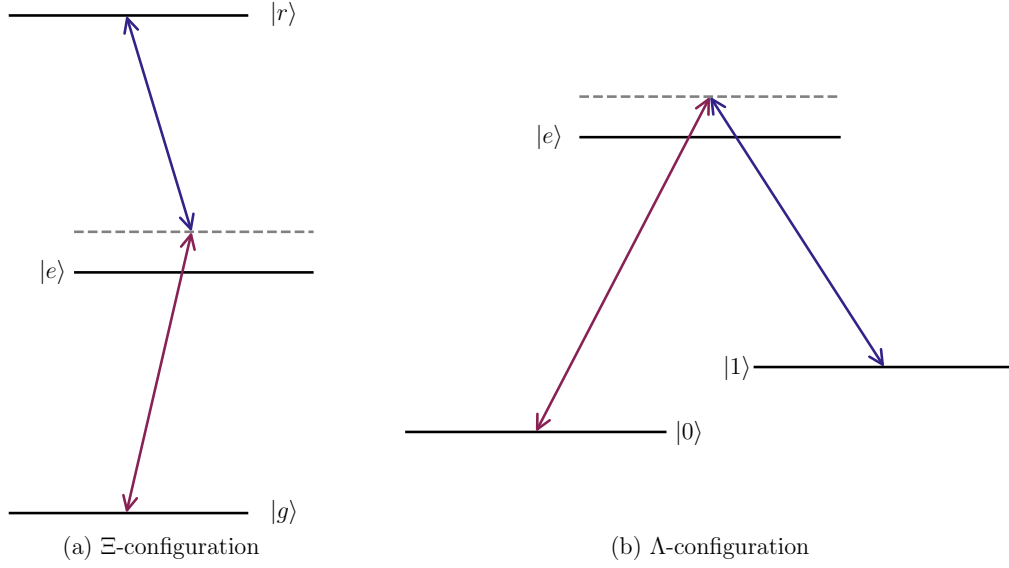


Figure 2.4: Three-level transition schemes. (a) The Ξ -configuration where excitation is performed from a low energy state $|g\rangle$ to a high energy state $|r\rangle$ via an intermediate excited state $|e\rangle$. (b) The Λ -configuration where excitation is performed from a low energy state $|0\rangle$ to a target state $|1\rangle$ via a higher energy state $|e\rangle$.

2.2.4 Three-Level Systems

As discussed in Sect. 2.1.1, the matrix elements for directly addressing Rydberg states with single photon excitation are weak due to the small overlap with the radial wave function between the ground and Rydberg states. Additionally, Table 2.1 shows that these transitions require lasers that operate at UV wavelengths which typically have lower power when compared to infrared laser systems. This motivates Rydberg excitation through a two-photon scheme where the system evolves from the ground state to the excited state via an intermediate state. Two-photon excitation is performed using a pair of lasers whose combined frequency matches the energy difference between the ground and target Rydberg state, but with a large detuning relative to an intermediate excited state. Under these conditions, the population can be coherently transferred to the Rydberg state with minimal population in the excited state, which suppresses decoherence due to scattering. Under these conditions, we can consider a Cs two-photon excitation scheme where, using the data from Table 2.1, we can see the $6s_{1/2} \rightarrow 7p_{1/2}$

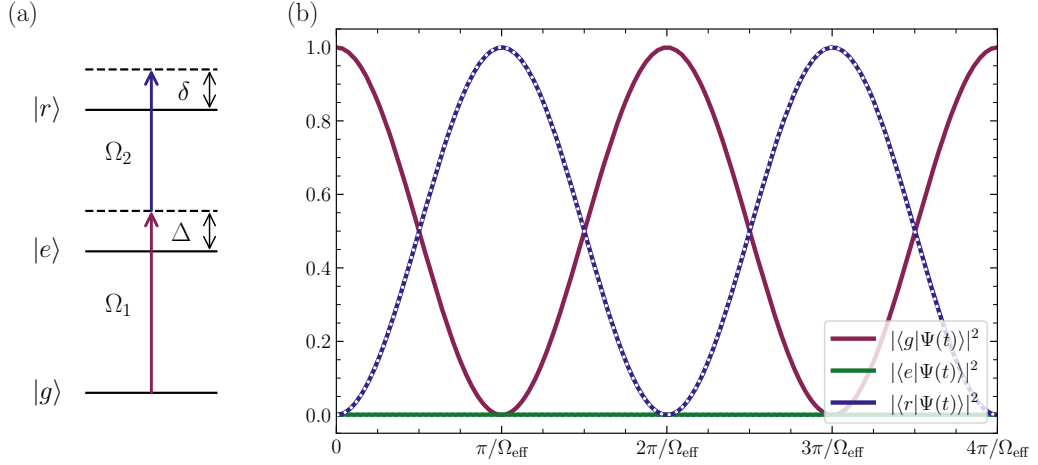


Figure 2.5: (a) Schematic of a quantised three-level atom interacting with a classical light field. (b) Rabi oscillations obtained through numerical QuTiP modelling for where $\Delta \gg \Omega_1, \Omega_2$ and $\Omega_1 = \Omega_2$ which show coherent Rabi oscillations between $|g\rangle$ and $|r\rangle$ with suppressed population in the $|e\rangle$ state; the white dashed line traces the solution obtained through adiabatic elimination presented in Eq. 2.43. The agreement between the numerical and analytical solution demonstrates the validity of adiabatic elimination.

transition has a comparatively strong transition strength at 459 nm, for which we can use a blue laser with lower power. The second photon transition can be performed on the $7p_{1/2} \rightarrow 60s_{1/2}$ which has a transition wavelength at 1041 nm for which higher laser power is available making up for the comparatively small dipole coupling strength. Selection rules of optical dipole transitions require a change in orbital angular momentum ℓ of ± 1 and as such, two-photon excitation allows the addressing of Rydberg s - and d -orbitals which are discussed further in Chapter 4.

In this subsection, we discuss the derivation of the effective two-level Hamiltonian by first considering the full three-level system before adiabatically eliminating the intermediate state. This analysis is not only valid for Rydberg excitation — more generally referred to as the ladder or Ξ -configuration as shown in Fig. 2.4a — but also for Raman transitions — often referred to as the Λ -configuration as shown in Fig. 2.4b.

Consider a three-level system where a ground state $|g\rangle$ with energy $\hbar\omega_g$ is dipole-coupled to an intermediate excited state $|e\rangle$ with energy $\hbar\omega_e$ driven by a laser with angular frequency ω_1 ; $|e\rangle$ is dipole-coupled to the Rydberg state $|r\rangle$ with energy $\hbar\omega_r$,

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driven by a second laser with angular frequency ω_2 . This system is depicted in Fig. 2.5a where the Hamiltonian can be represented as

$$\hat{H} = \hat{H}_A + \hat{H}_I, \quad (2.29)$$

where $\hat{H}_A$ is the atom Hamiltonian and $\hat{H}_I$ is the interaction Hamiltonian both of which can be treated in a similar manner to the two-level system discussed in Sect. 2.2.1. For the three-level atom, we can represent the interaction term under the rotating wave approximation for both light fields with

$$\frac{\hat{H}_I}{\hbar} = \frac{\Omega_1}{2} \left[e^{-i\omega_1 t} |g\rangle \langle e| + \text{h.c.} \right] + \frac{\Omega_2}{2} \left[e^{-i\omega_2 t} |e\rangle \langle r| + \text{h.c.} \right], \quad (2.30)$$

where Ω_χ is the single photon Rabi frequencies for the $|g\rangle \rightarrow |e\rangle$ transition when $\chi = 1$ and $|e\rangle \rightarrow |r\rangle$ transition when $\chi = 2$. The Rabi frequencies can be expressed in terms of the transition dipole moments $\boldsymbol{\mu}_\chi$ and the amplitude of the oscillating electric field $\mathcal{E}_\chi$ along the fields unit vector $\hat{\mathbf{e}}_\chi$ as

$$\Omega_1 = -\frac{\boldsymbol{\mu}_1 \cdot \hat{\mathbf{e}}_1}{\hbar} \mathcal{E}_1, \quad \Omega_2 = -\frac{\boldsymbol{\mu}_2 \cdot \hat{\mathbf{e}}_2}{\hbar} \mathcal{E}_2. \quad (2.31)$$

The general solution to the Schrödinger equation (Eq. 2.14) for this three-level system is

$$|\Psi(t)\rangle = \alpha_g(t) |g\rangle + \alpha_e(t) |e\rangle + \alpha_r(t) |r\rangle. \quad (2.32)$$

The coefficients in Eq. 2.32 can be represented as the coupled differential equations

$$\dot{\alpha}_g(t) = -i\omega_g \alpha_g(t) - i\frac{\Omega_1}{2} e^{i\omega_1 t} \alpha_e(t), \quad (2.33a)$$

$$\dot{\alpha}_e(t) = -i\omega_e \alpha_e(t) - i\frac{\Omega_1}{2} e^{i\omega_1 t} \alpha_g(t) - i\frac{\Omega_2}{2} e^{i\omega_2 t} \alpha_r(t), \quad (2.33b)$$

$$\dot{\alpha}_r(t) = -i\omega_r \alpha_r(t) - i\frac{\Omega_2}{2} e^{i\omega_2 t} \alpha_e(t). \quad (2.33c)$$

These can be rewritten to remove the time dependent terms with the transforms $\alpha_g(t) =$

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$\tilde{\alpha}_g(t)e^{i(\delta/2-\omega_g)t}$, $\alpha_e(t) = \tilde{\alpha}_e(t)e^{i(\delta/2-\omega_g-\omega_1)t}$, $\alpha_r(t) = \tilde{\alpha}_r(t)e^{i(\delta/2-\omega_g-\omega_1-\omega_2)t}$ to yield

$$\dot{\tilde{\alpha}}_g = -i\frac{\delta}{2}\tilde{\alpha}_g - i\frac{\Omega_1}{2}\tilde{\alpha}_e, \quad (2.34a)$$

$$\dot{\tilde{\alpha}}_e = +i\Delta\tilde{\alpha}_e - i\frac{\Omega_1}{2}\tilde{\alpha}_g - i\frac{\Omega_2}{2}\tilde{\alpha}_r, \quad (2.34b)$$

$$\dot{\tilde{\alpha}}_r = +i\frac{\delta}{2}\tilde{\alpha}_r - i\frac{\Omega_2}{2}\tilde{\alpha}_e, \quad (2.34c)$$

where we define $\Delta = \omega_1 - (\omega_g - \omega_e)$ as the first photon detuning and $\delta = \omega_2 - (\omega_e - \omega_r)$ as the second photon detuning, as shown in Fig. 2.5a. We can now rewrite the Hamiltonian in Eq. 2.29 as

$$\frac{\hat{H}}{\hbar} = \frac{\Omega_1}{2} \left[|g\rangle \langle e| + \text{h.c.} \right] + \frac{\Omega_2}{2} \left[|e\rangle \langle r| + \text{h.c.} \right] - \Delta |e\rangle \langle e| - \delta |r\rangle \langle r|, \quad (2.35)$$

where we impart all of δ on the Rydberg state for convenience.

Adiabatic Elimination

For $\Delta \gg \Omega_1, \Omega_2$, it can be assumed that $\dot{\tilde{\alpha}}_e \approx 0$ and Eq. 2.34b can be rewritten as

$$\tilde{\alpha}_e = \frac{\Omega_1\tilde{\alpha}_g + \Omega_2\tilde{\alpha}_r}{2\Delta}, \quad (2.36)$$

which allows the rewriting of Eq. 2.34a and Eq. 2.34c in the $\{|g\rangle, |r\rangle\}$ basis

$$\dot{\tilde{\alpha}}_g = -i \left[\frac{\delta}{2} + \frac{|\Omega_1|^2}{4\Delta} \right] \tilde{\alpha}_g - i \frac{\Omega_1\Omega_2}{4\Delta} \tilde{\alpha}_r, \quad (2.37a)$$

$$\dot{\tilde{\alpha}}_r = +i \left[\frac{\delta}{2} - \frac{|\Omega_2|^2}{4\Delta} \right] \tilde{\alpha}_r - i \frac{\Omega_1\Omega_2}{4\Delta} \tilde{\alpha}_g. \quad (2.37b)$$

In this representation, the differential equations look a lot like a two-level light-matter interaction as described in Sect. 2.2.1. We define an effective Rabi frequency

$$\Omega_{\text{eff}} = \frac{\Omega_1\Omega_2}{2\Delta}, \quad (2.38)$$

and the ground and Rydberg AC stark shift³ terms as

$$\Delta_g^{\text{AC}} = \frac{|\Omega_1|^2}{4\Delta}, \quad (2.39a)$$

$$\Delta_r^{\text{AC}} = \frac{|\Omega_2|^2}{4\Delta}, \quad (2.39b)$$

and an effective detuning as

$$\delta_{\text{eff}} = \delta + \Delta_g^{\text{AC}} - \Delta_r^{\text{AC}}. \quad (2.40)$$

Making the transforms $\tilde{\alpha}_g = \tilde{\alpha}'_g e^{-i(\Delta_{\text{eff}} + 2\Delta_g^{\text{AC}} - \delta)/2}$ and $\tilde{\alpha}_r = \tilde{\alpha}'_r e^{-i(\Delta_{\text{eff}} + 2\Delta_r^{\text{AC}} - \delta)/2}$ results in Eq. 2.37 becoming

$$\dot{\tilde{\alpha}}'_g = -\frac{i}{2} \left[\delta_{\text{eff}} \tilde{\alpha}'_g + \Omega_{\text{eff}} \tilde{\alpha}'_r \right], \quad (2.41a)$$

$$\dot{\tilde{\alpha}}'_r = +\frac{i}{2} \left[\delta_{\text{eff}} \tilde{\alpha}'_r - \Omega_{\text{eff}} \tilde{\alpha}'_g \right]. \quad (2.41b)$$

This allows for Eq. 2.41 can be written as an effective two-level Hamiltonian

$$\frac{\hat{H}_{\text{eff}}}{\hbar} = \frac{\Omega_{\text{eff}}}{2} \left[|g\rangle \langle r| + \text{h.c.} \right] - \delta_{\text{eff}} |r\rangle \langle r|. \quad (2.42)$$

Eq. 2.42 allows a three-level system to be considered under the analysis of a two-level light-matter interaction. We can therefore represent Ξ - and Λ -operations in the form of Eq. 2.20

$$P_r(t) = |\langle r | \Psi(t) \rangle|^2 = \frac{\Omega_{\text{eff}}^2}{\tilde{\Omega}^2} \sin^2 \left(\frac{\tilde{\Omega}}{2} t \right), \quad (2.43)$$

where $\tilde{\Omega} = \sqrt{\Omega_{\text{eff}}^2 + \delta_{\text{eff}}^2}$.

To check the validity of Eq. 2.42, Fig. 2.5b shows a numerical QuTiP model for the full three-level atom using Eq. 2.35 where $\Omega_1 = \Omega_2$ and $\Delta = 20\Omega_2$, and δ is set to ensure resonant two-photon transition; the initial state for this model is $|\Psi(0)\rangle = |g\rangle$. We can see that under these conditions, the $|\Psi(t)\rangle$ coherently evolves from $|g\rangle$ to the $|r\rangle$ with effectively zero population present in the $|e\rangle$ state. The white dotted trace over

³Discussed in more detail below in Sect. 2.3.1.

the $|\langle r|\Psi(t)\rangle|^2$ line in Fig. 2.5b the ground state evolution as predicted by adiabatic elimination of $|r\rangle$ (per Eq. 2.43), and which shows good agreement with the three-level numerical model.

2.3 Optical Dipole Traps (ODTs)

The confinement of atoms with light is a critical component of neutral atom quantum computing. In this section, we discuss how an ODT can be modelled in the two-level atom picture as introduced above. We then move on to discuss multi-level atoms and describe how this brings about the elimination of differential light shifts (referred to as magic trapping) which has important implications for operating tweezer arrays.

2.3.1 The AC Stark Effect

In Sect. 2.2.1, we introduced analysis of a two-level atom interacting with a classical light field which brings about the realisation of Rabi oscillations. Here, we take a different approach by studying the effects a light field has on an atom's energy levels, from which we can extract the AC Stark effect which is the undying principal governing optical tweezers.

Consider the two-level light-matter interaction Hamiltonian in the rotating frame (Eq. 2.19) in matrix form

$$\hat{H} = \hbar \begin{pmatrix} 0 & \Omega/2 \\ \Omega/2 & \delta \end{pmatrix}, \quad (2.44)$$

which has eigenenergies [144]

$$\lambda_{\pm} = \frac{\hbar}{2} \left[\delta \pm \sqrt{\delta^2 + \Omega^2} \right]. \quad (2.45)$$

For the case where $\Delta \gg \Omega$, we can perform a Taylor expansion to the second order on the square root terms in Eq. 2.45 from which we can define [145]

$$\Delta E_{\pm} = \pm \hbar \frac{\Omega^2}{4\delta}, \quad (2.46)$$

which is the AC Stark shift where the upper sign corresponds to the ground state energy shift, and the lower sign corresponds to the excited state energy shift. Fig. 2.6a illustrates the effects of a red detuned laser beam ($\omega < \omega_A$), where the ground and excited state experience opposite energy shifts. In the next subsection, we will see how the AC Stark effect can be utilised to create optical tweezers for trapping atoms.

2.3.2 Optical Tweezers

An optical tweezer can be created at the focus of a Gaussian beam resulting in an intensity at the focus given by

$$I(r, z) = I_0 \left(\frac{w_0}{w(z)} \right)^2 \exp \left[-\frac{2r^2}{w^2(z)} \right], \quad (2.47)$$

where z is the direction of beam propagation, r is the radial axis, $I_0 = \frac{2P_0}{\pi w_0^2}$ is the peak intensity, and $w(z) = w_0 \sqrt{1 + (z/z_R)^2}$ is the beam waist as a function of z with Rayleigh range $z_R = \pi w_0^2 / \lambda$ and focus waist w_0 .

We first consider a simple model of a two-level atom interacting with a far red detuned ODT. In the presence of a tightly focused laser beam, an atom will experience an energy shift from its unperturbed value, as per the AC Stark effect discussed in the previous subsection. This perturbation creates a specially dependent potential $U(r)$ which corresponds to a force $F(r) = -\nabla U(r)$ [146]. From the definition of the Rabi frequency (Eq. 2.13) it is clear that $\Omega^2 \propto \mu^2 |\mathcal{E}_0|^2$ and given that $I \propto |\mathcal{E}_0|^2$, the AC Stark shift for a two-level atom can be rewritten as [146]

$$\Delta E = \pm \frac{3\pi c^2}{2\omega_0^3} \left(\frac{\Gamma}{\delta} \right) I(r), \quad (2.48)$$

where $\Gamma \propto \mu^2$ is the atomic linewidth and w_0 is the beam waist at the focus. The photon scattering rate is given by

$$\Gamma_{\text{sr}}(r) = \frac{3\pi c^2}{2\hbar\omega_0^3} \left(\frac{\Gamma}{\delta} \right)^2 I(r). \quad (2.49)$$

Fig. 2.6b illustrates how the Gaussian beam described in Eq. 2.47 creates a trapping

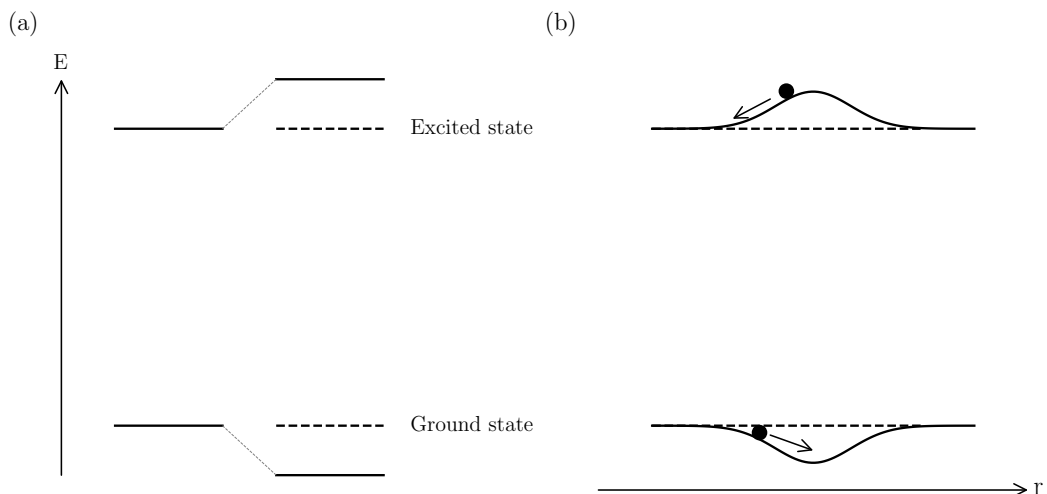


Figure 2.6: Schematic showing the light shifts from a red detuned ($\delta < 0$) ODT for a two-level atom. (a) Shows the energy shifts relative to the unperturbed state. (b) Shows the energy shift as a function of beam radius r revealing a trapping potential for the ground state and an anti-trapping potential for the excited state.

potential in the ground state and an anti-trapping potential in the excited state. During a typical trapping experiment, atoms will spend most their time in the ground state while confined by ODTs and any excitation to the excited state — to perform gates for example — happen in the dark so the anti-trap does warrant concern. However, cooling and imaging must take place with some confinement from the ODTs which will result in ejection of an atom when it is in the excited state under this two-level atom approximation. This necessitates the implementation of out of phase chopping between trap light and cooling/imaging light.

The two-level atom approximation is only valid when the net contribution of Stark shifts is dominated by a single transition, which is rarely the case. For a full picture, we must consider the light shift from all states with dipole couplings. The total Stark shift for a given state ζ can be calculated by summing the AC Stark shift contributions

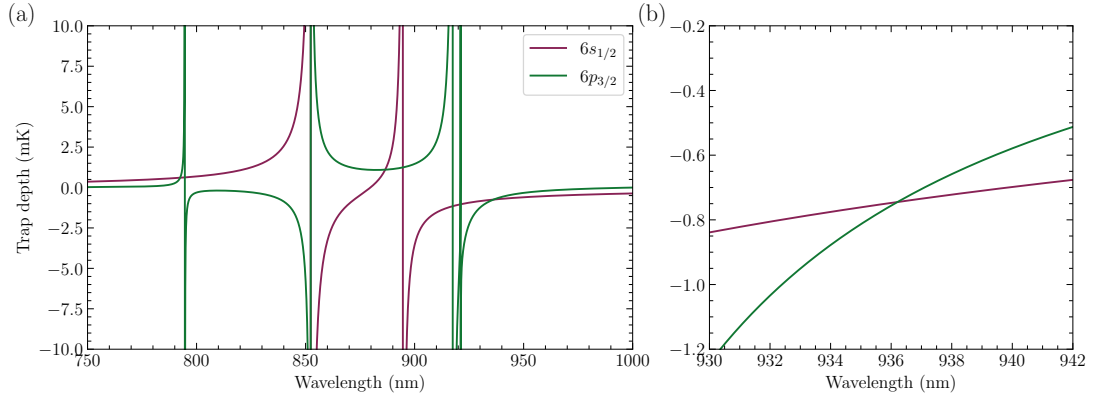


Figure 2.7: (a) The net contributions of the AC Stark shifts for all transitions up to $n = 10$ for the Cs $6s_{1/2}$ and $6p_{3/2}$ levels assuming π -polarised light with a waist of $1 \mu\text{m}$ and a power of 1.8 mW . (b) AC stark shifts highlighting a magic wavelength at $\sim 946 \text{ nm}$ where the differential shifts between the $6s_{1/2}$ and $6p_{3/2}$ states are zero.

over all transitions with the i th states as [147–149]

$$U(r) = \frac{3\pi c^2}{2} \sum_{\zeta \neq i} \frac{A_{\zeta,i}(2j_{\zeta} + 1)}{\omega_{\zeta,i}^3 \delta} \begin{pmatrix} j_{\zeta} & 1 & j_i \\ m_{j_{\zeta}} & q & -m_{j_i} \end{pmatrix}^2 I(r), \quad (2.50)$$

where $\omega_{\zeta,i}$ is the transition frequency for the $\zeta \rightarrow i$ transition and the Einstein A coefficient is defined as [150]

$$A_{\zeta,i} = \frac{\omega_{\zeta,i}^3}{3\pi\epsilon_0\hbar c^3} \frac{|\langle n_{\zeta}, \ell_{\zeta}, j_{\zeta} || \mu_q || n_i, \ell_i, j_i \rangle|^2}{2j_{\zeta} + 1}, \quad (2.51)$$

where $|\langle n_{\zeta}, \ell_{\zeta}, j_{\zeta} || \mu_q || n_i, \ell_i, j_i \rangle|$ is the reduced matrix element [151].

In Fig. 2.7, we compute the light shifts for the Cs $6s_{1/2}$ and $6p_{3/2}$ account for all dipole coupled states up to $n = 10$, calculated using Eq. 2.50. Fig. 2.7a shows the trap depth for states calculated for a dipole beam with $w_0 = 1 \mu\text{m}$ and laser power of 1.8 mW with π -polarisation ($q = 0$). We can see that the differential light shifts do not behave neatly as described in the two-level example which leads to some interesting behaviour. This is highlighted in Fig. 2.7b where at a wavelength of $\sim 936 \text{ nm}$, the differential light shifts between the $6s_{1/2}$ and $6p_{3/2}$ is zero, a phenomena often referred

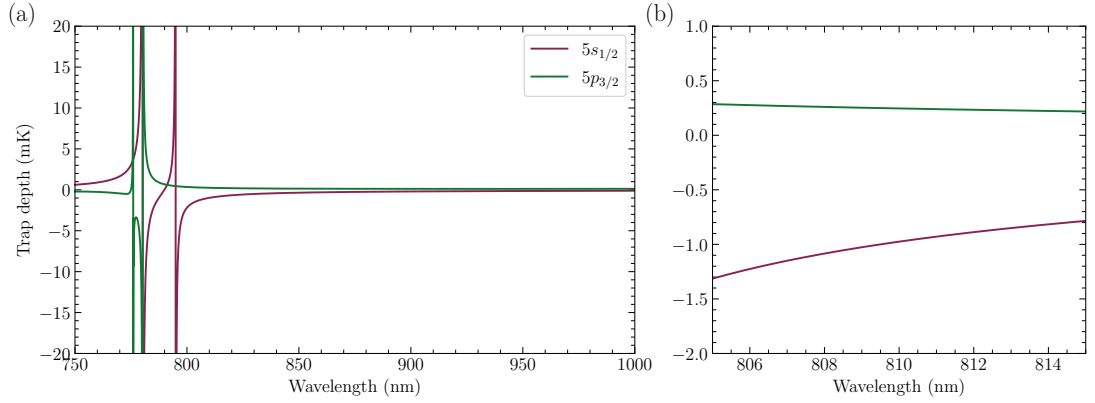


Figure 2.8: (a) The net contributions of the AC Stark shifts for all transitions up to $n = 9$ for the Rb $5s_{1/2}$ and $5p_{3/2}$ levels assuming π -polarised light with a waist of $1 \mu\text{m}$ and a power of 0.3 mW. (b) AC Stark shifts near the target wavelength (810 nm) for Rb trapping in this experiment; shows an attractive potential for the $5s_{1/2}$ state and a repulsive potential for the $5p_{3/2}$ state.

to as magic trapping. Using a magic wavelength has two useful consequences. First, both the ground and excited state are trapped by the ODT. Second, lack of differential shift means that cooling and imaging can be performed without need for out of phase chopping. Ref. [149] directly measures the magic wavelength for the Cs D_2 line on the $F = 4 \rightarrow F' = 5$ transition. They find that the magic wavelength exists between $\sim 933 \text{ nm}$ and $\sim 937 \text{ nm}$ dependent on which m_f level is addressed due to the tensor light shift of the hyperfine states.

We will now consider the same analysis to obtain the light shifts induced within a Rb atoms $5s_{1/2}$ ground state and $5p_{3/2}$ excited state. For this we will assume beam power of 0.3 mW with a $1/e^2$ radius of $1 \mu\text{m}$ and account for the light shift contributions for all dipole coupled states up to $n = 9$. The resulting shifts are shown in Fig. 2.8a where we can see an attractive potential in the ground state for all wavelengths $> 800 \text{ nm}$ while the excited state has a repulsive potential. In comparison to the Cs results in Fig. 2.7a, the ground state is repulsive on the wavelength range of $800 \text{ nm} - 850 \text{ nm}$, which makes this range ideal for species selective trapping of Rb.

For the experiment described in this thesis, we chose 810 nm as the trapping wavelength of Rb. The ground and excited state shifts are shown in Fig. 2.8b for wavelengths

about this value. The Cs trapping wavelength was chosen to be 940 nm which, as discussed in Chapter 8, was found to be near enough magic to perform imaging and cooling without the need for out of phase chopping. As discussed above, this wavelength creates an attractive potential in the Rb ground state, however, Ref. [18] finds that at ~ 940 nm, Rb experiences a roughly three times smaller trap depth when compared to Cs.

2.3.3 Overview

In this chapter, we have discussed some of the key principles utilised in Rydberg tweezer array experiments. We have discussed how high principal quantum number states can be reached via a two-photon excitation scheme, which allows for the use of transitions with advantageous wavelengths and stronger dipole matrix elements. Additionally, we have seen that the two-level atom can explain the concept of ODTs, however, consideration of the full multi-level nature of atoms enables the choice of wavelengths which offer *magic trapping* for efficient in-trap cooling and imaging of atoms. Finally, we have discussed incoherent scattering processes, which are key in many aspects of tweezer arrays experiments, such as qubit preparation and laser cooling.

Part I

Rydberg Interaction

Chapter 3

Rydberg-Atom Interactions

High fidelity multi-qubit gates are one of the key requirements for the realisation of a fault tolerant quantum computer. On neutral atom platforms, qubit entanglement is mediated by the Rydberg interaction which exists in states with high principal quantum numbers where the dipolar nature of neutral atoms becomes very apparent. Rydberg atoms are able to extend their influence to the micron distance scale where the phenomena of Rydberg blockade can affect the energy level of nearby atoms. Consider two atoms initially in the ground state with a small interatomic separation. If one of the atoms is excited to a Rydberg state with a resonant laser pulse, an energy shift is imparted on the second atom's Rydberg state. If the energy shift is far larger than the driving Rabi frequency, the second atom will be blocked from entering the Rydberg state. This is the key principle underpinning two- and multi-qubit gate sequences on neutral atoms systems.

Compared to other platforms, using Rydberg states of neutral atoms has several advantages. Firstly, the interaction can be switched on and off with relative ease through excitation and de-excitation pulses. Contrastingly, trapped ion qubits will always have coupling due to the nature of the strong Coulomb interaction. This allows for chains of up to 10 ions with all-to-all connectivity, however, as the system is scaled to more ions, the gate performance is reduced and requires re-optimisation of the protocol. Additionally, the limited range of the Rydberg interaction means gates can be performed with minimal cross-talk with qubits that are not part of the sequence. It is also possible to

locally address individual atoms by directing a driving laser on a site-by-site basis [14] or moving atoms from storage zones to interaction zones where gate light is applied to certain regions of the array [7].

To realise error correction on logical qubits, physical qubit fidelities need to reach $\mathcal{F} > 0.9999$ [114], but, current state-of-the-art Rydberg gates can currently attain $\mathcal{F} > 0.99$ for two-qubits [7, 14, 85, 100, 101]. Further enhancements of gate fidelities will be attainable with a combination of reduced gate times (which mitigate errors associated with the Rydberg state lifetimes), and through the engineering of target pair-states which reach stronger interaction regimes (minimising errors arising from imperfect blockade). The work presented in the next couple of chapters aims to identify Rydberg states which achieve the latter for multi-qubit gates between Rb and Cs atoms.

This chapter provides an overview of the Rydberg interaction and how it is exploited to realise entanglement between neutral atom qubits. A brief overview will be provided of gate sequences commonly used in Rydberg tweezer arrays with discussion towards implementation in a dual-species array.

3.1 Rydberg Blockade Gates

There are a few methods for performing multi-qubit gate operations using the Rydberg-atom interaction. The most common include blockade gates where only a single Rydberg excitation may exist, anti-blockade gates where multiple Rydberg excitations may take place during a sequence, and gates performed entirely within the Rydberg manifold between two or more Rydberg levels. The former two typically start and end with the qubits in the long-lived ground states and the Rydberg interaction is only used for entanglement. The latter can have one or both computational states in the Rydberg manifold but comes at the cost of increased errors associated with dephasing so is a less common approach. The experiments which this thesis is concerned with are qubits which interact under the blockade gate system as these are the most robust against Rydberg state errors.

The Rydberg blockade effect can be understood by considering a two-atom system

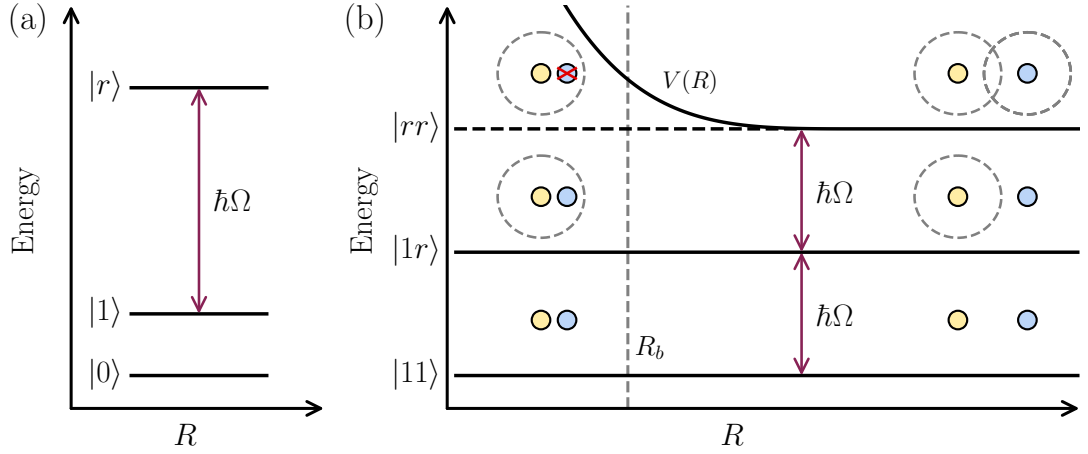


Figure 3.1: (a) Single atom energy levels with computational levels $|0\rangle$ and $|1\rangle$, and a single Rydberg level $|r\rangle$. Rydberg excitation is performed on the $|1\rangle \rightarrow |r\rangle$ transition. (b) The two-atom energy levels with double Rydberg state interaction energy $V(R)$ imparting an energy shift for small interatomic spacing R . The blockade region is defined by values of R less than the blockade radius R_b .

with interatomic separation R . Each of the atoms have the computational levels $|0\rangle$ and $|1\rangle$ encoded in their long-lived hyperfine ground state with the $|1\rangle$ state coupled to a single Rydberg level $|r\rangle$, as depicted in Fig. 3.1a.¹ This system can be represented by the following Hamiltonian

$$\begin{aligned} \hat{H} = & \frac{\hbar\Omega_A}{2} \left[|1\rangle \langle r|_A \otimes \hat{\mathbb{I}}_B + \text{h.c.} \right] \\ & + \frac{\hbar\Omega_B}{2} \left[\hat{\mathbb{I}}_A \otimes |1\rangle \langle r|_B + \text{h.c.} \right] \\ & + V(R) |r\rangle \langle r|_A \otimes |r\rangle \langle r|_B, \end{aligned} \quad (3.1)$$

where Ω_χ is the Rabi frequency of atom $\chi \in [A, B]$ and $V(R)$ is the interaction energy for both atoms being in the Rydberg state. The first and second terms correspond to the Rydberg excitation of atom A and atom B respectively. The full nature of the term $V(R)$ is provided in the next chapter, and here it is only presented as an interaction energy shift on the double Rydberg state.

¹In reality, many more levels may be used to realise a qubit experimentally, but to simplify analysis (both theoretically and computationally) other states are typically incorporated as parameters associated with $|0\rangle$, $|1\rangle$, and $|r\rangle$.

Chapter 3. Rydberg-Atom Interactions

By considering the case where atom B has been prepared in the $|r\rangle$ state, we can rewrite Eq. 3.1 for the evolution of atom A

$$\hat{H}_{\text{eff}} = \frac{\hbar\Omega_A}{2} \left[|1\rangle \langle r| + \text{h.c.} \right] + V(r) |r\rangle \langle r|. \quad (3.2)$$

Eq. 3.2 can be understood as a two-level light matter interaction where $V(R)$ can be considered an effective detuning on the $|r\rangle$ state of atom A. The time-dependent wave function of this system assuming atom A is initially prepared in the $|1\rangle$ state is,

$$|\langle r|\Psi(t)\rangle|^2 = \frac{\Omega_A^2}{\tilde{\Omega}^2} \sin^2\left(\frac{\tilde{\Omega}}{2}t\right), \quad (3.3)$$

where $\tilde{\Omega} = \sqrt{\Omega_A^2 + V(R)^2}$ is the generalised Rabi frequency. We can therefore define the Rydberg blockade condition as $V(R) \gg \Omega_A$ where atom A will be unable to be excited to the Rydberg state.

The interactions described by the Hamiltonians of Eq. 3.1 and Eq. 3.2 are shown schematically in Fig. 3.1b. In this picture, a single atom can be resonantly excited to a Rydberg state, however, the doubly excited Rydberg state can only be resonantly addressed at large values of R . The blockade radius $V(R_b) \approx \hbar\Omega$ defines the boundary at which blockade begins where for values of $R < R_b$, only a single Rydberg excitation may exist.

To understand how blockade can be used to realise a gate, consider a single target qubit initially prepared in the $|1\rangle$ state which undergoes a gate operation made up of a set of Rydberg pulses. The exact makeup of the Rydberg pulses is not important here and the only condition we are concerned with is that the target qubit is returned to the $|1\rangle$ state at the end of the gate sequence. The total evolution of the qubit is

$$|1\rangle_T \longrightarrow e^{-i\phi_T} |1\rangle_T, \quad (3.4)$$

where ϕ_T is the total phase accumulate by the target qubit. The value of ϕ_T is determined by a combination of the detuning, phase and amplitude of the Rydberg lasers. This allows any arbitrary phase to be realised for a single qubit.

Chapter 3. Rydberg-Atom Interactions

We now introduce a second control qubit which undergoes its own set of Rydberg pulses which may or may not be identical to the target qubit pulses. The phase of the two-atom wavefunction depends on the initial state of both the control qubit and target qubit. Just as we did for the single target qubit, we consider an unspecified set of gate pulses for the two-atom system. In this scenario, the computational basis states evolve as

$$|0\rangle_C \otimes |0\rangle_T \longrightarrow |0\rangle_C \otimes |0\rangle_T, \quad (3.5a)$$

$$|0\rangle_C \otimes |1\rangle_T \longrightarrow e^{-i\phi_{01}} |0\rangle_C \otimes |1\rangle_T, \quad (3.5b)$$

$$|1\rangle_C \otimes |0\rangle_T \longrightarrow e^{-i\phi_{10}} |1\rangle_C \otimes |0\rangle_T, \quad (3.5c)$$

$$|1\rangle_C \otimes |1\rangle_T \longrightarrow e^{-i\phi_{11}} |1\rangle_C \otimes |1\rangle_T, \quad (3.5d)$$

where $\phi_{\chi\chi}$ is the total phase accumulate during the gate sequence for $\chi \in \{0, 1\}$. The value of $\phi_{\chi\chi}$ is determined by both the initial states, and the pulses applied to the control and target qubits².

When engineering Rydberg gates, it is important to consider the total phase accumulated for the two-atom basis state. The natural encoding of a multi-qubit Rydberg gate is the $C^k Z$ gate where the total phase for each basis state (except the $|0\rangle^{\otimes N}$ state) is π . To illustrate, in a $k = 1$ system, the phase accumulated for each basis looks like

$$|0\rangle_C \otimes |0\rangle_T \longrightarrow |0\rangle \otimes |0\rangle, \quad (3.6a)$$

$$|0\rangle_C \otimes |1\rangle_T \longrightarrow e^{-i\pi} |0\rangle \otimes |1\rangle, \quad (3.6b)$$

$$|1\rangle_C \otimes |0\rangle_T \longrightarrow e^{-i\pi} |1\rangle \otimes |0\rangle, \quad (3.6c)$$

$$|1\rangle_C \otimes |1\rangle_T \longrightarrow e^{-i\pi} |1\rangle \otimes |1\rangle. \quad (3.6d)$$

Over the next few subsections, we will discuss some common gate sequences used on Rydberg systems to illustrate how the blockade effect can be utilised to set conditional phase rotations such as the CZ gate.

²Implicitly, $\phi_{00} = 0$.

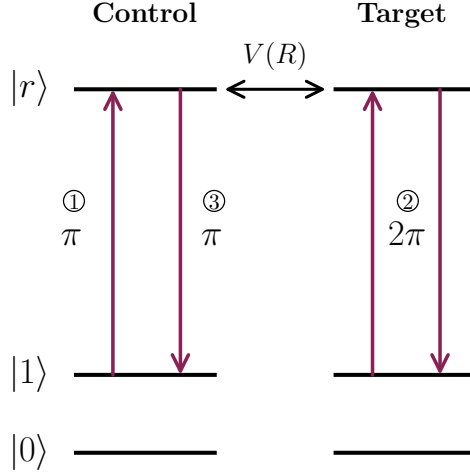


Figure 3.2: Three-pulse CZ gate scheme. (1) A π -pulse excites the control qubit from $|1\rangle$ to $|r\rangle$. (2) A 2π -pulse on the target qubit causes the accumulation of -1 phase if the control is initially in $|0\rangle$, whilst the target qubit will gain 0 phase if the control is initially in $|1\rangle$. (3) A final π -pulse returns the population of the control qubit from $|r\rangle$ to $|1\rangle$. This sequence implements a CZ gate.

The Three-Pulse Blockade Gate

The original CZ Rydberg gate presented in Ref. [83] considers two atoms³, a control and a target, which are located within each other's blockade radius and assumes each atom is individually addressable with resonant laser pulses. As shown in Fig. 3.2, three Rydberg pulses are applied to the atoms, (1) a π -pulse on the control qubit, (2) a 2π -pulse on the target qubit and (3) a π -pulse to the control qubit. For a single-photon excitation scheme each π -pulse imparts a $\pi/2$ phase on the qubit pair such that each orthogonal computational basis state evolves as [152]

$$\begin{aligned}
 |00\rangle &\xrightarrow{(1)} |00\rangle \xrightarrow{(2)} |00\rangle \xrightarrow{(3)} |00\rangle, \\
 |01\rangle &\xrightarrow{(1)} |01\rangle \xrightarrow{(2)} -|01\rangle \xrightarrow{(3)} -|01\rangle, \\
 |10\rangle &\xrightarrow{(1)} -i|r0\rangle \xrightarrow{(2)} -i|r0\rangle \xrightarrow{(3)} -|10\rangle, \\
 |11\rangle &\xrightarrow{(1)} -i|r1\rangle \xrightarrow{(2)} -i|r1\rangle \xrightarrow{(3)} -|11\rangle.
 \end{aligned} \tag{3.7}$$

³All though this gate sequence can be extended to k control qubits up to a single qubit rotation.

If Rydberg excitation is instead driven via a two-photon transition, an arbitrary two-qubit control phase gate can be realised by tuning the ratio between each photon's Rabi frequency. This arbitrary gate implementation can be represented in matrix form

$$C_\phi Z = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & e^{-2i\phi} & 0 & 0 \\ 0 & 0 & e^{-2i\phi} & 0 \\ 0 & 0 & 0 & e^{-2i\phi} \end{pmatrix}, \quad (3.8)$$

where ϕ is the phase accumulated during a resonant π -pulse. A CZ gate can be realised when $\phi = \pi/2$.

The three pulse gate is relatively simple to implement as it does not require chirping of detuning or laser amplitudes during the sequence. As such this gate has seen some use for implementing Rydberg gates [10, 153, 154].

However, this gate sequence is prone to errors, which means it is not suitable for reaching fidelities of > 0.9999 . One issue is that the control qubit spends nearly the entire gate sequence in the Rydberg state where spontaneous emission can lead to a large qubit dephasing. The time spent in the Rydberg state can be reduced by using higher Ω for the 2π -pulse on the target qubit, but will require strong Rydberg energy shifts to ensure the blockade condition is not broken. For scaling to $C^k Z$ gates, the associated errors with the target atom's Rydberg lifetime mean that errors become more acute for larger k [20]. For identical Rydberg interactions, cross-talk between control qubits is also increased for larger blockade. The latter point can be mitigated on a dual-species array by careful selection of Rydberg states, as discussed in the next chapter. For improved gate fidelities, other gate sequences should be used.

Global Gate Pulses

The three pulse blockade gate discussed in the previous section is highlighted as it can be naturally translated onto a dual-species system due to the requirement of independent qubit addressing. To reduce gate times, state-of-the-art schemes for single species qubits operate by illuminating qubits with the same laser such that all atoms experience the

same pulses which provides common-mode rejection of some errors. The two-qubit gate presented in Ref. [84] relies on the coupling between the $|1\rangle$ and $|r\rangle$ state with the $|00\rangle$ state remaining unperturbed during the sequence. Two Rydberg pulses with identical duration τ and detuning δ are applied sequentially to the atoms. τ is chosen such that a single pulse completes a full detuned Rabi-cycle for the $|11\rangle$ state while an initial $|01\rangle$ state will undergo an incomplete Rabi oscillation. Between the first and second pulse, a phase jump is implemented such that by the end of the sequence both the $|01\rangle$ and $|11\rangle$ states are returned to their original position on the Bloch sphere. At the end of the gate, the phase ϕ_{ij} is

$$\begin{aligned} |00\rangle &\longrightarrow |00\rangle, \\ |01\rangle &\longrightarrow e^{i\phi_{01}} |01\rangle, \\ |10\rangle &\longrightarrow e^{i\phi_{10}} |10\rangle, \\ |11\rangle &\longrightarrow e^{i\phi_{11}} |11\rangle, \end{aligned} \tag{3.9}$$

where by symmetry $\phi_{01} = \phi_{10}$. A CZ gate is realised when $\phi_{11} = 2\phi_{01} - \pi$ which is true when $\delta = 0.377\Omega$. The gate time for this sequence is $\sim 2.732\pi/\Omega$ which is far faster than the three pulse gate.⁴

Building on this protocol, further methods have been developed where a single common laser pulse is applied with a continuous phase profile which is chosen to achieve time-optimal operation of CZ gates [98]⁵. This principal has been extended to a parametrised version for which arbitrary control phase rotations can be achieved [85] and has been used in experimental state-of-the-art multi-qubit gates with fidelities of $\gg 0.99$ [7, 14, 85, 100, 101].

3.2 Dual-species Rydberg Gates

The gate sequences discussed in the previous section are highlighted as they have been implemented extensively experimentally [7, 10, 14, 85, 100, 101, 153, 154] so are ideal for benchmarking candidate Rydberg states for use as blockade gates. However, they

⁴The three pulsed gate has a gate time of $4\pi/\Omega$ assuming control and target qubits have the same Rabi frequency.

⁵This gate is $\sim 10\%$ faster than the gate presented in Ref. [84].

have some drawbacks for implementation at scale in a dual-species system.

The three pulsed π - 2π - π CZ gate can be implemented relatively easily on a dual-species platform due to the natural independent addressing from Rydberg levels and has been used to perform Rb-Cs gates previously [10]. However, it achieved rather low Bell state fidelities⁶ limited by dephasing during the relatively long time the control qubit spends in the Rydberg state.

The global gate pulses is highlighted as it has many of the highest gate fidelities achieved in Rydberg systems. The main limitation with implementation on a dual-species platform is the lack of common-mode rejection due to each atom requiring a separate laser for Rydberg addressing. That being said, the time optimal gate [98] has specifically seen recent implementation on a dual-species platform [130] perform with a non-Förster⁷ resonant pair state achieving a two-qubit fidelity of $\mathcal{F} = 0.975 \pm 0.002$. This result implies that these gates can provide relatively high gate fidelities on a dual-species platform. Further fidelity enhancement should be attainable by instead using stronger interacting Rydberg pair states, such as a Förster resonances [20, 132].

Although the common pulse gates provide a route to high gate fidelities, dual-species systems offer a unique opportunity to explore alternative methods of implementing gates. In this section, we shall discuss some novel proposals for using dual-species arrays to perform quantum computing.

CNOT_k via Rydberg Assisted Stark Shifts

The gate described in Ref. [126] realises a CNOT_k via a three step process. (1) A resonant Rydberg π -pulse on $|1\rangle \rightarrow |r\rangle$. (2) A Rydberg pulse with a detuning of Δ_r is applied to the target qubit's on the $|1\rangle \rightarrow |r\rangle$. Simultaneously, a microwave π -pulse detuned by δ and Rabi frequency Ω_T is applied to the target qubit's on the $|0\rangle \rightarrow |1\rangle$ transition. (3) Rydberg π -pulse identical to (1) is applied to the control qubit.

Consider this sequence applied in a $k = 1$ two qubit configuration. For an initial control qubit state of $|0\rangle_C$ there is no coupling to the Rydberg state, and the target

⁶ $\mathcal{F} \approx 0.49 \pm 0.02$ for the raw Bell state fidelity and $\mathcal{F} = 0.69 \pm 0.03$ for the SPAM corrected fidelity.

⁷Förster resonance is discussed in Sect. 4.1.

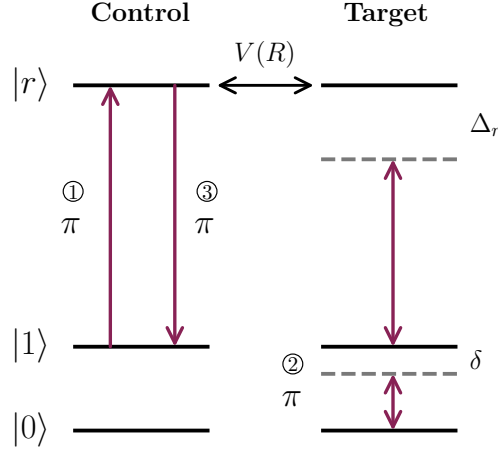


Figure 3.3: A CNOT gate where the Rydberg interaction shift imparted on the target qubit by the control qubit allows for the coherent transfer of the state from $|0\rangle$ to $|1\rangle$ with aid from a far off-resonant Rydberg laser

qubit remains in its initial state under the condition that $|\delta| \gg \Omega_T$. If instead the control qubit is initially in $|1\rangle_C$, there is an excitation to $|r\rangle_C$ and the target double Rydberg state will see an energy shift V . When the target qubit's Rydberg beam is turned on, there is no excitation pathway to the Rydberg state due to detuning Δ_r and instead the $|1\rangle_T$ state will experience an AC Stark shift of

$$\Delta_{AC} = \frac{\Omega_r^2}{4(\Delta_r + V)}, \quad (3.10)$$

where Ω_r is the Rydberg Rabi frequency of the target qubit. If the value of the microwave detuning δ is selected such that $\delta = -\Delta_{AC}$, a coherent population transfer of $|0\rangle_T \leftrightarrow |1\rangle_T$ will occur. This in effect means that the target qubit will flip if and only if the control qubit is Rydberg excited. To realise the CNOT gate in the computational basis, the second π -pulse is applied to the control qubit to return it to $|1\rangle_C$. In

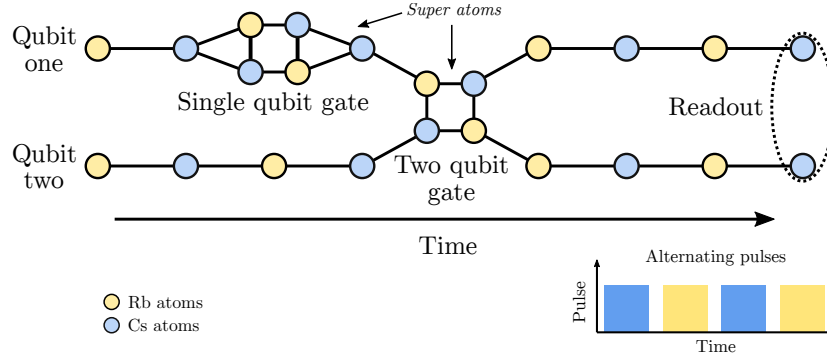


Figure 3.4: Schematic of the proposal for time-to-space implementation of an atom array quantum computer [155]. In this proposal, qubits are made up of alternating species called *wires*. Gates are engineered by the clusters of *super atoms*. Quantum information travels from left to right via a set of alternating Rydberg pulses as indicated by the inset in the lower right. The initial state is encoded in the first atom and the result of an algorithm is determined by the atoms at the end of the wire.

summary, each orthogonal computational basis will transform as

$$\begin{aligned}
 |00\rangle &\xrightarrow{\text{CNOT}} |00\rangle, \\
 |01\rangle &\xrightarrow{\text{CNOT}} |01\rangle, \\
 |10\rangle &\xrightarrow{\text{CNOT}} |11\rangle, \\
 |11\rangle &\xrightarrow{\text{CNOT}} |10\rangle.
 \end{aligned} \tag{3.11}$$

This gate sequence was originally proposed as an efficient method for fast qubit readout for quantum error correction on dual-species arrays by entangling many readout qubits to a single ancilla qubit. However, the gate sequence can also be utilised as a regular gate sequences for quantum computing. The total gate time is limited by the total time to perform a Rydberg 2π -pulse plus a microwave π -pulse. Faster gate times are attainable by instead using a two-photon Raman pulse for the target qubit rotation.

Time-to-space Quantum Gates

A novel idea for quantum computation on neutral atoms is suggested in Ref. [155] where the time component of a quantum circuit is transferred onto the atom geometry, as depicted in Fig. 3.4. This is realised by encoding single qubits as a 1D chains

built from alternating atomic species called wires. To perform quantum computation on a set of wires, the initial state of the system is encoded in the first atom of a wire and information is transferred down the wire via a series of sequential global Rydberg pulses. Multi-qubit interactions (wire-to-wire couplings) are realised with so called *super atoms* which consist of several Rb and Cs qubits between two wires which mediates the entanglement across the wires. Single qubit rotations can also be implemented by *super atoms* along individual wires.

This proposal is designed for qubits where the ground and Rydberg states act as the computational basis so is prone to errors associated with Rydberg scattering. Additionally, using wires to encode qubits scales poorly as physical qubits scale with the number of gates which is not generally the case in purely spacial logical qubits. However, the idea of using super atoms to perform entanglement could act as entanglement factories in a zone based architecture, similar to how magic state distillation⁸ can be implemented [156].

⁸Magic states allow for universal quantum computation by allowing non-Clifford rotations to be applied to logical qubits.

Chapter 4

Dual-species Rydberg Förster Resonances

It is common to address Rydberg levels in alkali atoms via a two-photon excitation scheme which avoids the use of UV wavelengths which are required for single-photon addressing of p -orbital Rydberg states. When using two-photon excitation in single species tweezer arrays, the natural choice of Rydberg levels lies in the s -orbital due to their near isotropic interactions [132]. To reduce laser overhead, it is typical to address the entire array with the same Rydberg lasers which means all atoms will be excited to the same Rydberg level. In the absence of external fields, low angular momentum Rydberg states don't have permanent dipole moments, which necessitates the use of high-principal quantum numbers $n > 70 - 80$. In a dual-species array, it is easier to address low-lying states which exhibit near resonant dipole-dipole (dd) interactions which permits richer control over the Rydberg-atom interaction.

Consider a surface code array geometry where Rb atoms are used as data qubits and Cs as ancilla qubits arranged in a checkerboard pattern as depicted in Fig. 4.1a. The surface code relies on strong nearest-neighbour interactions while next nearest neighbours on the diagonals should ideally experience no interactions. Given these requirements, it is necessary to identify Rydberg levels which exhibit strong interspecies interactions but weak intraspecies interactions. The typical nearest neighbours sepa-

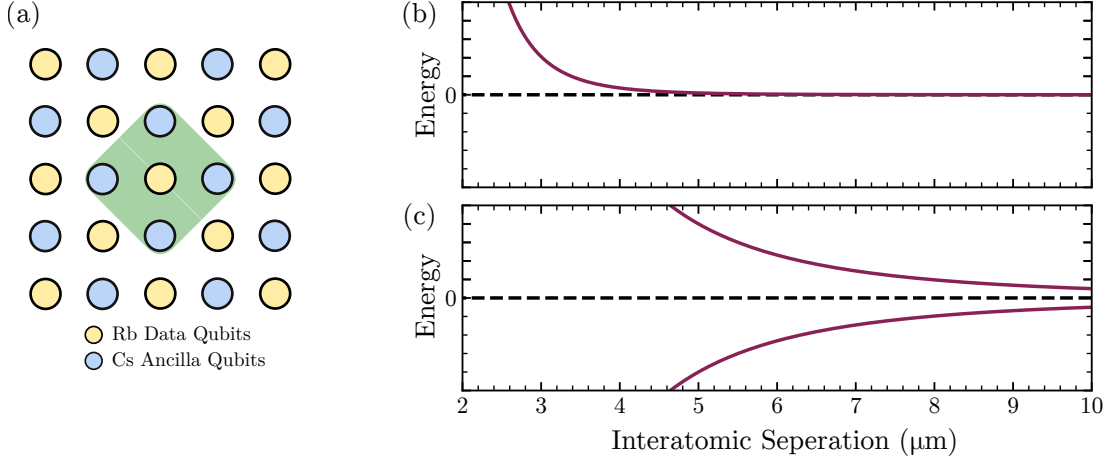


Figure 4.1: (a) Surface code of alternating of Rb and Cs atoms, desirable nearest neighbour interactions of the central Rb qubit is highlighted in green. (b) Illustration of vdW interaction found in from Eq. 4.5. $\Delta E \propto 1/R^6$ interactions makes these ideal for same-species couplings. (c) Illustration of dd type interaction found in Eq. 4.6. $\Delta E \propto 1/R^3$ interactions makes these ideal for Rb-Cs couplings.

rations in in tweezer arrays is $\sim 1 - 10 \mu\text{m}$ and at these ranges, there are two common types of Rydberg integrations. The first is the van der Waals (vdW) interaction regime which has Rydberg energy shifts of $\Delta E \propto 1/R^6$ for interatomic separation R as illustrated in Fig. 4.1b and is the most common interaction regime used in Rydberg arrays. The second interaction regime is the dd with energy shifts $\Delta E \propto 1/R^3$ which is illustrated in Fig. 4.1c and arise as a result of degeneracy in the pair states energy levels. The long tales of the dd interaction would make them suitable for enhancing interactions between Rb and Cs atoms, while the fast roll off of the vdW interaction would make it suitable for same species interactions.

This chapter will discuss dual species dd interactions between Rb and Cs atoms under a specific condition, known as Förster resonance, occurring when maximal degeneracy is present between Rydberg pair states. This builds on work previously carried out by in Ref. [132] by extending identification of Förster resonances to d -orbital Rydberg states and preform analysis via direct matrix diagonalisation to gleam insight beyond asymptotic approximations. Finally a gate sequence is modelled to demonstrate enhanced fidelities provided by Förster resonances in d -orbitals. The work presented

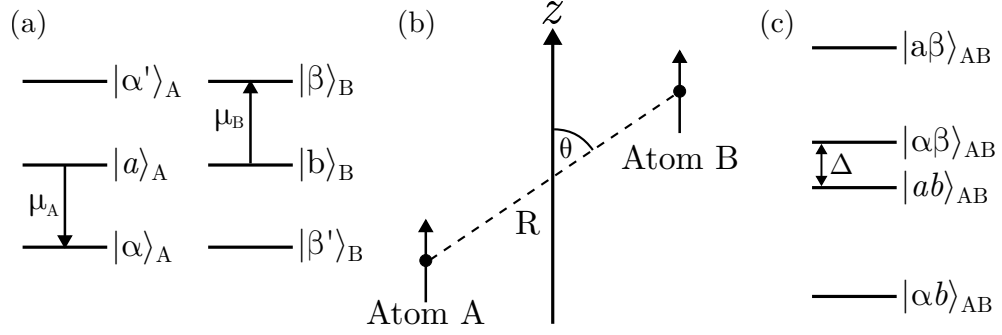


Figure 4.2: (a) Two atoms in the single atom basis, with target energy levels $|a\rangle_A$ and $|b\rangle_B$ with electric dipole μ_A and μ_B coupling to levels $|\alpha\rangle_A$ and $|\beta\rangle_B$. $|\alpha'\rangle_A$ and $|\beta'\rangle_B$ denote other Rydberg states which are not considered in this interaction picture. (b) Two dipolar atoms separated by R and angle θ with respect to the quantization axis z . (c) In the pair basis with target pair state $|ab\rangle_{AB}$ and nearest energy state $|\alpha\beta\rangle_{AB}$ separated by the smallest energy defect Δ .

in this chapter is expanded from work we published in Physical Review Research with the full manuscript available in Ref. [20].

4.1 Rydberg Atom Interactions

We begin by considering a pair of atoms, A and B, separated by distance R at angle θ with respect to the quantization axis, as illustrated in Fig. 4.2b. In the single atom basis, we define the target Rydberg states, $|a\rangle_A$ and $|b\rangle_B$ which have an electric dipole couplings μ_A and μ_B to the Rydberg states $|\alpha\rangle_A$ and $|\beta\rangle_B$ respectively, as shown in Fig. 4.2a. The dipole-dipole interaction between atom A and atom B is given by [157]

$$V(R) = \frac{\boldsymbol{\mu}_A \cdot \boldsymbol{\mu}_B - 3(\hat{\mathbf{n}} \cdot \boldsymbol{\mu}_A)(\hat{\mathbf{n}} \cdot \boldsymbol{\mu}_B)}{R^3}. \quad (4.1)$$

It is convenient to investigate the two atoms in the pair basis, as depicted in Fig. 4.2c, where the nearest pair states are separated by the smallest energy defect

$$\Delta = (E_\alpha - E_a) + (E_\beta - E_b), \quad (4.2)$$

where E_χ is the energy associated with level $\chi \in \{a, b, \alpha, \beta\}$.

We can consider the dipole-dipole interaction in the pair basis for two states $|ab\rangle$ and $|\alpha\beta\rangle$ by defining the Hamiltonian [158]

$$\hat{H} = \begin{pmatrix} 0 & V(R) \\ V(R) & \Delta \end{pmatrix}, \quad (4.3)$$

which has eigenenergies,

$$\Delta E = \frac{\Delta}{2} \pm \frac{\sqrt{\Delta^2 + 4V(R)^2}}{2}. \quad (4.4)$$

We can analyse Eq. 4.4 under two regimes. First we consider the case where $|\Delta| \gg V(R)$ where we perform an expansion of the second term of Eq. 4.4 for which first order energy shifts vanish and second order van der Waals (vdW) type perturbations dominate [159, 160]

$$\Delta E_{\text{vdW}} = -\frac{V(R)^2}{\Delta} = -\frac{C_6}{R^6}, \quad (4.5)$$

where the C_6 coefficient characterises the strength of the vdW interaction and scales as n^{*11} . In the case where $|\Delta| \ll V(R)$, where Eq. 4.4 becomes a resonant dipole-dipole (dd) interactions [160], which takes the form

$$\Delta E_{\text{dd}} = \pm V(R) = \pm \frac{C_3}{R^3}, \quad (4.6)$$

where the C_3 coefficient characterises the strength of the dd interaction and is proportional to n^{*4} . The case where $\Delta = 0$ is referred to as a Förster resonance which maximises the range at which dd interactions persist.

The magnitude of the C_3 coefficient between two Rydberg pair states can be approximated asymptotically for an interaction channel k using [132]

$$C_{3,k} = \frac{e^2}{4\pi\epsilon_0} \frac{\langle \alpha_k || r_A || a \rangle \langle \beta_k || r_B || b \rangle}{\sqrt{(2j_{\alpha_k} + 1)(2j_{\beta_k} + 1)}}, \quad (4.7)$$

where the target pair state $\langle \Phi_k || r_X || \Psi \rangle$ is the reduced matrix element for transitions between between states Ψ and Φ_k in the fine-structure basis for atom X, and j_ψ is the

total angular momentum quantum number associated to the state ψ .

For interspecies interactions, Förster resonance can be engineered by choosing relevant principal quantum numbers n and angular momentum quantum numbers ℓ for pair states which have small Δ and large dd coupling strengths $C_{3,k}$ [132]. Alternatively, Förster resonance can be induced with a dc electric fields [160–162] which shifts a near-resonant pair states to achieve $\Delta = 0$. Conversely, Förster resonances which exist at zero field are sensitive to external electric fields which necessitates the implementation of electric field nulling, as will be the case for the states discussed later in this chapter.

Because couplings between atoms of the same species is defined from a common target state, and Δ of a given pair is determined entirely by the inherent quantum defect for each species, there is typically only a single value of n which is near-resonant. This means interactions between the same species are typically restricted to in the weaker vdW regime. Fig. 4.3a shows the four nearest energy defects for ^{133}Cs s -orbital pair states where each atom addresses the same target state. It is clear that the value of n naturally yields a $\Delta = 0$. If we instead consider the addressing of dual-species Rydberg orbitals, we are no longer restricted to common target states which allows for the engineering of Förster resonances. This is shown in Fig. 4.3b where energy defects between the target state ($n_{\text{Rb}}d_{5/2} - n_{\text{Cs}}d_{5/2}$) and 4 near resonant states are shown. It is clear from the figure that by choosing different values n_{Cs} , we can obtain natural Förster resonances.

To get the energy shift of the atom’s double Rydberg state, we can sum over asymptotic dispersion coefficients for the relevant dipole-coupled pair states [132, 163],

$$\Delta E = - \sum_{n=0}^{\infty} \frac{C_n}{R^n}, \quad (4.8)$$

however for the typical atomic separations used in current tweezer array experiments ($3 \lesssim R \lesssim 15 \mu\text{m}$) accurate calculation of the resulting interatomic interactions requires exact diagonalisation of the pair state Hamiltonian. This allows for the accounting of higher-order dipole-quadrupole and quadrupole-quadrupole interactions, as well as other near resonant pair states rather than just the smallest energy defect. This is

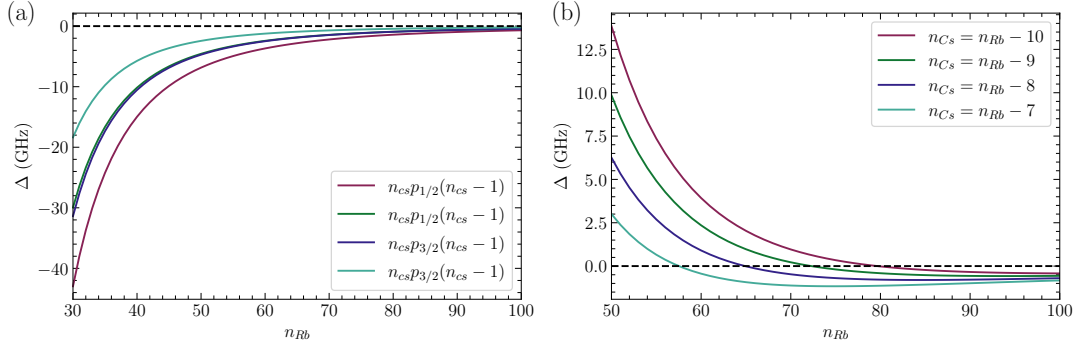


Figure 4.3: (a) Four nearest energy defects of the Cs-Cs target state $ns_{1/2} - ns_{1/2}$ commonly used for single level addressing in single species experiments. (b) Rb-Cs target state $n_{Rb}d_{5/2} - n_{Cs}d_{5/2}$ with Förster resonances at $\Delta = 0$ crossings.

discussed further in the next section.

4.2 Rubidium-Caesium d -state Förster Resonance

Suitable Förster resonant states for Rb and Cs have previously been found for the Rydberg s -orbitals [132], and were identified by asymptotic approximations of the pair potentials. This method can provide a good estimation of the Rydberg interaction strengths but will not present a complete picture of the pair potential as it does not account for higher order perturbations or other nearby pair states in the Rydberg manifold which may mix with the target pair state. Additionally, the work did not include other values of ℓ for the target state. We extend the work presented in Ref. [132] to identify suitable Förster resonant pairs of Rb and Cs in d -orbital angular momentum states.

To identify candidate Förster resonant pair states, we use Eq. 4.2 and Eq. 4.7 for calculating Δ and the channel coefficient $C_{3,k}$ respectively for all Rb-Cs dipole-coupled pair states from $d_{3/2}$ or $d_{5/2}$ target states. Due to the large number of possible Förster resonant pair states, we follow the approach of Ref. [132] to consider only resonances with $50 \leq n \leq 90$ that have $|C_{3,k}| > 1 \text{ GHz } \mu\text{m}^3$ and an energy defect $|\Delta| < 0.005$ of the level spacing. We identify resonant states in the Rb $d_{5/2}$ -Cs $d_{5/2}$, Rb $d_{3/2}$ -Cs $d_{3/2}$, Rb

$d_{3/2}$ -Cs $d_{5/2}$ and Rb $d_{5/2}$ -Cs $d_{3/2}$, each of which have 9 coupling channels k associated with them and are listed in Appendix A. With the aforementioned conditions, we identify 59 candidate Förster resonant pair states.

The approach used thus far has only considered the asymptotic approximation. To further reduce the total number of states to investigate, we move on to performed direct matrix diagonalisation of pair state Hamiltonian. This will provide a more accurate picture of the energy shifts allowing us to rule out states which do not present strong energy shifts. To perform matrix diagonalisation, we utilise an open-source library called PairInteraction which has a Python interface for studying Rydberg-atom interactions written with a C++ backend [164]. With PairInteraction, we can obtain a set of eigenvectors from which we can obtain the fractional overlap of the target state which admixes into other near resonant states. The library also provides the eigenenergy shifts of each Rydberg state.

Using PairInteraction, we now examine the characteristic pair potentials for the identified Förster resonances via direct matrix diagonalisation. Fig. 4.4 shows three example pair potentials identified as candidate sates using the asymptotic approach described above. Each one is calculated for $\theta = 0^\circ$ and considers atoms in magnetic sublevels $m_j^{\text{Rb}} = m_j^{\text{Cs}} = +3/2$. Fig. 4.4a shows the pair potential for the target state $76d_{3/2}$ -Cs $67d_{3/2}$ which has a resonance with the Rb $74f_{5/2}$ -Cs $65f_{5/2}$ state. Despite being Förster resonant, the two pair state branches do not diverge with the expected $\propto R^{-3}$ dependence which is desired for Rb-Cs interactions. As a result, the eigenstates present residual interactions at $\Delta E \approx 0$ and means that this pair state is a bad candidate for use in blockade gates. The same is true for the $74d_{3/2}$ -Cs $77d_{3/2}$ pair state shown in Fig. 4.4b where residual interaction exists near resonance. Of the three depicted states, only Fig. 4.4c demonstrates a clear Förster resonance with divergence of both pair state branches away from $\Delta E = 0$. The pair states in Fig. 4.4 shown that the asymptotic approach alone is not a robust method for identifying good Förster resonant pair states. Instead, we must devise an alternative method for selecting candidate states. A method for doing this is discussed in Sect. 4.2.2 for which we use blockade leakage as metric for characterising pair states.

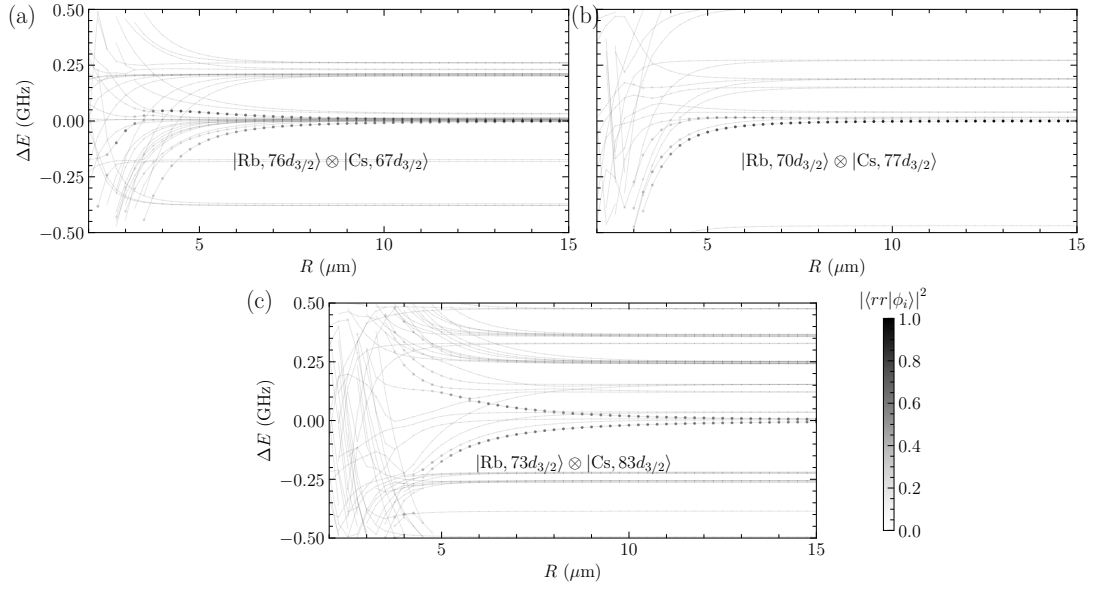


Figure 4.4: Example Förster resonant pair states for the Rb $d_{3/2}$ -Cs $d_{3/2}$ configuration calculated for $\theta = 0^\circ$ and $m_j^{\text{Rb}} = m_j^{\text{Cs}} = +3/2$. The faint lines indicate the eigenenergies as a function of R and the spots indicate the overlap with the target state. (a) Shows the pair state Rb $76d_{3/2}$ -Cs $67d_{3/2}$, which is resonant with Rb $74f_{5/2}$ -Cs $65f_{5/2}$; (b) shows Rb $70d_{3/2}$ -Cs $77d_{3/2}$, resonant with Rb $68f_{5/2}$ -Cs $79p_{1/2}$ and (c) Rb $73d_{3/2}$ -Cs $83d_{3/2}$, resonant with Rb $74p_{1/2}$ -Cs $81f_{5/2}$.

4.2.1 Angular Dependence

For Förster resonant interactions, the degeneracy of pair states is only lifted when the projection of total magnetic quantum number $M = m_j^A + m_j^B$ onto the internuclear axis is conserved $M = 0$ [163]. This means that careful consideration must be made for θ and m_j values when engineering Förster Rydberg interactions. Fig. 4.5 shows angular potentials for the Förster resonance between Rb $55d_{5/2}$ -Cs $63d_{5/2}$ pair state for different combinations of $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (5/2, \pm 5/2)$ and for angles $\theta = 0^\circ$ and 90° . Fig. 4.5a considers the $(5/2, 5/2)$ combination at 0° where the resonant pair states split as $M = 0$ resulting in the divergence of the resonant pair states. However, as shown in Fig. 4.5b, at $\theta = 90^\circ$ the $(5/2, 5/2)$ combination has a $\Delta M = \pm 2$ which causes the target state to mix with other eigenstates weakening the interaction. The strong Förster interaction can be recovered by flipping one of the atoms dipoles as depicted in Fig. 4.5d where the

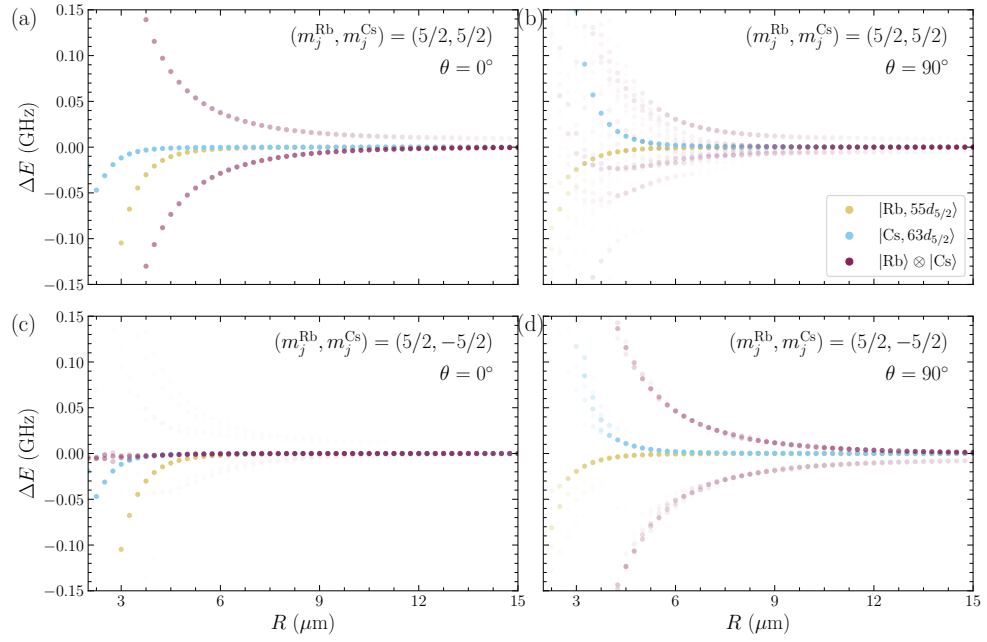


Figure 4.5: Angular dependence of the Rb $55d_{5/2}$ -Cs $63d_{5/2}$ for different θ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}})$ combinations. (a) $\theta = 0^\circ$, $(5/2, 5/2)$; (b) $\theta = 90^\circ$, $(5/2, 5/2)$; (c) $\theta = 0^\circ$, $(5/2, -5/2)$; (d) $\theta = 90^\circ$, $(5/2, -5/2)$.

$(5/2, -5/2)$ combination returns the interaction to the $\Delta M = 0$ condition. Finally, if the dipoles are anti-aligned with $(5/2, -5/2)$ and $\theta = 0^\circ$, the strong Förster interaction is eliminated due to the presence of a Förster zero interaction channel [165], as shown in Fig. 4.5c.

Fig. 4.5 also shows the Rb $55d_{5/2}$ - $55d_{5/2}$ and Cs $63d_{5/2}$ - $63d_{5/2}$ interaction curves both of which have flat interactions at $R \gtrsim 5 \mu\text{m}$ while displaying vdW C_6/R^6 like shifts at $R \lesssim 5 \mu\text{m}$. The Rb-Rb pair state displays weaker C_6 interactions than the Cs-Cs pair state despite being at a larger quantum number. This is due to the smaller Rb quantum defect yielding smaller interaction shifts in d -orbitals. This exemplifies the advantage of using dual-species to engineer Rydberg-atom interactions as at $\sim 6 \mu\text{m}$ there is a strong interspecies shifts while the intraspecies shifts are negligible.

Fig. 4.6 shows the pair potentials as a function of θ at $R = 6 \mu\text{m}$ which suggest that irrespective of the sign chosen for m_j^{Cs} , angles between $0^\circ < \theta < 90^\circ$ displays a complex energy landscape with residual interaction about $\Delta E = 0$. However, Fig. 4.6a

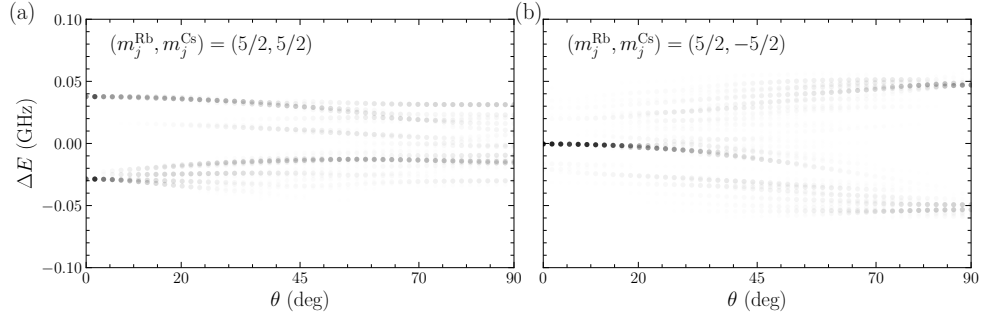


Figure 4.6: Angular dependence of the Rb $55d_{5/2}$ -Cs $63d_{5/2}$ for different values of θ with $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (5/2, \pm 5/2)$ calculated at $R = 6 \mu\text{m}$. (a) Reveal that strong Förster resonance is present at $\theta = 0^\circ$ for the $(5/2, 5/2)$ configuration; (b) reveals strong Förster resonance at $\theta = 90^\circ$ for

shows that a strong Förster resonance exists at $\theta = 0^\circ$ for the $(5/2, 5/2)$ configuration, while Fig. 4.6b show a strong Förster resonance at $\theta = 90^\circ$ for the $(5/2, -5/2)$ configuration. This would suggest that an array experiment must operate at either 0° or 90° depending on the chosen m_j^{Rb} and m_j^{Cs} configuration in order to exploit Förster resonant interactions.

Rydberg excitation schemes to reach s - and d -orbital Rydberg states are typically performed with circular polarised light propagating along the quantisation axis to ensure σ^+ - or σ^- -transitions are driven in a two-photon process. For experiments in 2D tweezer arrays, it is optimal to have isotropic Rydberg interaction in the plain of the array which is only possible at $\theta = 90^\circ$. This requires the quantisation axis to be defined out-of-plane to the array which proves an engineering challenge as Rydberg beams will have to propagate through the lens axis. It is instead easier experimentally to align the quantisation axis in-plane to the array. This would mean that the Förster resonant interactions described in this section will be of the form shown in Fig. 4.5d along one axis of the array, while the orthogonal axis will experience reduced interactions as shown in Fig. 4.5c.

4.2.2 Blockade Leakage

In Sect. 4.2, 59 combinations of Rb-Cs pair states were identified as being Förster resonant. As discussed at the start of this section, only a handful of these states will demonstrate good divergence away from resonance as the interatomic separation R is reduced. Here, we define the effect of blockade leakage as a metric to identify states which possess desirable Förster resonance characteristics which can be utilised for dual-species gate operations.

For high-fidelity gate operation, a Rabi frequency Ω is chosen to ensure the blockade condition is met ($V(R) \gg \Omega$) which ensures only a single excitation can be created. For Rydberg-atom interactions of the vdW type, the target eigenstate typically does not admix into other eigenstates, and as such the Rydberg manifold can be approximated as a single energy level. In the case of resonant dd interactions the pair potential includes coupling to many eigenstates $|\phi_i\rangle$ with eigenenergies ϵ_i all of which must be considered when considering Förster resonant Rydberg states. Eigenstates with an ϵ_i close to a driving laser's resonance can break the blockade condition allowing partial double excitation of the Rydberg state and will suppress the fidelity of a quantum gate. It is therefore not only necessary to consider the blockade effect of just the target Rydberg pair state $|r_t^A, r_t^B\rangle$, but the collective effect due to the weak couplings to other $|\phi_i\rangle$ as a result of blockade leakage.

To model the effects of blockade leakage, we define a Hamiltonian for this coupled two atom system where atom A is initially in the computational ground state and atom B is initially in the target Rydberg state is

$$\hat{H} = \hbar \sum_{i=1} \left(\left[\frac{\eta_i \Omega}{2} |\phi_i\rangle \langle 1^A, r_t^B| + \text{h.c.} \right] - \epsilon_i |\phi_i\rangle \langle \phi_i| \right), \quad (4.9)$$

where Ω is the Rabi frequency for coupling $|1^A\rangle \leftrightarrow |r_t^A\rangle$ and $\eta_i = \langle r_t^A, r_t^B | \phi_i \rangle$ denotes the relative overlap of the target pair state $|r^A, r^B\rangle$ with the eigenstate $|\phi_i\rangle$. The time evolution of the wavefunction is

$$|\Psi(t)\rangle = e^{-i\hat{H}t/\hbar} |\Psi_0\rangle, \quad (4.10)$$

where $|\Psi_0\rangle = |1^A, r_t^B\rangle$ is the initial state of the system. The probability of double Rydberg excitation is obtained by summing over all eigenstates with overlap with the target state

$$P_{rr}(t) = \sum_i |\langle \phi_i | \Psi(t) \rangle|^2, \quad (4.11)$$

while the probability of atom A remaining in the ground state is

$$P_{1r}(t) = |\langle 1^A, r_t^B | \Psi(t) \rangle|^2. \quad (4.12)$$

PairInteraction was used to obtain all ϵ_i , $|\phi_i\rangle$ and η_i for each of the 59 Förster resonant states identified in Sect. 4.2. These parameters were then used to build the Hamiltonian described in Eq. 4.9 for each Förster resonance. To reduce the number of states to consider without losing information about blockade leakage, we only consider eigenstates where $\eta_i > 10^{-3}$. The Rb-Cs interactions are calculated for $R = 6 \mu\text{m}$ at $\theta = 90^\circ$ and $(+m_j^{\text{Rb}}, -m_j^{\text{Cs}})$. To characterise the blockade, we consider a resonant π -pulse where $\Omega/2\pi = 2 \text{ MHz}$ and use Eq. 4.12 to obtain the probability of atom A remaining in the ground state $P_{1r}(t_\pi)$ which ideally is maximised. Results for all 59 states are presented in tables in Appendix B but Table 4.1 shows all pair states which present minimal blockade leakage — where $P_{1r}(t_\pi) > 0.999$. Eq. 4.7 asymptotic approximation for the $C_{3,k}$ coefficient and fitted¹ $\tilde{C}_{3,k}$ coefficient are also provided in the table. The $C_{3,k}$ and $\tilde{C}_{3,k}$ are not necessary identical due to the direct matrix diagonalisation taking into account the more complex nature of the interaction with higher order interactions and nearby eigenstates mixing. Additionally, the equivalent Rb-Rb and Cs-Cs interactions were investigated for next nearest neighbour interactions ($R = \sqrt{2} \cdot 6 \mu\text{m} \approx 8.5 \mu\text{m}$) to obtain $P_{rr}(t_\pi)$ which is the probability of a double Rydberg excitation after a π -pulse. The corresponding $\tilde{C}_{6,k}$ coefficients are also fitted and presented in the table.

Table 4.1 shows that there are many Rb-Cs Förster resonant pair states which offer a high suppression of blockade leakage and are suitable candidates for dual-species gate operation. While pair states with high principal quantum numbers ($n \gtrsim 65$) may

¹Fitting for $\tilde{C}_{3,k}$ and $\tilde{C}_{6,k}$ performed at $4 < R < 15 \mu\text{m}$.

Rb - Cs Förster Resonances		$\Delta/2\pi$	$C_{3,k}$	Rb-Cs at $R = 6 \mu\text{m}$		Rb-Rb at $R = 8.5 \mu\text{m}$		Cs-Cs at $R = 8.5 \mu\text{m}$	
Target Pair	Resonant Pair	(MHz)	(GHz μm^3)	$P_{1r}(t_\pi)$	$\tilde{C}_3$ (GHz μm^3)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)
*55 $d_{5/2}$ - 63 $d_{5/2}$	56 $p_{3/2}$ - 61 $f_{7/2}$	12.66	-2.24	0.9969	-11.30	0.9989	21.5	0.9840	-98.8
62 $d_{5/2}$ - 71 $d_{5/2}$	63 $p_{3/2}$ - 69 $f_{5/2}$	-9.53	-3.66	0.9997	15.79	0.9855	64.1	0.7747	-416.8
69 $d_{5/2}$ - 79 $d_{5/2}$	70 $p_{3/2}$ - 77 $f_{7/2}$	-19.14	-21.97	0.9998	26.34	0.7855	192.8	0.0013	-1448.9
70 $d_{5/2}$ - 80 $d_{5/2}$	71 $p_{3/2}$ - 78 $f_{7/2}$	13.69	-23.21	0.9990	-29.03	0.7442	223.9	0.0238	-1812.1
76 $d_{5/2}$ - 87 $d_{5/2}$	77 $p_{3/2}$ - 85 $f_{5/2}$	-19.30	-8.41	0.9997	40.58	0.1156	450.0	0.0388	-4482.5
*52 $d_{5/2}$ - 60 $d_{3/2}$	53 $p_{3/2}$ - 58 $f_{5/2}$	15.19	-6.70	0.9964	-8.37	0.9997	12.2	0.9958	-45.0
*59 $d_{5/2}$ - 68 $d_{3/2}$	60 $p_{3/2}$ - 66 $f_{5/2}$	13.40	-11.23	0.9992	-14.36	0.9951	41.6	0.9410	-206.1
66 $d_{5/2}$ - 76 $d_{3/2}$	67 $p_{3/2}$ - 74 $f_{5/2}$	11.12	-17.74	0.9997	-22.84	0.9176	114.4	0.4467	-692.9
79 $d_{5/2}$ - 91 $d_{3/2}$	80 $p_{3/2}$ - 89 $f_{5/2}$	-12.53	-36.95	0.9999	49.07	0.0070	619.2	0.0052	-3596.5
80 $d_{5/2}$ - 92 $d_{3/2}$	81 $p_{3/2}$ - 90 $f_{5/2}$	7.42	-38.76	0.9999	-51.12	0.0233	705.2	0.0144	-5756.9
*53 $d_{3/2}$ - 60 $d_{5/2}$	54 $p_{1/2}$ - 58 $f_{5/2}$	-4.41	-1.97	0.9968	7.12	0.9989	19.7	0.9938	-57.00
76 $d_{3/2}$ - 87 $d_{5/2}$	77 $p_{3/2}$ - 85 $f_{5/2}$	6.49	2.81	0.9986	-52.91	0.0129	723.9	0.0186	-5038.3
65 $d_{3/2}$ - 74 $d_{3/2}$	66 $p_{1/2}$ - 72 $f_{5/2}$	-13.07	-17.02	0.9974	16.67	0.7416	234.6	0.8297	-375.1
73 $d_{3/2}$ - 83 $d_{3/2}$	74 $p_{1/2}$ - 81 $f_{5/2}$	-0.70	-27.26	0.9995	28.54	0.1360	452.0	0.1082	-2156.1
80 $d_{3/2}$ - 91 $d_{3/2}$	81 $p_{1/2}$ - 89 $f_{5/2}$	-13.66	-39.62	0.9998	48.51	0.0440	1263.6	0.0234	-3178.2
81 $d_{3/2}$ - 92 $d_{3/2}$	82 $p_{1/2}$ - 90 $f_{5/2}$	4.39	-41.55	0.9997	-47.30	0.0590	1411.2	0.0153	-3836.9
72 $s_{1/2}$ - 70 $s_{1/2}$	71 $p_{3/2}$ - 70 $p_{1/2}$	-7.90	8.51	0.9992	9.82	0.0040	-1106.6	0.4005	-605.3

Table 4.1: The strongest Rb-Cs d -state Förster resonances identified as having $P_{1r}(t_\pi) \sim 0.999$ calculated at $R = 6 \mu\text{m}$, $\theta = 90^\circ$ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (+j^{\text{Rb}}, -j^{\text{Cs}})$ for a single-photon Rabi frequency of $\Omega/2\pi = 2 \text{ MHz}$. The interspecies blockade error is calculated for Rb-Rb and Cs-Cs interactions at $8.5 \mu\text{m}$. The asymptotic energy defect Δ and channel coefficient $C_{3,k}$ are provided along with the effective $\tilde{C}_3$ and $\tilde{C}_6$ coefficient obtained by fitting the pair potentials for $4 < R < 15 \mu\text{m}$. States labelled with an asterisk (*) are suitable for multi-qubit gates with strong interspecies but weak intraspecies couplings. Also shown is the only s -state resonance achieving $P_{1r}(t_\pi) > 0.999$.

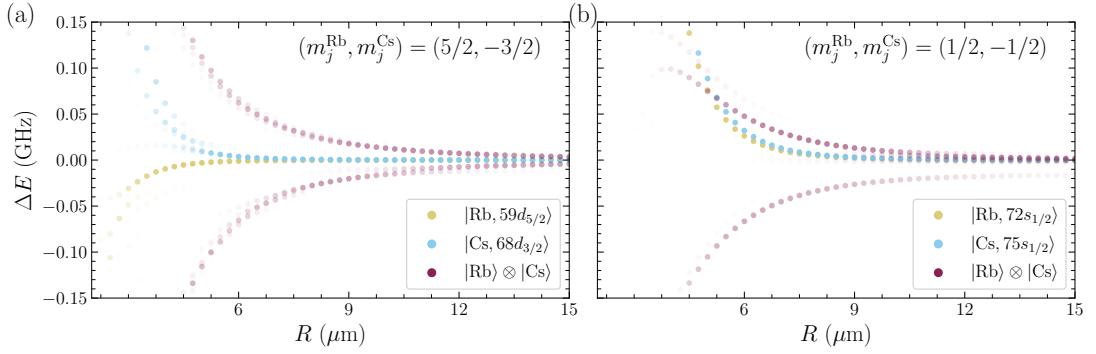


Figure 4.7: The pair potential for the best d -state (a) and best s -state (b) Förster resonance as identified as possessing the lowest blockade leakage between Rb-Cs atoms at $R = 6 \mu\text{m}$ (see Table 4.1).

demonstrate good blockade leakage suppression, there is also a large residual inter-species interactions which will increase the coupling between Rb-Rb and Cs-Cs atoms at $8.5 \mu\text{m}$ separation. This makes these states poor choices for gates with > 2 qubits. Increasing Ω will reduce the intraspecies couplings but this will also increase leakage error for the Rb-Cs interactions.

States in Table 4.1 denoted by an asterisk have the additional constraint that Rb-Rb interactions and Cs-Cs interactions have $P_{rr}(t_\pi) > 0.9$ which meets the overall aim of strong Rb-Cs couplings with reduced intraspecies couplings. The $59d_{5/2}$ - $68d_{3/2}$ is highlighted at the *best*² state due to possessing the highest $P_{1r}(t_\pi) = 0.9992$ for Rb-Cs interactions. The state also offers $P_{rr}(t_\pi) = 0.9951$ for Rb-Rb interactions and $P_{rr}(t_\pi) = 0.9410$ for Cs-Cs interactions. The relatively high blockade of the Cs-Cs interaction may pose a limit on the ascertainable gate fidelity. This could be mitigated by increasing the Ω associated with Cs excitation or by careful consideration of atom positioning ensuring only a single Cs atom is present when performing a gate sequence. The potential curves for this state are shown in Fig. 4.7a where the Rb-Cs interactions are denoted by the red curves and show strong divergence of the two near resonant

²Best under the conditions set for testing blockade leakage. Different nearest neighbour separations would require recalculation of blockade leakage. We chose $R = 6 \mu\text{m}$ for Rb-Cs separation as this offered good blockade leakage for at least one of the candidate states and is experimentally realisable in the planned experimental system (see Part II).

pair states with a $|\Delta E| > 50$ MHz at $R = 6 \mu\text{m}$. The yellow pair potential indicates interactions for Rb-Rb which shows negligible energy shift at $R = 8.5 \mu\text{m}$. The blue pair potential shows the interaction of Cs-Cs which has a small energy shift at $R = 8.5 \mu\text{m}$ which is the cause of the slightly blockaded interaction born out in the $P_{rr}(t_\pi)$ associated with it.

Comparison with s -orbital Förster Resonance

The analysis discussed above was applied to s -orbital Rydberg states to identify potential Förster resonances at lower angular momentum. Work similar to this has previously been carried out but only by accounting for the asymptotic approximation of Eq. 4.7 [132]. In total 9 Förster resonant pair states were identified and are shown in the Appendix B. The blockade leakage is also calculated for each candidate pair state where $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (1/2, -1/2)$, $R = 6 \mu\text{m}$ and $\theta = 90^\circ$. Under these conditions and only considering purely Rb-Cs interaction, s -orbital pair states have larger blockade leakage when compared candidate d -orbitals. For implementing gates, these results imply that s -orbitals will yield lower two-qubit gate fidelities than their d -orbital counter parts. When considering same species interactions there are no s -orbital states which simultaneously possess strong Rb-Cs couplings with well suppressed blockade leakage for Rb-Rb and Cs-Cs interactions. This will further reduce the attainable gate fidelities for ≥ 3 qubit gates when using s -orbitals.

For comparison Fig. 4.7b shows Rb $72s_{1/2}$ -Cs $75s_{1/2}$ which has a strong Förster resonant interaction for the Rb-Cs pair states, however, the strength of the divergence is clearly smaller than that of the d -state shown in Fig. 4.7a. This is likely due to the asymptotic $C_{3,k}$ coefficient of the s -state ($|C_{3,k}| = 9.65 \text{ GHz } \mu\text{m}^3$) is comparable to the d -states ($|C_{3,k}| = 11.23 \text{ GHz } \mu\text{m}^3$) even despite the lower smallest energy defect ($\Delta/(2\pi) = 2.7$ MHz) for s -states and $\Delta/(2\pi) = 13.4$ MHz. This highlights why blockade leakage is a better method for comparing Förster resonant states as it takes into consideration the larger number of states which effect Rydberg energy shifts, and not just the interactions through a single defect channel. When comparing the Rb-Rb and Cs-Cs interactions in Fig. 4.7, the s -orbitals have stronger couplings at $8.5 \mu\text{m}$ than

d -states which will mix unwanted blockade into the same species interactions.

4.2.3 Gate Performance

The blockade leakage error discussed in the previous section only describes the fractional coupling to the double Rydberg state due to imperfect blockade for Rb-Cs interactions, and the leakage of population out of the double Rydberg state due to residual intraspecies couplings for Rb-Rb and Cs-Cs interactions. Although these are important for characterising and comparing different pair states, this method does not take into account spontaneous emission from the Rydberg state or intermediate state scattering during the two-photon excitation to the Rydberg state. All of these processes must be taken into account to garner an accurate estimation of gate fidelity. Here we will aim to obtain an estimation of gate fidelity using interaction data obtained through Pair-Interaction and taking into account Rydberg decay and intermediate state scattering.

We chose to use the three pulse gate described in Sect. 3.1 to performing multi-qubit gate modelling for the Rb $59d_{5/2}$ -Cs $68d_{3/2}$ Förster resonant target state. As previously noted, the π - 2π - π pulse sequence is not an optimal pulse sequence due to the length of time the control qubit is left in the Rydberg state, however, it provides a simple approach to extend to C^kZ gates yielding the correct output up to a single qubit rotation. This pulse sequence has also been demonstrated experimentally for Rb-Cs multi-qubit gates using an s -orbital Förster resonant states [10], and can act as a benchmark for modelling d -orbital Förster resonance gate fidelities.

Modelling was performed with the method presented Ref. [22] which reduces computational complexity by only considering each atom with three computational levels, two ground states $|0\rangle$ and $|1\rangle$, and a single Rydberg state $|r\rangle$. To incorporate the effects of spontaneous decay from the intermediate state, P_e^1 and P_e^r weightings are introduced as dissipation terms to define the residual excited state population due to coupling between the $|1\rangle$ and $|r\rangle$ transitions respectively. The parameters associated with two-photon excitation are determined by performing adiabatic elimination of excited state

hyperfine levels f_e requiring calculation of the following terms

$$\begin{aligned}
 \Omega_R &= \sum_{f_e} \frac{\Omega_1^{f_e} \Omega_{f_e}^r}{2\Delta_{f_e}}, \\
 \delta_{AC} &= \sum_{f_e} \frac{|\Omega_1^{f_e}|^2 - |\Omega_{f_e}^r|^2}{4\Delta_{f_e}}, \\
 P_e^1 &= \sum_{f_e} \frac{|\Omega_1^{f_e}|^2}{4\Delta_{f_e}^2}, \\
 P_e^r &= \sum_{f_e} \frac{|\Omega_{f_e}^r|^2}{4\Delta_{f_e}^2},
 \end{aligned} \tag{4.13}$$

where Ω_R is the effective two-photon Rabi frequency with single-photon Rabi frequencies $\Omega_1^{f_e}$ and $\Omega_{f_e}^r$, Δ_{f_e} is the single-photon detuning with respect to the intermediate state hyperfine levels, and δ_{AC} is the differential AC Stark shift. The effective single atom Hamiltonian describing atom-light interactions is given by

$$\begin{aligned}
 \hat{H}_{\text{eff}} &= \frac{\hbar}{2} [-iP_e^1 \Gamma_e |1\rangle \langle 1| \\
 &\quad + \Omega_R (|r\rangle \langle 1| + |1\rangle \langle r|) \\
 &\quad - (\delta + i\Gamma_r + P_e^r \Gamma_e) |r\rangle \langle r|],
 \end{aligned} \tag{4.14}$$

where Ω_R is the effective two-photon Rabi frequency, δ is the two-photon detuning and $\Gamma_{e,r}$ are the excited and Rydberg state linewidths. With multiple intermediate hyperfine states for two-photon excitation, the phase after a 2π -pulse $\phi_{2\pi}$ can take any value on the interval $[0, 2\pi)$ and depends upon the relative laser power and the intermediate state detuning [153]. To ensure the wave function acquires $\phi_{2\pi} = \pi$, for each species we model the single atom excitation without adiabatic elimination and adjust the power ratio of the two lasers to tune to achieve a phase shift error of $|\phi_{2\pi} - \pi| < 10^{-5}$.

The parameters $\Omega_1^{f_e}$ and $\Omega_{f_e}^r$ are calculated assuming a centre of mass intermediate state detuning $\Delta_{f_e} = 10$ GHz and a targeted two-photon Rabi frequency of $\Omega_R/(2\pi) = 5$ MHz. Schematic of the excitation scheme for Rb and Cs are shown in Fig. 4.8a and Fig. 4.8b respectively. Rydberg excitation from $|1\rangle_{\text{Rb}} = |5s_{1/2}, F=2, m_f=0\rangle$ to $|r\rangle_{\text{Rb}} = |59d_{5/2}, m_j=5/2\rangle_{\text{Rb}}$ via the $|e\rangle_{\text{Rb}} = |6p_{3/2}\rangle$ intermediate state is modelled

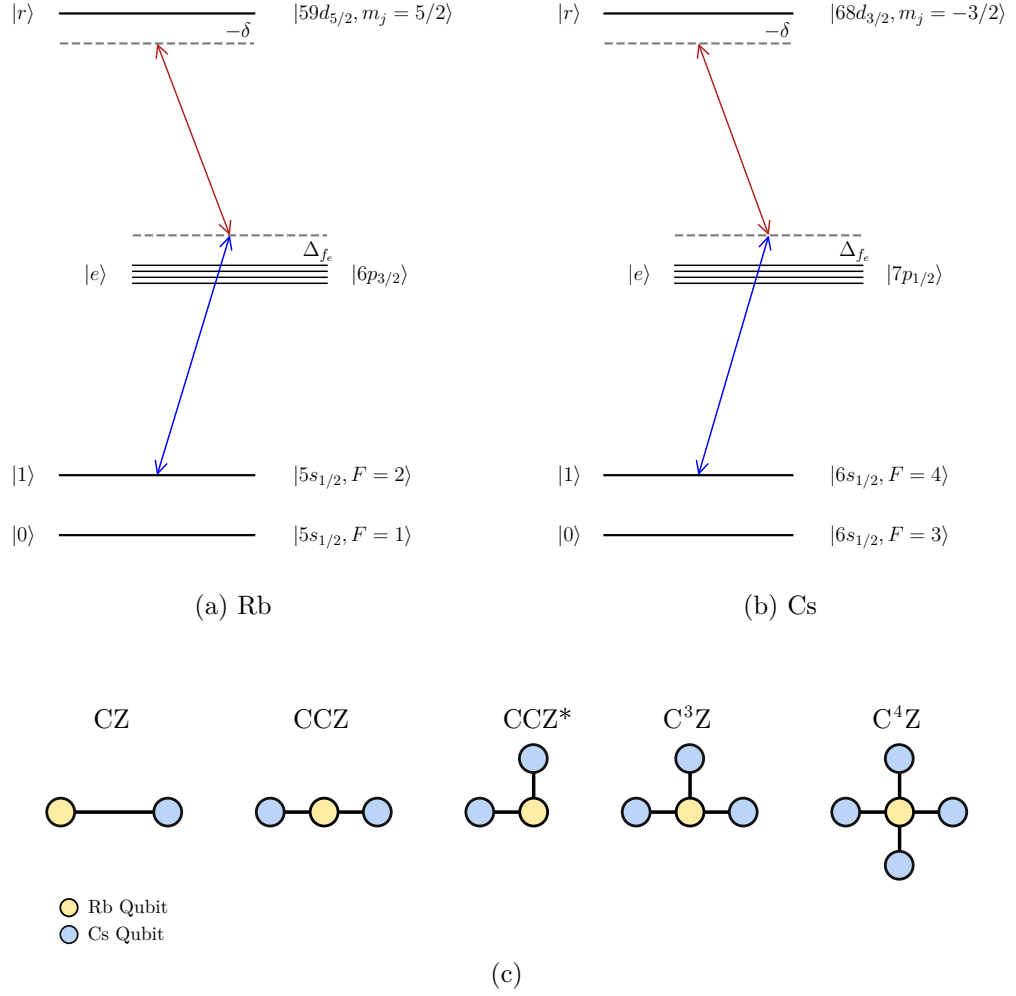


Figure 4.8: (a) Rydberg excitation scheme for Rb; (b) Rydberg excitation scheme for Cs. (c) Shows the three configurations dual species gates where Förster interaction are indicated by the black line between Rb (yellow) and Cs (blue) qubits.

with $\Gamma_e/(2\pi) = 1.4$ MHz and $\Gamma_r/(2\pi) = 0.67$ kHz with a pair of σ^+ -polarised photons. For 10 μm waist laser beams, choosing a power of 5.2 mW for the $|1\rangle_{\text{Rb}} \rightarrow |e\rangle_{\text{Rb}}$ transition and 240 mW for the $|e\rangle_{\text{Rb}} \rightarrow |r\rangle_{\text{Rb}}$ transition results in $\phi_{2\pi}^{\text{Rb}} = \pi$ shift with $P_e^1 = 5.0 \cdot 10^{-4}$ and $P_e^r = 1.7 \cdot 10^{-4}$. Performing similar calculations for Cs with the excitation pathway $|1\rangle_{\text{Cs}} = |6s_{1/2}, F=4, m_f=0\rangle$ to $|r\rangle_{\text{Cs}} = |68d_{3/2}, m_j = -3/2\rangle_{\text{Cs}}$ via the $|e\rangle_{\text{Cs}} = |7p_{1/2}\rangle$ intermediate state is modelled with $\Gamma_e/(2\pi) = 1.0$ MHz and $\Gamma_r/(2\pi) = 0.68$ kHz with a pair of σ^- -polarised photons. A power of 20 mW and 150

mW for the $|1\rangle_{\text{Cs}} \rightarrow |e\rangle_{\text{Cs}}$ and $|e\rangle_{\text{Cs}} \rightarrow |r\rangle_{\text{Cs}}$ respectively. This led to a $P_e^1 = 5 \cdot 10^{-4}$ and $P_e^r = 1.2 \cdot 10^{-4}$ for Cs for a $\phi_{2\pi}^{\text{Cs}} = \pi$.

Using the above excitation parameters and the Hamiltonian $\hat{H}_{\text{eff}}$ from Eq. 4.14, we first model a two-qubit CZ gate with the π - 2π - π pulse sequence described in Sect. 3.1. This qubit layout is shown in the first schematic of Fig. 4.8c. The generalised initial state for N qubits is

$$|\Psi\rangle_N = \frac{1}{N} \sum_{i,j,\dots,N=0}^1 |i,j,\dots\rangle, \quad (4.15)$$

which under ideal conditions will evolve after a C^kZ gate as

$$|\Psi\rangle_N \xrightarrow{C^kZ} |\Psi\rangle_I, \quad (4.16)$$

where $|\Psi\rangle_I$ is the qubits state after an ideal gate. For a two-qubit CZ gate

$$|\Psi\rangle_I = \frac{1}{2} (|00\rangle - |01\rangle - |10\rangle - |11\rangle). \quad (4.17)$$

We first consider a system with Rb as the control qubit and Cs as the target qubit for an interatomic separation of $R = 6 \mu\text{m}$ for which the atoms experience an interaction shift of 65.6 MHz. This results in a gate fidelity of $\mathcal{F}_{CZ} = 0.9953$. When instead the Cs is set as the control, and Rb is set as the target qubit $\mathcal{F}_{CZ} = 0.9955$. These two-qubit gate fidelities are competitive with state-of-the-art demonstrations of two-qubit gates [7, 14, 85, 100, 101]. It should be noted that the cited examples used gate sequences such as the Levine-Pichler [84] and the time-optimal [98] gates described in Sect. 3.1 which are less prone to errors when compared to the three pulse gate. Higher $\mathcal{F}_{CZ}$ would be attainable if the chosen d -orbital Rydberg states were tested against a more optimised gate sequence, however, this was beyond the scope of the work presented here. Realising the fidelities found here for the CZ gates would prove a challenge experimentally due to factors not considered in the model, such as laser stability or errors associated with the angular dependence of the interaction (see Sect. 4.2.1). Hence the fidelities stated here are an upper limit to what would be expected in an experiment.

Extending the gate sequence to a CCZ gate where the target Cs qubit is along a

Gate configuration	Control qubits	$\mathcal{F}_{C^kZ}$
CZ	Rb	0.9953
CZ	Cs	0.9955
CCZ	Rb	0.994
CCZ^*	Rb	0.983
C^3Z	Rb	0.988
C^4Z	Rb	0.913

Table 4.2: Summary of C^kZ gate fidelities. (*) indicates CCZ gate is performed in an L-shaped geometry as opposed to a line geometry.

line with two Rb control qubits $6\ \mu\text{m}$ either side of it (as shown in Fig. 4.8c) has a gate fidelity of $\mathcal{F}_{CCZ} = 0.994$. If one of the targets is moved to the orthogonal axis such that the three qubits form a triangle where the Rb-Rb interatomic distance is $8.5\ \mu\text{m}$, the gate $\mathcal{F}_{CCZ} = 0.983$. Extending to a C^3Z gate where the three control Rb atoms form a triangular configuration, as shown in Fig. 4.8c has $\mathcal{F}_{C^3Z} = 0.988$. Finally, a C^4Z gate with a single target Cs qubit in the centre yields a $\mathcal{F}_{C^4Z} = 0.913$. Table 4.2 provides an overview of the fidelities for each configuration.

4.3 Implementation of Förster Resonant d -states

This chapter has discussed the origins of Förster resonant Rydberg interactions and how they can be specifically exploited in a dual-species system to extend the range and strength of interactions. An extended overview of the work from Ref. [20] was provided in this chapter where identification of candidate d -orbital Förster resonant states between Rb and Cs Rydberg atoms was performed. These states display strong Rb-Cs interactions while intraspecies couplings are suppressed as required for multi-qubit C^kZ on chequerboard type geometries. The identified states have improved interaction characteristics when compared to s -orbital Förster resonances which have previously been identified and used on Rb-Cs experimental platforms [10, 132].

Despite the apparent advantages, d -orbital Rydberg states will have some challenges to reach the expected enhanced gate fidelities presented in this chapter. The anisotropic interaction discussed in Sect. 4.2.1 means the quantisation axis and array geometries must be set to ensure interactions occur at $\theta = 90^\circ$. For optimal gate interaction in

the array plane, the quantisation axis must be normal to the array — i.e. through the lenses. This requires both the bias field and Rydberg laser beams to be aligned along the lens axis to ensure purely σ^+ - or σ^- -transitions are used when performing Rydberg excitation, which motivates a future move to an objective with a hole. If instead the quantisation axis is defined along one of the axes of the array, a strong Förster resonant interaction can still be achieved in a line limiting the multi-atom gates to two or three qubits³. This configuration is still workable on a scalable platform on a zone based architectures as described in Ref. [7]. The reduction in blockade as angles away from 90° for $(m_j^{\text{Rb}} > 0, m_j^{\text{Cs}} < 0)$ configuration means that errors in atom positioning must be minimised. This can be achieved if both species are in a common array with well defined trap to trap location, or using by two separate arrays with separate spatial light modulators (SLMs) which provides fine control of array positioning by tweaking parameters of each hologram independently.

The work presented in this chapter identified Förster resonant states which provide strong interspecies couplings but weak intraspecies couplings. If instead strong Rb-Rb or Cs-Cs interactions are desired, a different set of states should be identified. However, due to the increased complexity of choosing different energy levels for separate same species atoms, the likely method for achieving Förster resonance would be to tune $\Delta = 0$ via external electric fields [160]. We did not investigate Rb-Rb or Cs-Cs Förster resonances as we do not intend work in a regime where either of these would be used.

³See the *CZ* and *CCZ* configuration of Fig. 4.8c.

Part II

Experimental Design

Chapter 5

Laser Setup

This chapter will discuss the laser systems required for achieving single atoms in traps. We will begin with a brief overview of laser cooling for magneto-optical traps (MOTs) and polarisation gradient cooling (PGC). This chapter will then move on to discuss the laser system that this experiment will use to realise a dual-species MOT. The setup of the MOT system requires us to account for optics which can operate at both rubidium's and caesium's respective wavelengths. This necessitates a system which combines both species' lasers to ensure that the MOTs overlap within the experiment. The method for doing this is presented here. We will end this chapter by briefly discussing the optical dipole trap (ODT) laser system for each atomic species.

5.1 Laser Cooling

The first step in achieving single atoms is to use laser cooling to remove kinetic energy from a room temperature cloud and bring the effective atom temperature low enough to enable trapping in ~ 1 mK deep ODTs. The MOT has been the workhorse of cold atom physics since first being produced in 1987 [166], and performs the bulk of cooling during the molasses stage in cold neutral atom experiments. We utilise two different MOTs for this platform. The first is a 2D MOT and loads Rb or Cs from background vapour in the high pressure side of our vacuum systems, and is discussed in Sect. 6.1. The second MOT is situated at the cryogenic stage within the low-pressure science

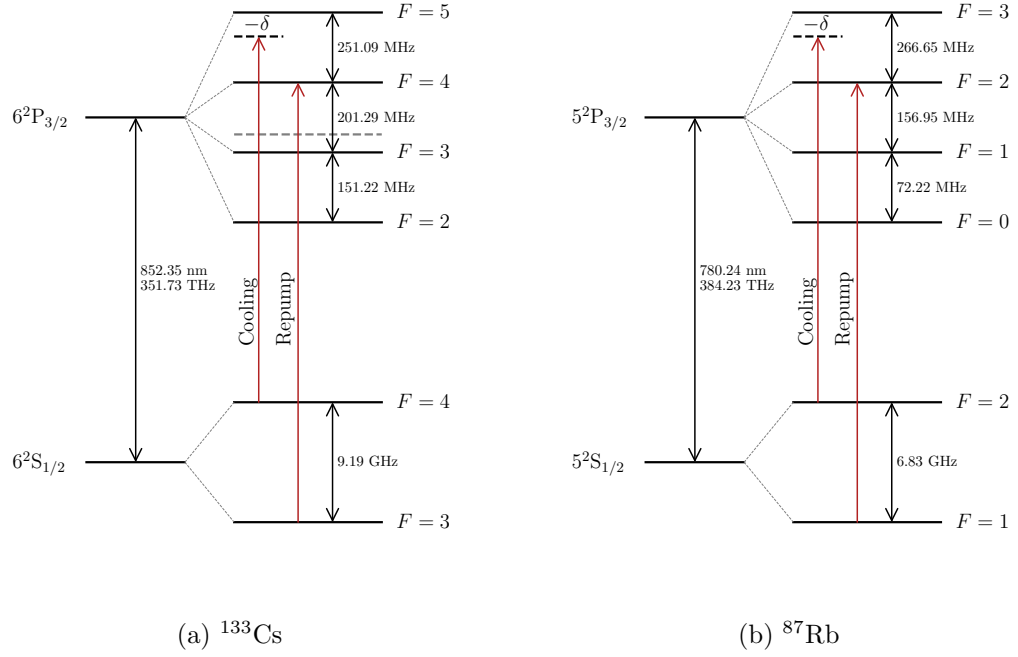


Figure 5.1: D₂ transition hyperfine structure with frequency separations marked for caesium (a) and rubidium (b). Cooling and repump are denoted by the red arrow with cooling detuning $-\delta$. The $F = 3 \rightarrow F' = 2 - 4$ crossover is marked by the grey dashed line for caesium. [78, 79]

chamber where it receives atoms from the 2D MOT and cools them for loading ODTs, and is discussed in Sect. 6.2.

There are two main components to a MOT, an optical molasses and a quadrupole field. The optical molasses forms when counter-propagating laser beams red-detuned from resonance create a velocity dependent force to slow the atoms down [167]. The addition of a magnetic quadrupole field — either by an anti-Helmholtz coil pair or permanent magnets — creates a position dependent light interactions arising from spatially varying Zeeman shifts in the atom's magnetic sublevels. This results in localisation of the slowed atoms about the magnetic field zero [168]. To insure repeated photon scatterings by the atoms, laser cooling requires a closed atomic transition.

Alkali atom MOTs are predominantly produced on the D₂ line corresponding to the transition from the S_{1/2} ground state to the first P_{3/2} excited state, where a closed transition exists. Fig. 5.1a shows the D₂ hyperfine structure for ¹³³Cs where the cooling

is performed on the $F = 4 \rightarrow F' = 5$ transition with laser frequency detuned by $-\delta$. Although this is a closed transition, off-resonant scattering on the $F = 4 \rightarrow F' = 4$ is still possible where atoms can decay into the $F = 3$ state due to spontaneous emission. This removes atoms from the cooling cycle and causes rapid depletion of the MOT, which requires an additional laser to ensure dark atoms are returned to the cooling cycle. It is typical to perform this repumping resonantly on the $F = 3 \rightarrow F' = 4$ transition so atoms can decay back into the $F = 4$ state where they are returned to the cooling cycle. The equivalent MOT cycle for ^{87}Rb is shown in Fig. 5.1b where cooling is performed red-detuned from the $F = 2 \rightarrow F' = 3$ transition and repumping performed resonantly on the $F = 1 \rightarrow F' = 2$ transition.

The lower bound on temperature in a MOT is defined by the Doppler limit [78, 79],

$$T_D = \frac{\hbar\Gamma}{2k_B}, \quad (5.1)$$

where $\hbar = h/(2\pi)$ is the reduced Planck's constant, Γ is the natural linewidth of the cooling transition and k_B is the Boltzmann constant. The natural linewidth for the ^{133}Cs cooling transition is $2\pi \cdot 5.234$ MHz which gives a limit of $T_D = 126$ μK [79]. The natural linewidth for the cooling transition in ^{87}Rb is $2\pi \cdot 6.065$ MHz which gives a limit of $T_D = 146$ μK [78]. These values for T_D are of order 10% the trap depth of the tweezers, and may limit the trap lifetime due to the finite heating rate of the traps [169]. Cooling below the Doppler limit can be achieved by employing a stage of polarisation gradient cooling (PGC).

PGC is performed in molasses with no quadrupole field. When two counter-propagating circular-polarised beams of light overlap, they produce a standing wave of linear polarisation which rotates in a helical pattern. Consider an ensemble of atoms at rest, for which the polarisation will look linear in the atom's frame. Atoms in this regime will experience π -polarised light to the excited state. The Clebsch-Gordan coefficients suggest the most likely decay paths will cause atoms to populate the $m_f = 0$ magnetic-sublevel with the greatest probability, and no net force will be experienced by the atoms. Consider instead an ensemble of atoms moving along the axis of beam propagation. In

this regime the atoms will see a fictitious magnetic field leading to a Raman induced population imbalance in the magnetic sublevels with a net force against the direction of propagation, hence cooling the atom [170].

The temperature limit of PGC cooling has the following relation,

$$T_{\text{pgc}} \propto \frac{\hbar\Omega}{k_B|\delta|}, \quad (5.2)$$

where Ω is the Rabi frequency. Therefore, colder temperatures can be reached during PGC either by reducing laser intensity I or increasing δ [170].

Cooling once atoms are loaded into the ODTs can be performed with stages of single atom PGC [121, 123, 171], and additionally with Raman sideband cooling [172] or resolved sideband cooling [173] which takes atoms to their motional ground state. The vibrational state does not limit the ability to create clock state qubits and as such this system will not implement sideband cooling.

5.2 Cs MOT Lasers

5.2.1 Cooling Laser

Fig. 5.2a shows a schematic of the laser cooling setup which consists of a pair of independent cooling and repump lasers, required due to the large ground-state hyperfine splitting (9.19 GHz in Cs). The Cs cooling laser is a homebuilt external-cavity diode laser (ECDL) with a 150 mW diode at 852 nm (Thorlabs: L852P150) [174]. To enable precise frequency control, a double-pass AOM is used to frequency shift the laser before locking to the $F = 4 \rightarrow F' = 5$ transition using polarisation spectroscopy (pol-spec)[175, 176]. The Cs laser is then boosted from 20 mW to 2.4 W using a tapered amplifier (TA) (m2k: TA-0850-2000-DHP) to ensure sufficient power for both the 2D MOT and the 3D MOT optics. AOMs after the TA provide independent control of the 2D MOT and 3D MOT detunings. The 2D MOT and 3D MOT are coupled into separate 1 m patch fibres taking the light to the combination section described in Sect. 5.4.

Chapter 5. Laser Setup

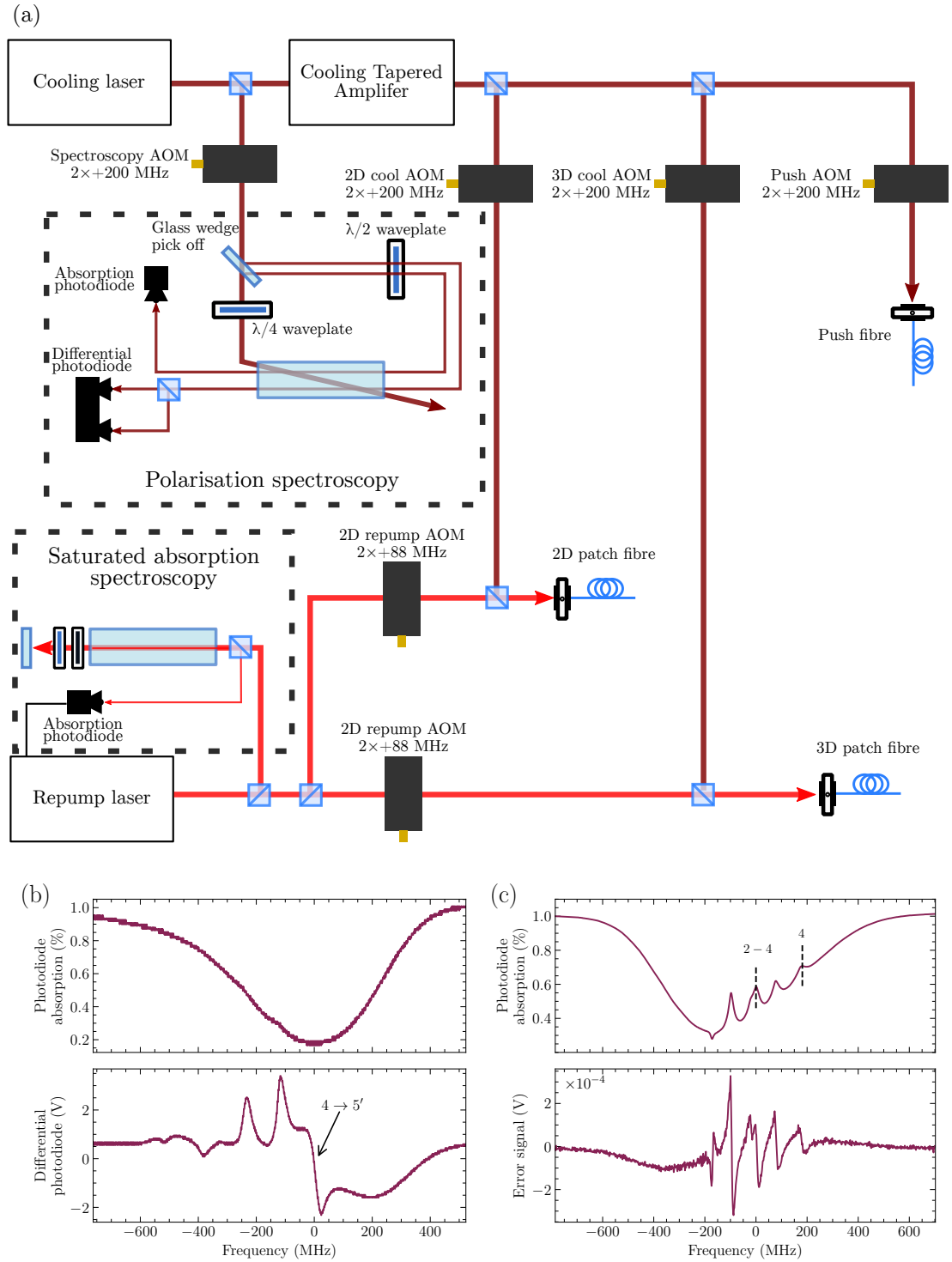


Figure 5.2: (a) Schematic of the cooling and repump laser systems. All acousto-optic modulators (AOMs) are set in a double pass configuration, however, are shown as single pass for simplicity. (b) Cs cooling locking absorption photodiode signal (upper plot) and differential photodiode signal (lower plot). (c) Cs repump locking absorption photodiode signal (upper plot) and error signal (lower plot).

The upper box in Fig. 5.2a shows the pol-spec layout. A glass wedge picks off two weak probe beams from the transmitted pump beam. The pump beam is circularly polarised by a $\lambda/4$ waveplate and passed through the vapour cell in one direction. The Zeeman degeneracy is lifted with a small magnetised Allen key situated next to the vapour cell. A $\lambda/2$ waveplate sets the polarisation of the two probe beams to be linear at 45° , which then pass through the glass cell overlapping with the pump beam. A photodiode measures the absorption signal of one of the probe beams which is shown in the upper plot of Fig. 5.2b for the D_2 line from the $F = 4$ ground state. A polarising beamsplitter (PBS) decomposes the second probe beam which is then measured with a differential photodiode to produce an error signal, shown in the lower plot Fig. 5.2b. The use of pol-spec for locking the cooling laser does not provide absolute frequency reference when setting the MOT detuning; alternatively modulation transfer spectroscopy [177] would provide dispersive line shapes with zero background, however, we did not implement this here as pol-spec provided a stable lock.

5.2.2 Repump Laser

For Cs repumping we use a TOPTICA DL Pro at 852 nm which can output up to 150 mW. The laser is locked directly to the $F = 3 \rightarrow F' = 2 - 4$ cross-over transition using saturated absorption spectroscopy (sat-spec) [178]. The repump for the 2D MOT and 3D MOT are frequency shifted with two AOMs to bring the frequency of the laser on resonance with the $F = 3 \rightarrow F' = 4$ repump transition.

The sat-spec optics layout is shown in the lower box of Fig. 5.2a and consists of a strong pump beam which is passed through a vapour cell. After the first pass through the cell, a neutral density filter reduces the laser intensity to act as a probe beam. A $\lambda/4$ waveplate and retro-reflecting mirror ensure the probe beam is counter-propagating to pump beam with orthogonal polarisation. A PBS directs the probe beam to a photodiode with the resultant signal used to generate the error signal. The raw photodiode signal and error signal from the Cs sat-spec is shown in Fig. 5.2c for the repump transition.

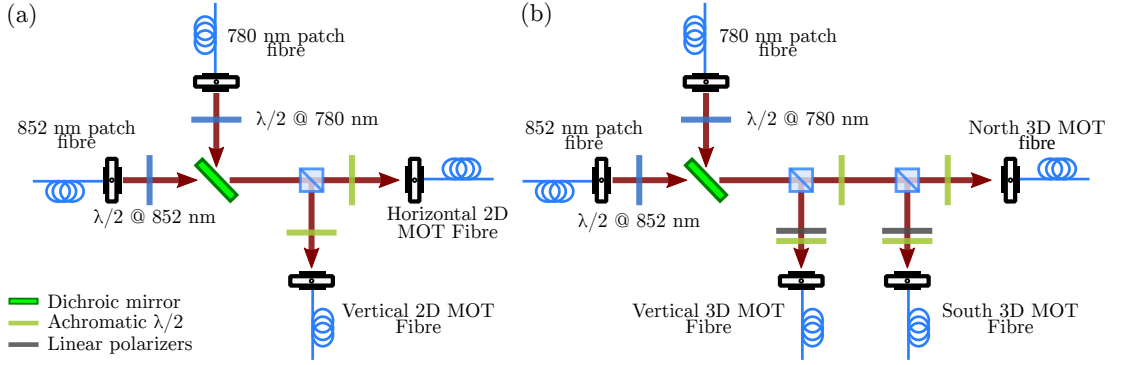


Figure 5.3: Laser beam combination for 852 nm and 780 nm the 2D MOT (a) and 3D MOT (b).

5.3 Rb MOT Lasers

The Rb MOT laser set-up is a mirror image of the Cs setup described in Sect. 5.2. Both Rb cooling laser and repump laser are homebuilt ECDLs with a 70 mW diode (Eagleyard: EYP-0790-03010-1500-SOT02-0000) used for cooling and a 120 mW (Sharp: GH0781JA2C) diode used for repump. The TA for amplifying the cooling light is a COHERENT I15N-TA-0785-3000-DHP which outputs a power of 1.1 W during normal operation.

5.4 MOT Beam Combination

To simplify the alignment of the MOT beams on the science table, it is desirable to combine the 852 nm laser beams and 780 nm laser beams into a single optical fibre. This could be achieved either with beam splitting fibres which offer a compact design, or with free space optics offering improved flexibility. We chose to go with the latter.

Fig. 5.3a shows the 2D MOT beam combination, 852 nm and 780 nm exit separate patch fibres and are passed through a $\lambda/2$ waveplate for setting the power splitting between the horizontal 2D MOT beam fibre and vertical 2D MOT beam fibre. A dichroic (Thorlabs: DMLP805) overlaps the 852 nm and 780 nm beams and a PBS splits powers for coupling into the vertical and horizontal 2D MOT fibres. We employ achromatic waveplates which are $\lambda/2$ at both 852 nm and 780 nm (EKSMA Optics:

467-4210) for polarisation matching into the 10 m long fibres. These fibres take the light to the 2D MOT on the science table.

Fig. 5.3b shows the 3D MOT beam combination, again, 852 nm and 780 nm come out of separate patch fibres and are passed through a $\lambda/2$ waveplate to set power splitting between the vertical 3D MOT beam and horizontal 3D MOT beams. A dichroic overlaps the 852 nm and 780 nm beams and a PBS then picks off some power for coupling into the vertical 3D MOT fibre. The remaining light passes through an achromatic $\lambda/2$ waveplate for setting the power balance of the horizontal 3D MOT beams, a second PBS splits the light for coupling into each of the horizontal MOT fibres. Achromatic waveplates are used immediately before the long fibres for polarisation matching. In this configuration we found that the south horizontal MOT beam was subject to polarisation noise. This was due to the reflection port of the PBS having a polarisation extinction ratio (PER) of 20:1 which caused ellipticity of the light coupled into the fibre and hence poor polarisation matching¹. To rectify this, we added a linear polariser (Thorlabs: LPNIRE100-B) set to maximise transmission, ensuring good linear polarisation. The achromatic waveplate was then set to match the light's polarisation into the fibre. This eliminated the polarisation noise.

The push beams are not combined in a specified combination section and are instead each coupled into the same fibre immediately after the AOM shown in Fig. 5.2a.

5.5 Trapping Lasers

For this experiment, we chose to trap each atomic species in separate arrays of ODTs with the wavelengths selected to ensure species-selective trapping. Analysis presented in Ref. [18] suggests that the light shifts of Rb produce an attractive potential in the wavelength range of 800 nm – 850 nm while the light shifts of Cs in this wavelength range has a repulsive potential². Therefore, a trapping wavelength in this range is ideal for ensuring species-selective trapping for the Rb arrays. When considering Cs-selective

¹It is common for PBS manufactures to guarantee good PER on their transmission port but make no promises about their reflection port. Alternatively a Wollaston prism can provide good PER for both s- and p-polarisation. Due to budget constraints, we chose to take the polariser approach.

²Due to the $\sim 5\text{--}10\text{ }\mu\text{m}$ trap spacing, the effects of light shifts from neighbouring traps is negligible.

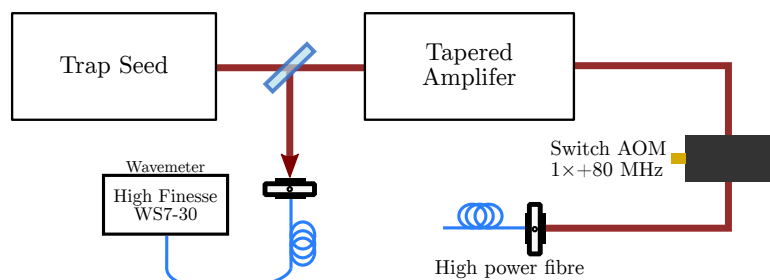


Figure 5.4: Schematic of the trap laser setup.

trapping, wavelengths far red-detuned from the D_1 and D_2 transitions have attractive ground state potentials for both Rb and Cs. For the 900 nm – 950 nm wavelength range, however, the Rb trap depths will be approximately four times smaller than that of Cs. This means Rb is effectively untrapped for suitably small trap depths. As discussed in Sect. 2.3 and from analysis presented in Ref. [149], at wavelengths between ~ 935 nm and ~ 938 nm, the Cs cooling transition is magic³, which will allow for in-trap cooling and imaging without the need for out of phase chopping between trap and imaging light.

The wavelength ranges for each species have been demonstrated in previous dual-species arrays [10, 127, 128]. The magic wavelength for Cs has specifically been utilised in the generation of Rb-Cs molecules in tweezer arrays as these systems require minimisation of light shift due to the large dipole moment [18].

5.5.1 Cs Trap Laser

Fig. 5.4 provides a schematic of the trap laser path. The seed laser is a homebuilt ECDL with a diode (Eagleyard: EYP-RWE-0980-08020-1500-SOT02-0000) pulled to 940 nm. During operation, the diode outputs 25 mW. An anti-reflective (AR) wedge picks off a small amount of light which is coupled into a wavemeter for wavelength monitoring. The remaining light is used to seed a 4 W TA (Coherent: I15N-TA-0950-3000-DHP). The spatial mode out of the TA is non-Gaussian and requires beam shaping optics immediately after the output to ensure good coupling into an optical

³Magic here means the differential light shift between the ground and excited states is zero, as described in Sect. 2.3.

fibre. The TA has broad amplified spontaneous emission (ASE) extending to 921 nm where the excited state of the cooling transition has a dipole allowed coupling to the $6^2D_{3/2}$ state. Excitation to this state will cause atom loss during imaging. A 940 nm bandpass filter (Thorlabs: FBH940-10) which has a bandwidth of 10 nm was installed to mitigate this effect. The light is single passed through an 80 MHz AOM for switching. Finally, the light is coupled into a high-power single mode optical fibre (OZ Optics Ltd: HPUC-2A3AHPC-1040-P-8AS-11) to ensure that the light incident on the spatial light modulator (SLM) has a good Gaussian mode.

The trap laser and components described here are all placed on a single breadboard allowing for the setup in one location before moving the system to the main science table. During experimental operation, the breadboard sits on the science table.

5.5.2 Rb Trap Laser

The layout for the Rb trap laser is nearly identical to the layout of the Cs trap laser. The Rb trap seed laser is a homebuilt ECDL with a diode (Eagleyard EYP-RWE-0840-06010-1500-SOT02-0000) pulled to 810 nm. During operation, the diode outputs 25 mW. An AR wedge picks off a small amount of light which is coupled into a wavemeter for wavelength monitoring. The remaining light is used to seed a 3 W TA (Coherent I15N-TA-0810-3000-DHP).

Chapter 6

Vacuum System

An essential requirement for laser cooling and trapping is the operation within an ultra-high vacuum (UHV) environment to suppress atom losses due to collisions with background contaminants. This project deviates from the typical approach of room temperature vacuum systems that achieve trap lifetimes of order 10 s. Motivated by recent developments with trap lifetimes extending to thousands of seconds [9, 123] a new experimental platform has been developed integrating a closed-cycle cryostat into the vacuum system.

A render of the vacuum system is shown in Fig. 6.1. The science chamber is the heart of the experiment and where atoms are cooled in a 3D MOT before they are loaded into the tweezer arrays. The science chamber is maintained at UHV by a 75 L/s ion pump with additional pumping provided by periodic firing of a titanium sublimation pump (TSP). The closed cycle cryostat provides a third source of vacuum pumping at temperatures of < 30 K when the effects of cryopumping occurs. The bottom of the cryostat extends into the main chamber and houses the high-numerical aperture (NA) lens used for focusing tweezer arrays and fluorescence imaging, and superconducting coils to generate the quadrupole fields for the 3D MOT.

The 2D MOT chamber is the high pressure side of the vacuum system and sources the Rb and Cs atoms from two ampoules and transports atoms towards the science chamber. The chamber is kept under UHV by a 5 L/s ion pump. The 2D MOT chamber and science chamber are separated by a gate valve that can be closed to

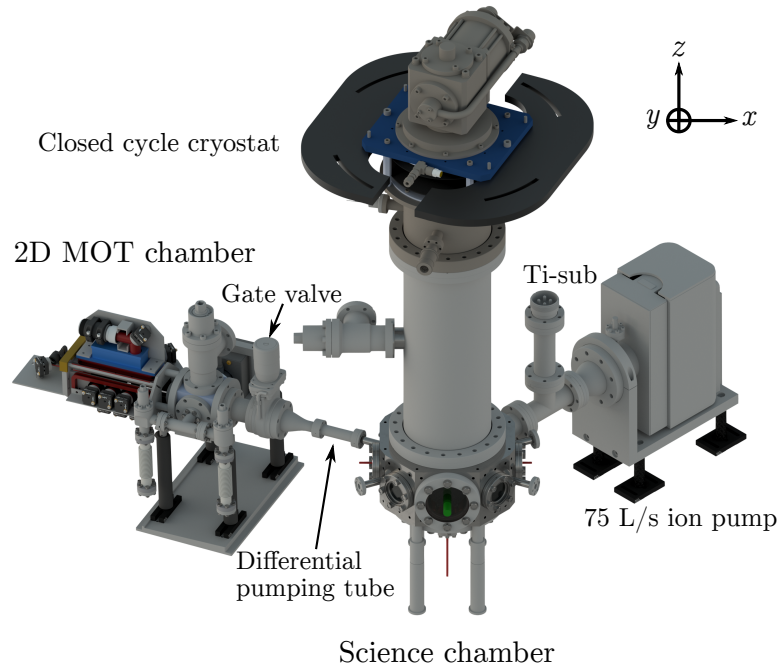


Figure 6.1: Computer aided design (CAD) rendering of the vacuum system highlighting the main sections and key components.

isolate each of the vacuum systems. Atoms are transferred from the 2D MOT to the 3D MOT via a differential pumping tube between the two chambers.

6.1 2D MOT

The 2D MOT vacuum system is shown in Fig. 6.2a and is an improved design from the permanent magnet MOT used by the square project [26]. The vacuum chamber consists of a DN40CF 6-way cube which holds a 5 L/s ion pump, a DN16CF tee and a glass cell where the 2D MOTs are created. A 112 mm long metal rod with a 3 mm hole is connected to a drilled out blank to act as a differential pumping tube. It is attached to the cube face opposite the glass cell. The ampoules reside in flexible metal bellows attached to the tee via an all metal gate valve; Rb and Cs ampoules are attached to opposite ports – as shown in Fig. 6.2b. Ampoules sit inside a hollow metal cylinder to ensure they crack when the bellows are bent, shown by the cut away view in Fig. 6.2c. The partial pressure of Cs or Rb in the 2D chamber can be adjusted by opening and

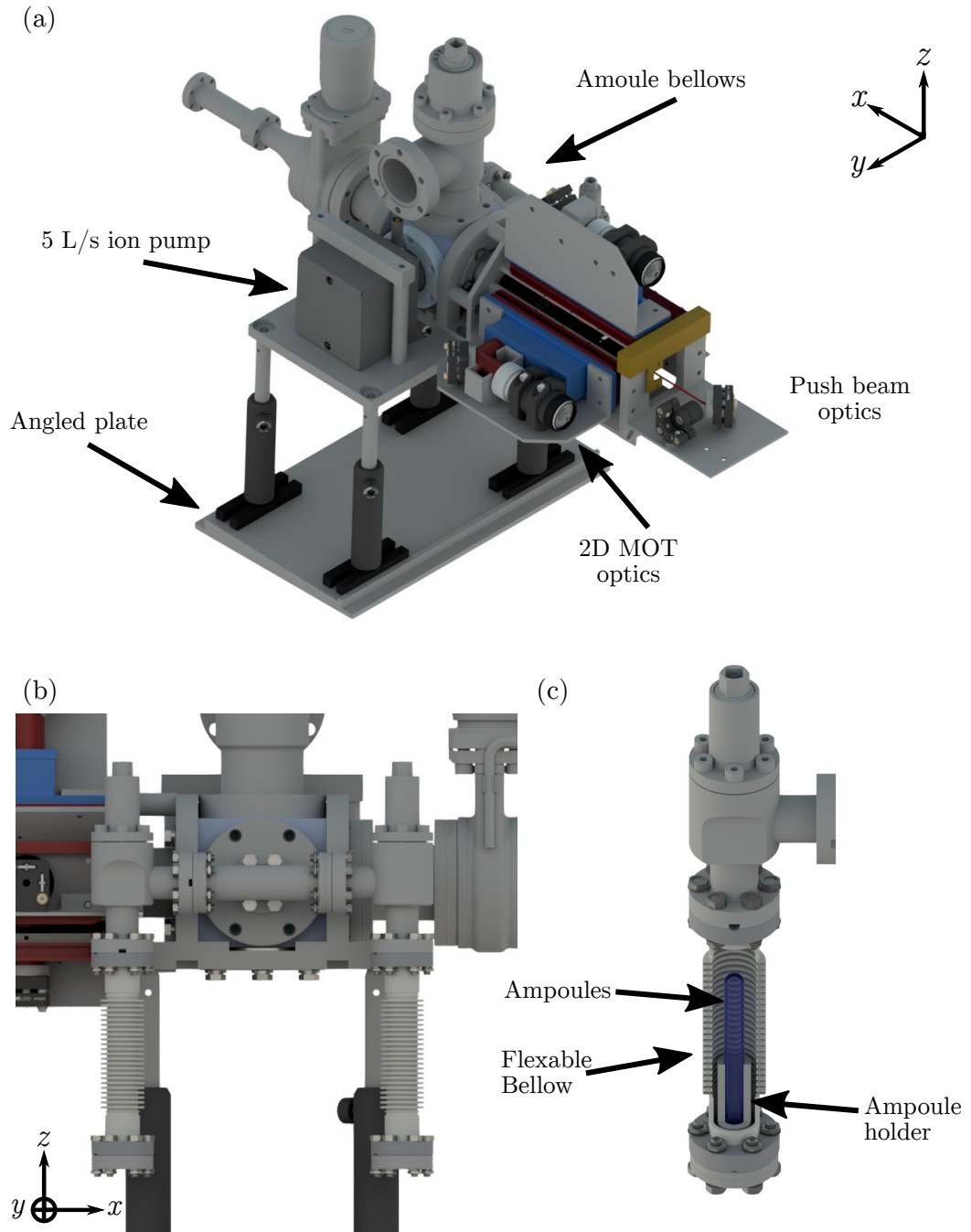


Figure 6.2: (a) Render of the 2D MOT vacuum system. (b) Render of the reducing tee which holds the bellows containing the Cs and Rb ampoules. (c) Sliced view of the ampoule cracker, the ampoule resides in a metal holder inside the bellow to ensure it cracks when the bellow is bent.

closing the respective valves on the tee. When we cracked the Cs ampoules, it took over a month to see fluorescence in the glass cell. This is because Cs will stick to surfaces not previously coated so is slow to diffuse through a vacuum system particularly when there is poor conductance, as is the case for the 2D chamber. Dispensers would have been an alternative atom source to avoid the slow diffusion, but require periodic breaking of vacuum to replace [179].

An aluminium plate holds the cube, ion pump and support structure for the 2D MOT optics which surround the glass cell. The 2D MOT consists of three regions with quarter waveplates held by the cage outside the glass cell. Around the waveplates resides 4 rectangular shim coil formers which produce 3 G/A at the atoms. Two sets of 5 permanent magnets sit in rectangular plates on opposite sides of the glass cell at a 45° angle relative to the horizontal and vertical MOT beams, as shown in Fig. 6.3b. The orientation of the permanent magnets ensure a flat field gradient is produced at the centre of the glass cell. The plates sit on tracks to allow adjustment of the 2D MOT's magnetic field gradient which is set to 17 G/cm.

The optics for the MOT are all from Thorlabs and sit in a combination of 3D printed holders and aluminium L-plates. Fig. 6.3a shows the optics layout for the 2D MOT, with light sourced from the combination optics as discussed in Sect. 5.4. The fibre collimators consist of a fibre plate and a 50 mm lens inside a lens tube for a $1/e^2$ radius of 5 mm to produce a large trapping volume. A single mirror reflects the light into the 3D printed holder for the non-polarising beamsplitter (NPBS) cubes which split the light into three regions. The first NPBS has a splitting ratio of 70:30 (reflection:transmission) with the transmission port going to the first MOT region, the second NPBS has a 50:50 splitting ratio with the reflection port creating the second MOT region; a gold coated prism reflects the remaining light to create the third MOT region. Each MOT beam goes through a $\lambda/4$ waveplate which is achromatic at 852 nm and 780 nm to create the required circular polarisation at the MOT. The MOT beams are retroreflected through a second achromatic $\lambda/4$ waveplate and a mirror. This optical system is implemented on two orthogonal axes to produce the 2D MOTs. The 2D MOT cooling light is red-detuned by $\delta = -4.9\Gamma$ with 60 mW at 852 nm for each axis for a

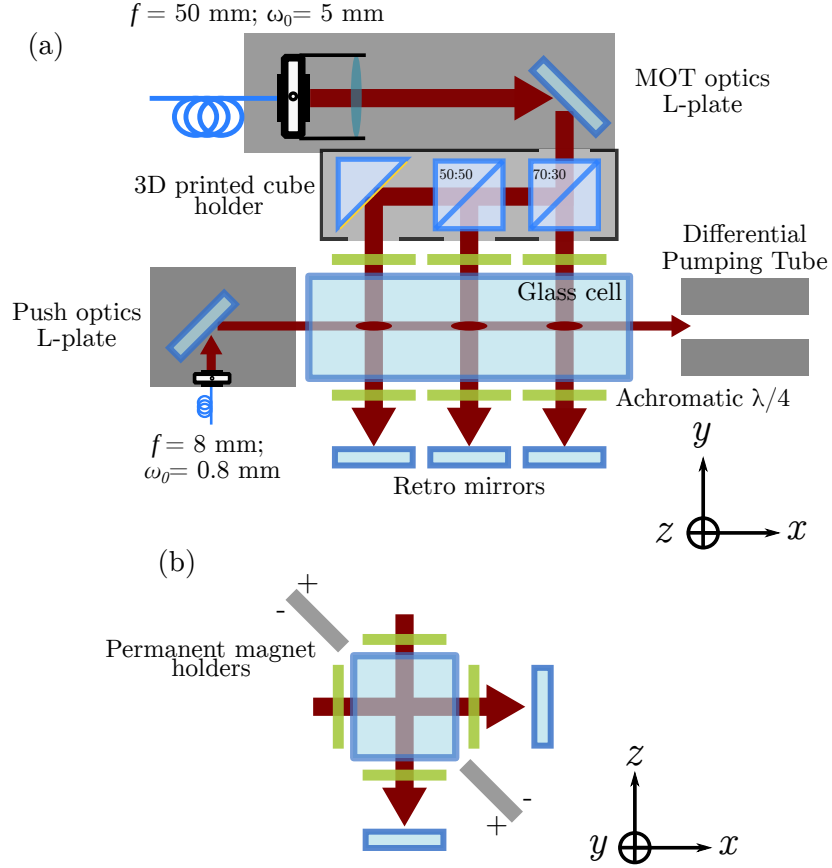


Figure 6.3: (a) Schematic of the 2D MOT optics and push beam. (b) Schematic of the permanent magnets orientation relative to the 2D MOT beams.

total of 20 mW per region beam. The total repump power is 2 mW per axis.

An L-plate to the rear of the glass cell holds the optics for the push beam. The fibre plate and 8 mm collimating lens sit in a lens tube and ensures a $1/e^2$ radius of 0.8 mm for the push beam. The fibre and lens tube sit in a kinematic mounts which, along with a single mirror, allows for alignment through the differential pumping tube.

The 3D printed material used in the 2D MOT section was tough black PLA which was found not to transmit in the near infrared. The lids for the cube holders were printed on a flat glass plate which left one face of the PLA with a glossy finish. The formation of the 2D MOT was disrupted when the glossy surface of the lids faced the NPBS. It was found that the glossy surface increased light reaching the second MOT region by 1.2%. When the lids were printed in a manner to avoid glossy surfaces, the

reflectivity was reduced to 0.8% (the same as no lid) which did not effect the formation of the 2D MOT.

To image the 2D MOTs we used a Raspberry Pi Module 3 camera without an IR filter. We chose a Raspberry camera for its small footprint and cost-effectiveness. The Raspberry Pi software¹ provides a web interface for monitoring the MOT but does not allow for accurate characterisation. The camera has successfully imaged fluorescence at 852 nm and 780 nm, and has imaged a Cs MOT and a Rb MOT.

6.1.1 Velocity Class Selection

After initial the setup of the vacuum system, we were unable to demonstrate loading of the 3D MOT from the 2D MOT. The MOT loading was limited in part by the differential pumping tube and cryogenic mount acting as a velocity filter for parabolic trajectories of hot atoms from the 2D MOT. Modelling presented in Appendix C suggests that if the 2D MOT is sat flat on the table, only atoms with a velocity of $\gtrsim 20 \text{ ms}^{-1}$ will reach the 3D MOT. Previous modelling, presented in Ref. [180], suggests the bulk of atoms from a 2D MOT have a velocity between 7 ms^{-1} and 20 ms^{-1} , which implies our system is unsuitable for capturing sufficient atoms. One solution would be switching from a push beam to a detuned dipole beam which would act as an atomic waveguide to ensure relevant velocity classes reach the 3D MOT trapping region, as previously demonstrated in Ref. [123]. This method was not feasible as it required major rework of the optical systems to implement.

Alternatively, introducing a tilt to the 2D MOT offers a method for ensuring slower atoms can reach the trapping region on a parabolic trajectory. Previous models suggest that for a range of beam intensities, a 10 G/cm quadrupole field gradient and $\delta = -3\Gamma$, a Cs MOTs capture velocities would be $< 18 \text{ ms}^{-1}$ [181, 182]. Therefore, the tilt of the 2D MOT should optimise for velocities around $< 20 \text{ ms}^{-1}$. Modelling presented in Appendix C suggests introducing a 2° tilt to the 2D MOT chamber ensured atoms in the velocity range $> 15 \text{ ms}^{-1}$ could reach the trapping region. We did not perform modelling for Rb as the system with the tilt was able to load the 3D MOT with no

¹Motion: <https://github.com/Motion-Project/motion>

further adjustments having optimised the angle for Cs.

Implementing a tilt to the system was a viable solution as the 2D MOT chamber and science chamber are connected via a flexible bellow. To introduce the tilt, an aluminium baseplate with a 2° gradient was machined and placed under the 2D MOT as indicated in Fig. 6.2a.

6.2 Cryogenic Stage

To achieve cryopumping², the cryogenic stage must reach suitably low temperatures to freeze-out background contaminants including, N₂, O₂, CO, He and H which contribute to the residual background pressure in the system. These gases cryopump at different temperatures with N₂, O₂, and CO sublimation temperatures between 20-30 K [183]. H₂ specifically can permeate the crystalline structure of a stainless steel vacuum system making it one of the main contributions to unavoidable vacuum leaks [184]. At 4 K (the temperature of liquid He) only H₂, and He are not cryopumped, however, these elements will lose kinetic energy when they come in contact with cryogenic surfaces. These elements are of lesser concern due to their smaller mass.

The cryostat used in this experiment was repurposed from work previously carried out as presented in Ref. [185] where it was used in a hybrid system with Rydberg atoms in tweezers for interfacing with a superconducting resonator. The cryo-cooler is a Sumimoto RDK-408D2 coldhead combined with a custom ultra-low vibration (ULV) interface developed by ColdEdge with He as its cryo-coolant. The cryostat has cooling powers of 0.81 W at 4 K, 5.68 W at 10 K and 12.8 W at 20 K at the second stage cooling system. The cold-head is fed He with a compressor (Sumitomo: F-70H). Fig. 6.4 shows a cross-sectional view of the cryostat and low vibration interface. The cryostat is a Gifford-Mahon (GM) cryostat which operates on the following principle [186, 187]: during operation, a rotary valve in the cold-head switches between the high pressure and low pressure gas lines fed from the compressor. When the rotary valve aligns with the high pressure gas line, the displacer moves up and the gas expands, isothermally

²Discussed in detail in Sect. 8.4.1.

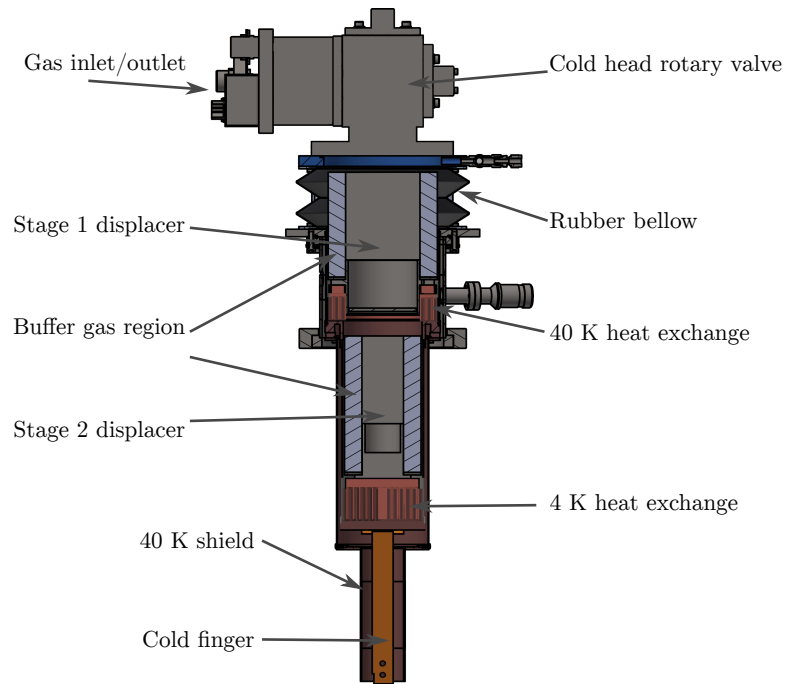


Figure 6.4: Cross-sectional view of the cryostat showing the atmospheric side and UHV.

cooling. When the rotary valve aligns with the low pressure line, the cold gas flows through the heat exchange and carries heat away from the cold stage. The cold-head switches between these two states continuously providing the cooling cycle for the cryostat.

It is paramount to minimise vibrational coupling between the science table and the cold-head as failing to do so would result in optical instabilities effecting the ability to maintain alignment of critical laser paths. To reduce this coupling, this cryostat is designed with a low vibration interface where the cold-head is mechanically decoupled from the top of the vacuum system. The science tables canopy suspends the cold-head above the UHV stage and the heat exchangers are physically separated. A rubber bellow connects the top of the vacuum system to the cold-head and the space between the rubber bellow and heat exchange is filled with a buffer gas of high purity He drawn from a separate canister. The partial pressure of the buffer gas region is 0.4 PSI. As the cold-head cools towards 4 K, the buffer gas starts to condense in the heat exchange. The buffer gas and condensed He is what draws heat from the cold finger and shield

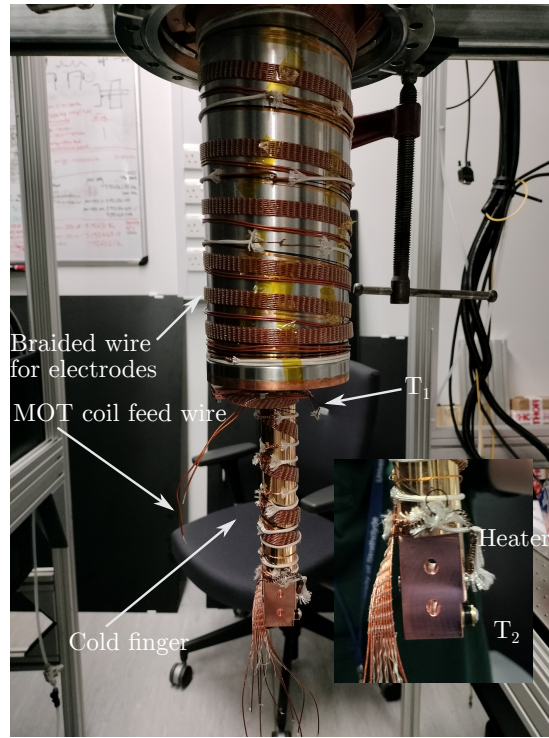


Figure 6.5: Photograph of the in-vacuum portion of the cryostat.

stage attached to the two heat exchangers. To further reduce vibrational coupling, the compressor is raised above the floor away from the science table.

The in-vacuum cryogenic stage consists of two parts. Fig. 6.5 shows a photograph of 4 K stage which consists of an oxygen-free high-conductivity (OFHC) copper cold finger that is attached to the lower heat exchange and extends to the centre of the science chamber. There are two temperature sensors, one just below the heat exchanger (T_1) and one at the bottom of the cold finger (T_2). Wires for electrodes and superconducting coils are wrapped around the second stage displacer and cold finger to ensure thermal coupling before reaching the copper mount discussed in Sect. 6.3. The stage connected to the upper heat exchanger completely shields the cold finger and operates at 40 K. The shielding is plated with a $1\ \mu\text{m}$ layer of gold for improved emissivity.

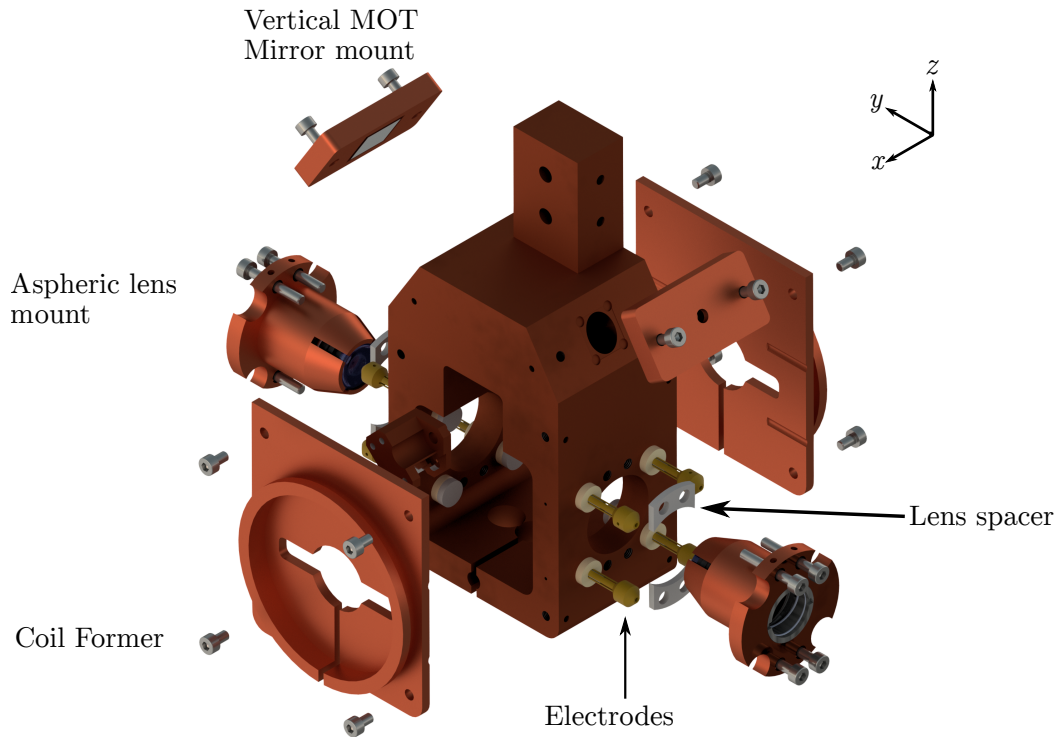


Figure 6.6: Exploded CAD render of the cryogenic experimental stage attached to the bottom of the cold finger for holding the aspheric lens mount, superconducting coil former, field-nulling electrodes and MOT mirrors.

6.3 Cryogenic Mount

The central component of this experiment is the mount at the base of the cold finger which houses all the equipment needed to project and load the tweezer arrays. Fig. 6.6 shows an exploded view of this stage, colloquially referred to as the ‘bird house’. The design derives from the SQuAre experiment’s in-vacuum lens mount, discussed in Ref. [26], with alterations made to account for the limited optical access due to the cryostat. The mount is custom machined from OFHC copper (S&D Non-Ferrous³) chosen due to the material’s high thermal conductivity for improved cryogenic performance. Previous tweezer array experiments with cryogenic systems have opted to use copper beryllium (Cu-Be) alloy which has improved ductility and reduced electrical conductivity, the latter of which is important when minimising Eddy currents induced through coil switching [9, 121]. We chose not to go with Cu-Be due to a matter of cost,

³<https://www.sdnferrous.co.uk/>

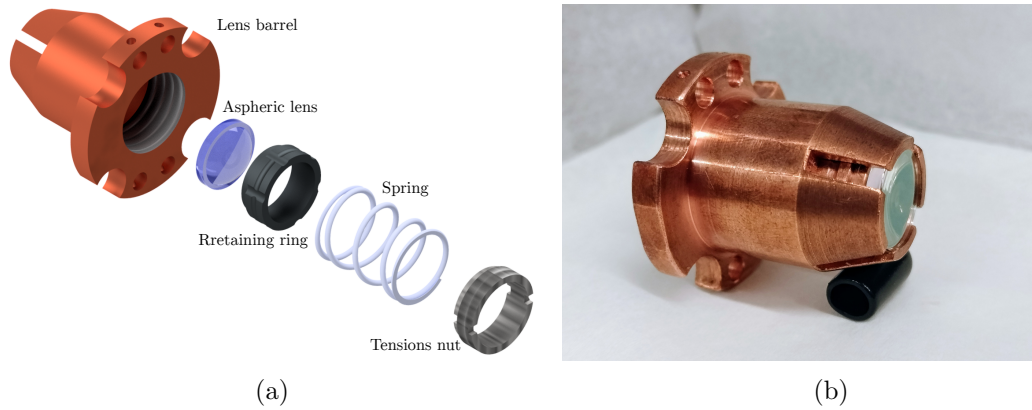


Figure 6.7: (a) An exploded CAD rendering of the lens mount. (b) Photograph of the lens mount before gold coating.

and instead chose to put a cut through the bird house to break the free path for Eddy currents to form.

In-vacuum high-NA aspheric lenses are used to focus down the trap beams and to perform fluorescence imaging. The lenses reside in OFHC copper mounts and are attached to the bird house. Fig. 6.7 shows an exploded CAD view and a photograph of the assembled lens mount. The aspheric lens' (Geltech 355561 [188]) focal length is ~ 10 mm with an $NA = 0.45$ at 852 nm. The flat side of the lens is coated with a bilayer layer of Nb_2O_5 and highly conductive aluminum-doped zinc oxide (AZO) (PlasmaQuest) which stops the build up of charge near Rydberg atoms. The conductivity of the coating increases as the coating thickness increases, but this will also reduce the optical transmission and, hence, increase the heat load on the cryostat, and could cause a time-dependent thermal lensing⁴. To balance these, we chose an AZO coating thickness of 124 nm and a Nb_2O_5 coating thickness of 94 nm which provided a transmission of $> 90\%$ at 750 nm - 860 nm and $> 85\%$ at 940 nm with a sheet resistance of $< 250 \Omega/\square$. Any residual charge will be cancelled by the 8 electrodes around the lens mounts. A retaining ring and spring (Lee Spring: LC 049HJ 03S316) holds the lens flush to the front of the lens mount with tension on the spring created by an externally threaded nut. The spring ensures the lens stays flush with the mount during the contraction of the copper and glass during the cooldown. Before being placed in-vacuum

⁴No evidence of thermal lensing was observed while operating the arrays discussed in Chapter 8.

and cooled to cryogenic temperatures, the precise separation of the lens must be set. This is exacerbated when accounting for the thermal contraction of the copper mount, glass lens and the change in refractive index of the lens. The thermal contraction of the copper components can be characterised by the fractional length change

$$\Delta L/L = (L_{293\text{ K}} - L_{\chi}) / L_{293\text{ K}}, \quad (6.1)$$

where $L_{293\text{ K}}$ is the length of copper at 293 K and L_{χ} is the length of copper at some temperature χ [189]. At cryogenic temperatures, the lenses must be separated by 14.068 mm. For copper at ~ 5 K, a $\Delta L/L = 0.3276\%$ is expected⁵ and implies a room temperature separation of 14.114 mm, not accounting for changes in lens shape or refractive index of the glass. The lenses are made of dense lanthanum flint glass and have a room temperature refractive index⁶ of 1.79 at 852 nm. Accounting for the Sellmeir coefficients of the glass, the expected change in refractive index from 300 K to 4 K is $\Delta n = 1.924 \cdot 10^{-3}$ assuming dn/dT is constant with temperature.

The focal length of the lens pair at room temperature was modelled in Zeemax [190] for an 852 nm collimated input beam. In this configuration, the beam should focus at ~ 2 m behind the second lens. The spacing of the lenses was set out of vacuum by incrementally shaving material off the OFHC copper lens spacers (shown in Fig. 6.6) using a lathe. The lenses were placed in the bird house with the lens separators in place, and a beam profile was taken of a collimated laser beam at 852 nm after it had passed through the lens pair. Zeemax was used to determine how much material should be removed from the spacers, and half the required amount was removed to avoid an overshoot. This process was repeated until the beam profile converged on the focal distance corresponding to the desired separation.

Larger field of views (FOVs) can be achieved by using a high NA microscope objective as has been demonstrated both in room-temperature [8, 74] and cryogenic [122, 123] systems. A microscope has not yet been implemented directly on a cryostat which would prove a large engineering challenge. Initial experiments in our system do not require

⁵From page 65 of <https://www.osti.gov/servlets/purl/4219483>

⁶From <http://www.cdgmgd.com/webapp/pdf/D-ZLaF52LA.pdf>

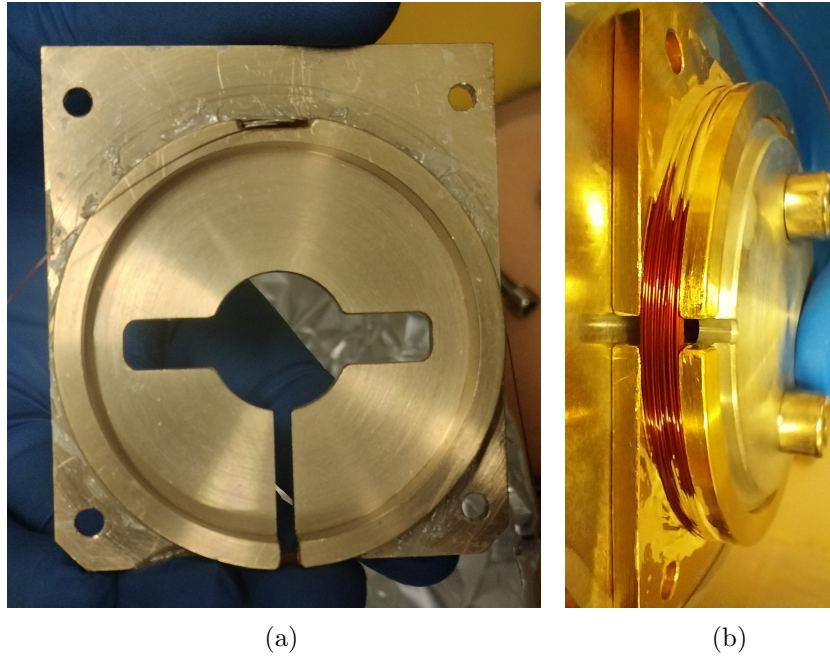


Figure 6.8: (a) Photograph of the superconducting coil former with the cut for optical access. (b) Photograph showing the superconducting wire at the cut to stop free paths for Eddy currents; epoxy can be seen away from the cut.

large arrays, so the aspheric lens provides a suitable choice for initial characterisation of cryogenic Rb and Cs arrays, and interspecies Rydberg interactions.

A MOT requires large quadrupole fields with gradients across the trapping volume in the 10 – 20 G/cm range. To produce such fields, one can have small coils with relatively low currents in close proximity to the trapping region, or large coils with higher currents further way from the trapping region. The latter may require water cooling which increases complexity and reduces optical access. Following the design work presented in this thesis, a team demonstrated an out of vacuum coil system for a cryogenic tweezer array experiment by mounting the coils in a dimpled viewport[123]. This design maintained good optical access while keeping the coils at close separation. Superconducting wire becomes an option at cryogenic temperatures and allows for small coil separation with low additional heating of the cryostat. This has previously been demonstrated in the first cryogenic tweezer array experiment [9, 121]. As our expected base temperature will be < 10 K (so lower than the critical temperature T_c of many

common superconducting materials), we chose to implement superconducting coils in our system.

The extreme bend radius for our coil mount makes high temperature superconductors ($T_c > 77$ K) a poor wire choice due to reduced performance under stress [191]. Niobium Titanium (NbTi) ($T_c = 9$ K) and Niobium Tin (Nb₃Sn) ($T_c = 10$ K) are two common low temperature superconductors with wide availability [189]. Due to the better mechanical properties of NbTi, we chose this material for our MOT coils.

We chose wires with 54 filaments of NbTi in a copper cladding with a Cu to NbTi ratio of 1.3:1 (Supercon Inc: SC-54S43-0.229mm-HML). The wire is insulated with heavy Polyimide enamel with total wire diameter of 0.254 mm. Each coil has 93 turns (14 axial by 7 radial), and has an internal radius of 30 mm. Each coil uses 24 m of wire with an additional 1 m of wire for thermal coupling. The coils are wound to a former, shown in Fig. 6.8a; the former is made from OFHC copper and has a slit to ensure optical access for horizontal MOT and Rydberg beams, and to allow the 3D MOT to receive atoms from the 2D MOT. The two coil formers are attached to the bird house as shown in Fig. 6.6. The coils have an internal separation of 52 mm to ensure a uniform field gradient of 10 G/cm when supplied with a 1.3 A current.

A silver based two part epoxy (Masterbond: EP21TDCS-LO) is used to maintain wire tension. This epoxy has shown operation at a range of temperatures (from < 40 K to > 450 K) and vacuum pressures ($< 10^{-10}$ mbar to 10^{-2} mbar) in one experimental system [192]. The proven operation of this epoxy over a range of temperatures will ensure the coils will not unspool as material deforms during multiple warm-ups and cool-downs of the cryostat. Fig. 6.8b shows the superconducting wire after winding to the former. The epoxy is conductive, so is not applied at the cut in the former. Thermal coupling of the superconducting wires is performed by winding it about the cold finger as shown in Fig. 6.5. About halfway up the cold finger, the wires are barrel connected to 1 mm thick Kapton-coated copper wire. The copper wire is thermally anchored to the upper stages of the cryostat. An electrical feed through at the top of the vacuum system allows connection to an external coil driver. To reduce thermal coupling to room temperature surfaces, the two coils are connected together in an anti-Helmholtz

configuration on the cold finger itself.

Alternative designs for MOTs in tweezer experiments have been demonstrated in recent work where the cooling stage is performed elsewhere and a transport beam is used to bring the atoms between the lenses [193–195]. The addition of an atom transport system within a cryogenic system would remove the need to have quadrupole coils near the OFHC copper, eliminating Eddy currents induced by magnet field switching. However, such a setup would have increased optical complexity.

One factor which can limit the base temperature of a cryostat is radiative heating from room temperature surfaces surrounding cryogenic components. This effect can be mitigated by implementation of a shield which completely surrounds 4 K stage. Shielding would also reduce the black-body radiation (BBR) light shifts which can increase the Rydberg lifetimes by a factor of < 4 depending on the chosen Rydberg state [124, 196]. Previous cryogenic tweezer array experiments have implemented shielding [9] which offered enhanced lifetimes, however, limited optical access and required reengineering to implement Rydberg beams [122]. We chose to forgo a full shield and kept the previously installed 40 K shielding which covered the cold finger but not the copper mount. To increase the emissivity, we applied a 1 μm thick gold-plating to the main copper mount, lens mount, mirror mounts and coil former. To limit the BBR light shifts, we designed the mount to limit the solid angle of room temperature surfaces facing the atoms.

6.4 3D MOT Optics

For maximal trapping volume in a 3D MOT, it is typical to use three pairs of orthogonal counter-propagating beams⁷. However, this is not possible in our system due to limited optical access caused by the separation of lenses and position of the cryostat above the MOT which rules out beams entering from the top of the chamber. As a result, the 3D MOT has a total of 3 beams, a pair of counter-propagating horizontal beams and a single vertical beam which makes two passes across the MOT before being retroreflected

⁷In reality, a N dimensional MOT requires a minimum of $N + 1$ beams to perform cooling and trapping [197].

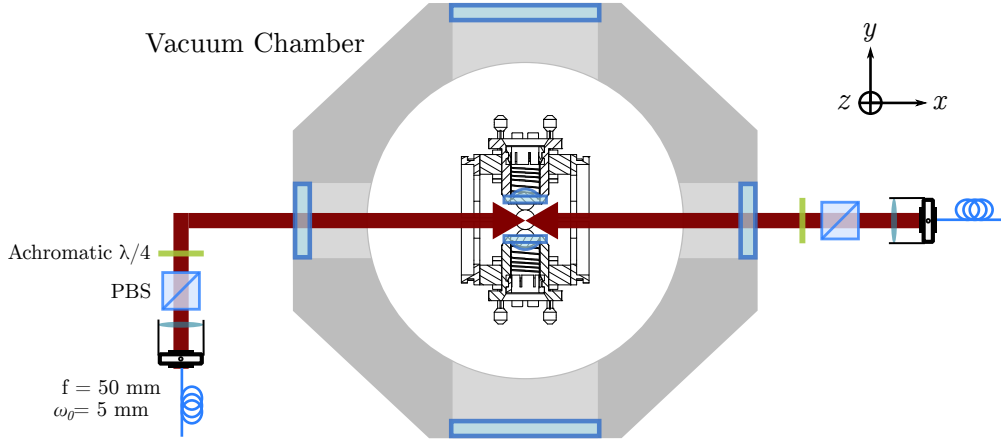


Figure 6.9: Schematic of laser path for the horizontal MOT beams.

back onto itself. The horizontal MOT beams pass through the trapping region at a right angle to the lenses as illustrated in Fig. 6.9. The light is sourced by fibre from the cooling table discussed in Sect. 5.4. The fibre collimator consists of the fibre plate and a 50 mm lens inside a lens tube for a $1/e^2$ radius of 5 mm. Each beam passes through a PBS to ensure good p-polarisation. Circular polarisation for the MOT is set by an achromatic $\lambda/4$ waveplate before the light enters the chamber.

The vertical beams can only enter the vacuum chamber from underneath and cannot be retroreflected inside the chamber as this would require a waveplate in-vacuum likely attached to the cryogenic stage. Instead, after the light enters the chamber, it makes a two passes across the trapping region as illustrated by the red arrows in Fig. 6.10. Between the first and second pass, the beam is reflected off two silver coated mirrors (Thorlabs: ME05S-P01) which ensure the correct handedness of polarisation for each pass across the MOT. After the second pass, the light exits the chamber and is retroreflected by a mirror mounted on a piezo-actuated kinematic mount. Each reflection off the internal mirrors will degrade the beams circularity due to differential reflectivities between s- and p-polarisations. Results in Appendix D suggests the degradation is small enough that it should not limit this experiment's ability to load a 3D MOT. For the optical alignment of the input beam, the fibre collimator and mirrors are placed on a translation stage situated below the vacuum chamber, a schematic of this is shown in Fig. 6.11a. Fig. 6.11b shows a schematic of the optics on the translation stage. The

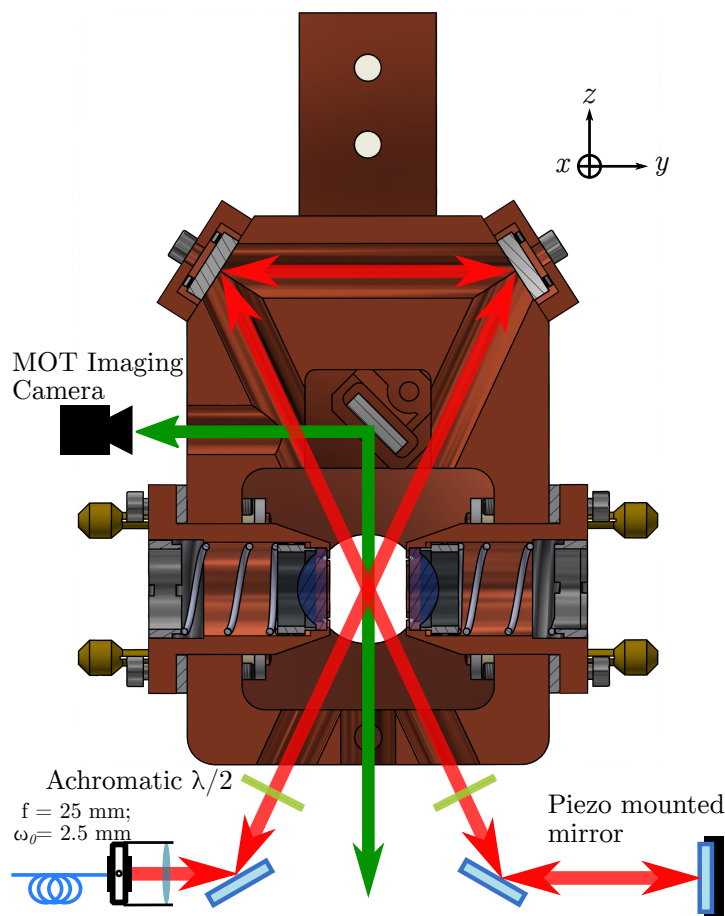


Figure 6.10: Vertical slice the cryogenic mount with MOT beam paths indicated by the red arrows; an additional beam path is indicated in green.

beam is collimated with a 25 mm lens for a $1/e^2$ radius of 2.5 mm. The PBS ensures the beam has good p-polarisation. A glass wedge between the PBS and fixed mirror redirects some light into a photodiode for power monitoring. The fibre collimator and final mirror are attached to kinematic mounts for aligning the beam through the chamber.

6.5 Vacuum Assembly

The first stage of the vacuum build involved dismantling of the old system [198] and repurposing its main chamber and cryostat. Working on the cryostat required its suspension between the two optical tables, a process described in Ref. [185]. The copper mount was assembled, shown in Fig. 6.12a, and attached to the cold finger,

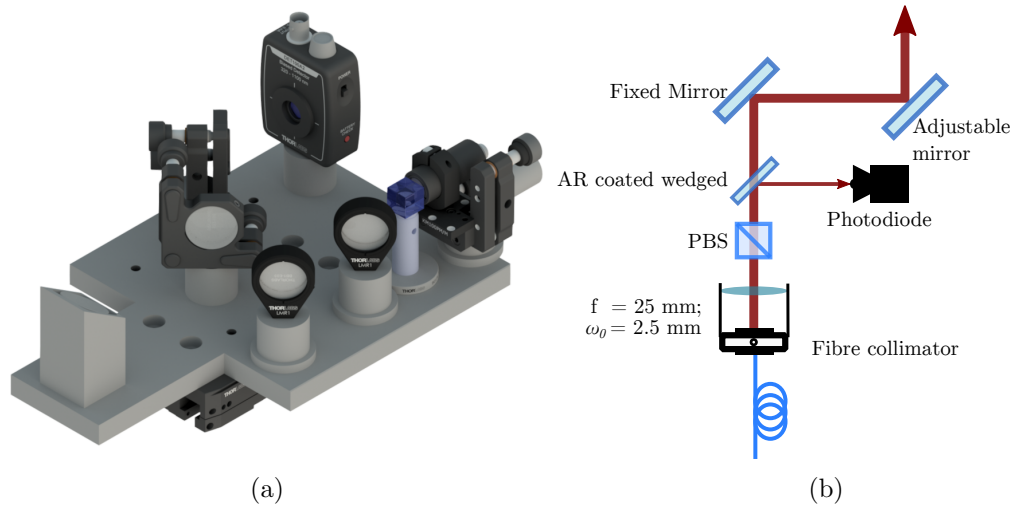


Figure 6.11: (a) CAD rendering of vertical fibre launch translation stage. (b) Schematic of optics on translation stage.

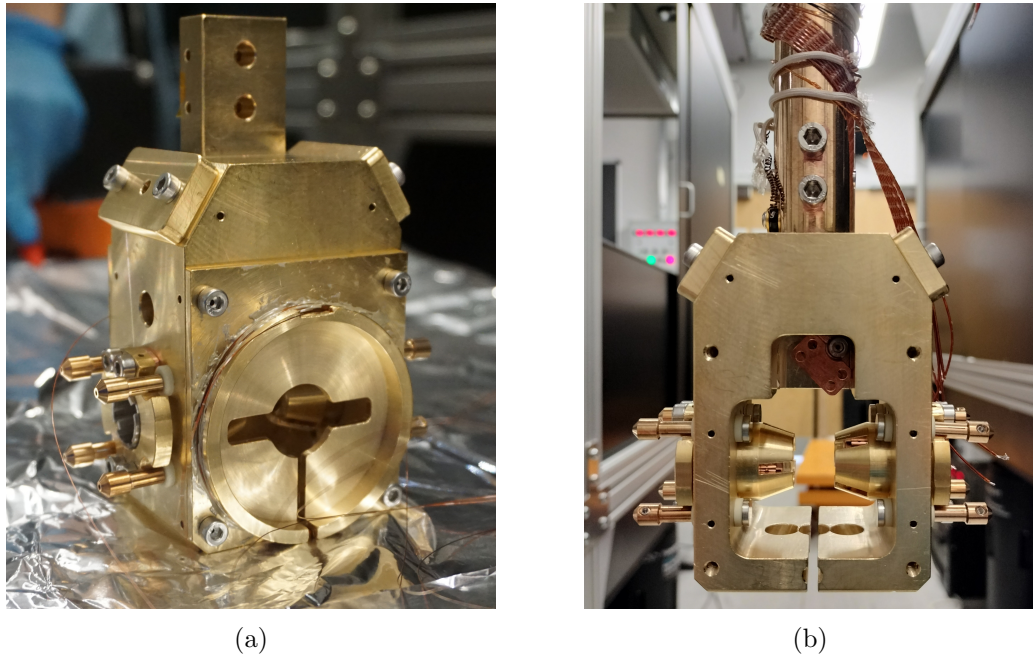


Figure 6.12: Gold coated OFHC copper mount assembly (a) and attached to cryostat without the superconducting coil former (b).

shown in Fig. 6.12b. The 8 electrodes for electric-field nulling were connected to the copper mount by the thermally anchored ribbon cable. The superconducting coils were connected in an anti-Helmholtz configuration, verified by applying a small current and measuring the resulting field with a gauss meter. The main chamber was attached after the cryogenic stage was complete. The coldhead was removed from the vacuum system to prepare for the bake process. After assembling the science chamber, it was placed on the main optical table where the 2D MOT was attached.

After the build, a turbo pump brought the vacuum system to a pressure of $1.7 \cdot 10^{-6}$ mbar. The system was then baked at a temperature of $100 - 150$ °C for 14 days. After cooling down from the bake, the system reached a base temperature of $2.6 \cdot 10^{-8}$ mbar. The coldhead was reattached and the ion pumps for the science chamber and 2D MOT chamber were turned on which reached a base pressure of $3.8 \cdot 10^{-10}$ mbar and $2.2 \cdot 10^{-9}$ mbar respectively.

After a few months of operation, the system was partially rebuilt due to challenges in observing atoms loading from the 2D MOT into the 3D MOT. Due to the restricted optical access of the cryostat setup, dispensers were installed in the science chamber to enable independent alignment and verification of the 3D MOT and 2D MOT. The two chambers were isolated with the gate valve before the science chamber was brought up to atmosphere. The opportunity was taken to also install the TSP between the 75 L/s ion pump and science chamber. Following the rebuild, the science chamber was baked for a second time. We chose to perform this second bake with the cold head in place, to avoid removing all the optics that had subsequently been assembled around the vacuum chamber. This limited the bake temperature near the cold head to < 80 °C. We chose to heat the top vacuum system around the coldhead to 45 °C and ion pump to 100 °C for a total of 10 days. The pressure on the turbo pump at the start of the bake was $2.5 \cdot 10^{-6}$ mbar. The base pressure after the bake was $1.7 \cdot 10^{-8}$ mbar. After switching on the ion pump, a base pressure of $1.2 \cdot 10^{-9}$ mbar was achieved. The TSP was then fired for 5 rounds⁸ after which a base pressure of $3.7 \cdot 10^{-11}$ mbar was reached. Fig. 6.13

⁸Each round consisted of firing the TSP at 48 A for a duration of 60 s a total of 5 times with 10 minutes between each firing. Each round was performed with the 75 L/s ion pump turned off.

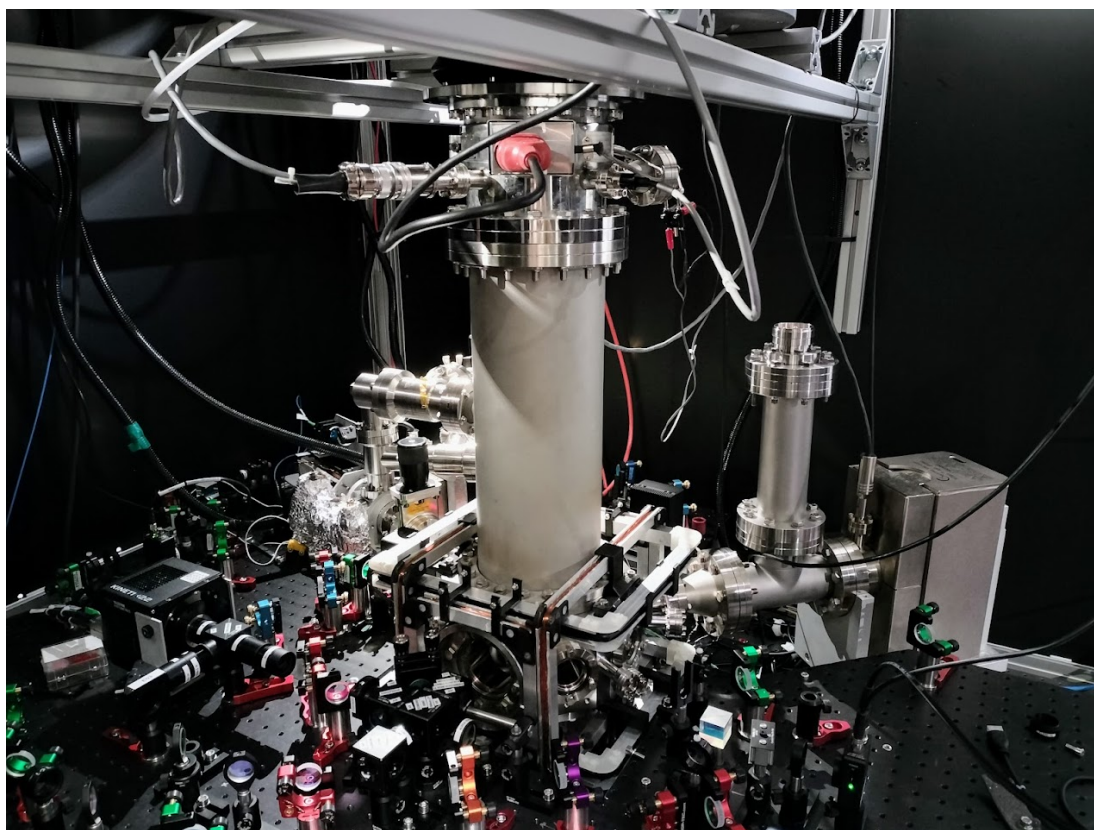


Figure 6.13: Complete vacuum system after the second build.

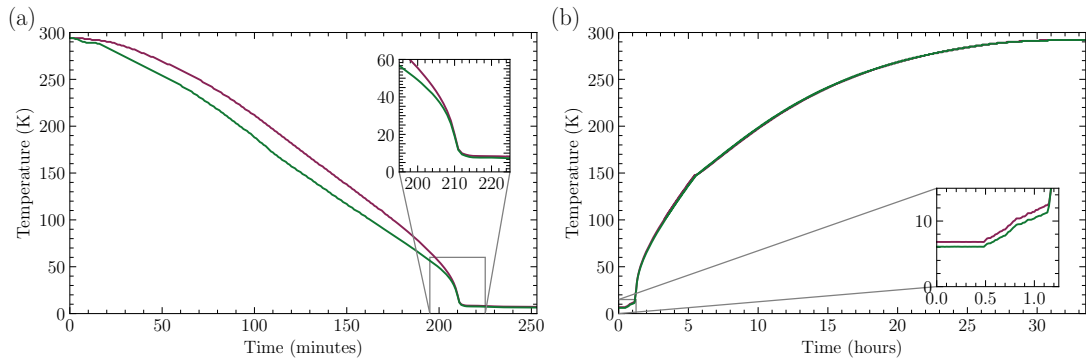


Figure 6.14: Measurements of the cryostat cooldown (a) and warm-up (b) for the upper temperature sensor (green) and lower temperature sensor (red).

is a photo of the complete vacuum system following the second build of the experiment.

6.6 Cryostat Performance

In preparation for bringing the system to cryogenic temperatures, the cryostat is connected to the compressor via feed lines. The low vibration interface must be purged of air to ensure ice does not form in the two heat exchangers. The purge is performed by connecting a 200 bar cylinder of high purity He and inflating the bellow with a partial pressure of 0.4 PSI. The release valve at the top of the interface is then opened to deflate the bellows. This inflation/deflation step is repeated a total of 10 times. After the purge the bellow is inflated a final time and the compressor is turned on. Measurements of the cryostat cooldown are presented in Fig. 6.14a and recorded with a temperature controller (Lakeshore: Model 335) connected to the two sensors on the cold finger. It takes fewer than 4 hours for the system to reach a base temperature of 6.1 K at the upper temperature sensor and at 6.8 K for the lower temperature sensor. Clear evidence of cryopumping was apparent when the temperature would reach ~ 30 K when the pressure reading on the 75 L/s ion pump rapidly falls from the 10^{-10} mbar range to $< 1.0 \cdot 10^{-11}$ mbar. This base temperature could be improved by implementation of a cryogenic shield as discussed in Sect. 6.3, however, ~ 7 K is low enough to take advantage of the cryopumping effect. Following the installation of the TSP, prior to the cooldown, the ion pump read a base-pressure of $1 \cdot 10^{-11}$ mbar at room temperature.

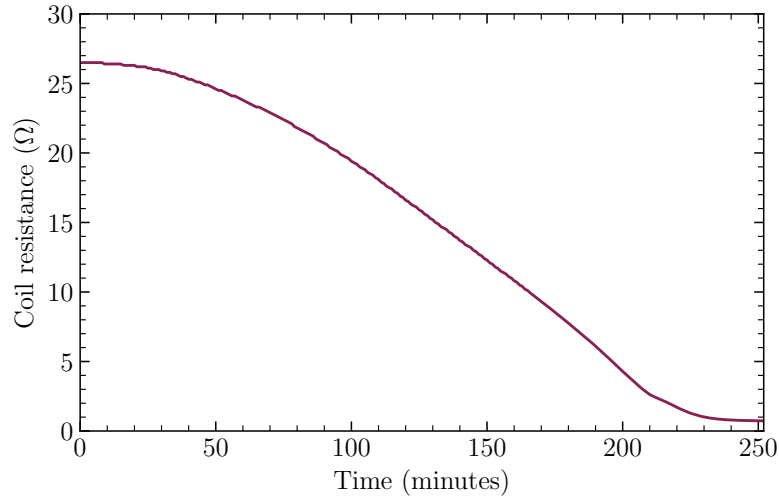


Figure 6.15: Resistance of the in-vacuum coil chain during the cooldown of the cryostat.

After the cooldown, the system remained below the ion pump's base reading for several months before observing any increase due to saturation of the cryogenic surfaces.

Fig. 6.14b shows temperature traces for the warm-up of the cryostat. Due to the build up of sublimates on cryogenic surfaces, it is important to start the warm-up process slowly, otherwise a spike in pressure may choke the ion pump. To do this, we step the heater over the course of ~ 1 hour till the temperature reaches ~ 12 K; this will release some of the impurities off the cryogenic surfaces. Once the ion pump brings the pressure back down, we turn off the heater and compressor simultaneously and the cryostat starts to warm-up. To aid the warm-up process, the heater is turned back on once the temperature reaches ~ 30 K. It takes the cryostat ~ 30 hours to reach room temperature.

6.6.1 Superconducting Coils

Fig. 6.15 shows the measured resistance of the in-vacuum portion of the coil wiring during the cryostat cooldown. At room temperature, the resistance is 26.5Ω which falls to $730 \text{ m}\Omega$ at the base temperature, due to the associated increase in conductivity of copper. To check for evidence of superconductivity, the cryostat heater was slowly

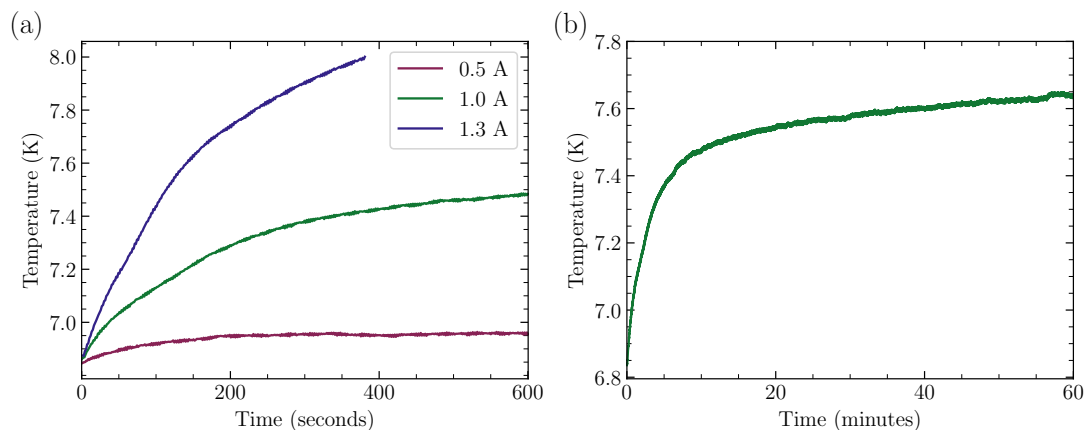


Figure 6.16: Temperature measurements of the cryostat for coils under load for 10 mins (a) and 1 hour (b).

raised to 12 K. No discontinuity in resistance was observed. The expected jump in resistance should be of the order $10 \text{ m}\Omega$, and the method is insufficiently sensitive due to the large length of standard copper wire which also changes resistance over these temperature scales. A direct 4-point wire measurement of the coil would be necessary to achieve this sensitivity.

To see if the system could perform under load, we carried out a series of tests in which the current passed through the coil was incrementally increased. Fig. 6.16a shows the affects applied currents had on coil temperature. As $T_c \approx 9 \text{ K}$ for NbTi, we chose an upper limit of 8 K above which the current through the coils would be cut. This value was chosen to provide headroom as we have no temperature sensor at the coils, only at the top of the bird house. The target current of the coils for a MOT is 1.3 A; this current heats the coils to 8 K in just over 6 minutes. This timescale is acceptable as the coils are only supplied current for a short time during a tweezer experimental sequence. When looking for a MOT however, we would want to run the coils CW. If the coils are supplied 1 A (8 G/cm field gradient) the green trace in Fig. 6.16a indicates the coils remained $< 8 \text{ K}$ for the entire 10 minutes of mesurement. Fig. 6.16b shows the coils with 1 A through them for a total of 1 hour, and shows that we can run for considerable time staying below the 8 K limit.

Part III

Towards Cryogenic Tweezer Arrays

Chapter 7

Cryogenic Caesium MOT

This chapter will discuss the caesium magneto-optical trap (MOT) of this experiment which forms the first step required for loading single atoms into traps. This system consists of two MOTs, a 2D MOT residing in the high pressure side of the vacuum chamber and a 3D MOT which forms at the cryogenic stage within the ultra-high vacuum (UHV) science chamber. The 2D MOT is an improved design from the one built for the Square project [26] with alteration made in regards to the optics surrounding the chamber. The 3D MOT could not follow the previous experiment's design because of the reduced optical access from above the chamber due to the cryostat. Optical access for MOT beams and for imaging were further reduced as the cryogenic mount was also designed to take advantage of the effects of cryopumping.

This chapter will discuss the challenges involved in loading the 3D MOT due to the system design constraints. We will then discuss MOT optimisation after the challenges of loading were overcome. The polarisation gradient cooling (PGC) stage is then discussed which is the final step in cooling before loading the Cs tweezer array — discussed in the next chapter. We will also discuss additional challenges presented in regards to the response of cryogenic copper to quadrupole coil switching which will require additional controls to be implemented for future operation of the system.

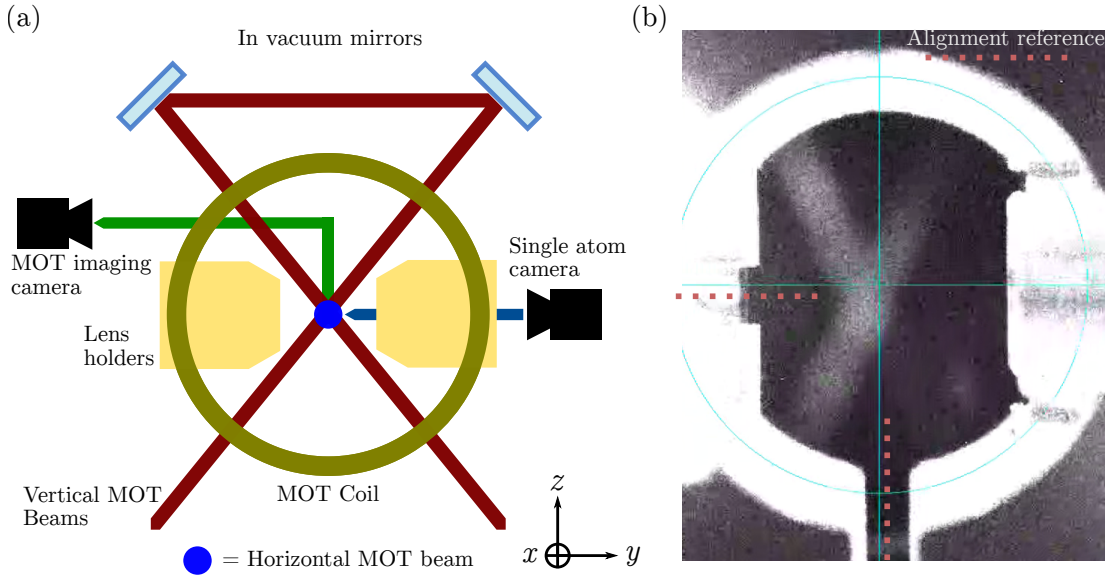


Figure 7.1: (a) Schematic of the MOT beam alignment relative to the lenses and coils. (b) Fluorescence image of the vertical 3D MOT beams with markers indicating alignment guides (lens centre: vertical; and cryogenic mount cut: horizontal) for setting the beam overlap position.

7.1 MOT Alignment

As discussed in Sect. 6.3 the 3D MOT is designed to form at the centre of the cryogenic stage at the bottom of the cold finger. Due to the restricted optical access of the high numerical aperture (NA) lenses, a pair of MOT beams with a waist of 2.5 mm cross 25° from the vertical axis in the mid plane of the lenses, while an additional beam with waist 5 mm crosses through the coil axial axis. A schematic of this is shown in Fig. 7.1a. The beam alignment was initially performed using the mechanical constraints of the chamber, with the horizontal MOT beams aligned with a pair apertures attached to the view ports on either side of the main chamber, and the vertical beams aligned by maximising transmission through the chamber after two bounces off the internal mirrors.

After performing the alignment using the mechanical constraints and verifying operation of the 2D MOT, we attempted to load a caesium 3D MOT without success. This persisted even after verifying the beam polarisation relative to the quadrupole

field orientation with a polarimeter and Gaussmeter, and verifying the beam alignment of the 3D MOT, the 2D MOT and push beam.

The cause of the issue was found to be optical alignment due to the cold finger not running concentric to the vacuum system's large reducing tee. As a result, the copper mount sat off centre in the science chamber with a slight angle. This meant that the mechanical references were no longer valid leaving the 3D MOT beams poorly overlapped at the centre of the copper mount.

To overcome these challenges, two alterations were made to the 3D MOT. The first change was to temporarily switch from two counter-propagating horizontal MOT beams to a single retro reflected beam. This would reduce the complexity in optical alignment while looking for a MOT. In this configuration the horizontal axis will become prone to standing waves, but given these are the larger MOT beams, this effect can be mitigated by offsetting the retroreflecting mirror alignment. The experimental data from the rest of this thesis was recorded with the 3D MOT's horizontal beam in the retro configuration. Long term, the system will return to two independently aligned counter-propagating beams to allow state selective readout with a blow away beam and for optical access of the Rydberg lasers.

The second alteration made to the experiment was the addition of dispensers to the science chamber as discussed in Sect. 6.5. One dispenser was installed for each species. This reduced the number of state space dimensions to explore by removing the parameters associated with the 2D MOT. The dispensers were used to guarantee the presence of atoms within the 3D MOT overlap volume, enabling optimisation of beam alignment, and to position the MOT at the centre of the lenses via magnetic field nulling. Once these checks were performed with dispensers, we would move on to loading atoms exclusively from the 2D MOT. To optimise the optical alignment, the vertical MOT beam was reduced in size with an iris to improve the contrast of the atom fluorescence when a camera is set to high exposure, as shown in Fig. 7.1b. The fluorescence can be used to walk the vertical MOT beams roughly between the lenses. The top-down MOT (the green path in Fig. 7.1a) camera is used to align the vertical MOT beams on the horizontal axis. After optical alignment, a Cs MOT was

quickly achieved with the dispensers and finally, it was shimmed between the lenses in preparation for loading atoms from the 2D MOT.

7.2 MOT Loading

In the absence of a push beam, atoms will disseminate from the 2D MOT's non-trapping axis as a thermal distribution and only a small subset of atoms will arrive at the 3D MOT within a trappable velocity class. The inclusion of a push beam set to the ideal detuning provides a momentum kick to some atoms travelling in the wrong velocity class and will maximise atoms that are cooled in the 3D MOT [199]. Alignment of the push beam is typically done by ensuring it passes through the 3D MOT and exits the chamber via a view port opposite the 2D MOT. However, due to the addition of the tilt described in Sect. 6.1.1, our push beam scatters off the front face of the cryogenic mount which complicates alignment. Ensuring that the push beam passes through the differential pumping tube without scattering is a slow process and involved visually inspecting the scatter off the cryogenic mount until the beam looked as though it had passed through without scatter.

Despite the alignment of the push beam, the 3D MOT was not visible on the top-down camera, however, a MOT was visible through the in-vacuum lenses by using the camera nominally set up for single atom imaging. The top-down axis is imaged through an 8 mm hole on the side of the copper mount which opens 62 mm from the trapping region¹. The small aperture for this imaging axis was chosen to maximise cryogenic surfaces facing the atoms, but consequently results in limiting the solid angle for imaging. As such, for a weak MOT, few fluorescence photons will reach the camera sensor. When accounting for reflection off the internal mirror, the solid angle of the copper mount aperture², transmission of the camera lens, and the quantum efficiency (QE) of the camera, the photon collection efficiency was found to be $\sim 0.04\%$ — assuming isotropic scattering of light by the atoms in the MOT.

The camera counts on the single atom imaging axis can still be used to optimise

¹See Fig. 6.10 in Sect. 6.4.

²The copper mount itself was found to limit the light cone reaching the camera.

the MOT, but this cannot be a proxy for atom number as we are only imaging the MOT with a limited depth of focus. Parameters of the 2D MOT, 3D MOT and push beam were optimised for brightness using the single atom camera. In an optimised regime, the MOT was visible on the top-down camera but with a poor signal-to-noise ratio. The next optimisation performed was the positioning of the 2D MOT vacuum chamber. This was done by making small adjustments in angle or lateral positions of the chamber and comparing the relative brightness of the MOT. These adjustments improved the signal-to-noise on the top-down camera enough to begin optimisation by converting camera counts to atom number.

Atom number optimisation was performed on the 3D MOT with 5 mW of cooling power in the horizontal axis, 3 mW of cooling power in the vertical axis, 1 mW of repump power on the vertical axis and 50 μ W of repump power on the horizontal axis. The push beam operates at 350 μ W. The atom number was determined by fluorescence imaging [200, 201] on the top-down axis with a CCD camera (The Imaging Source: DMK33GX273). Alternatives, such as absorption imaging [202] were not feasible due to the limited signal on the top-down axis. Optimisation was performed with the superconducting coils operating at 1.3 A (10 G/A). The optimal detunings of the push, 2D MOT and 3D MOT were found to be -5.6Γ , -4.9Γ , and -4.0Γ respectively.

The MOT loading time constant τ can be extracted by fitting the atom number as a function of load duration t [203]

$$N(t) = N_{\text{sat}} (1 - \exp[-t/\tau]), \quad (7.1)$$

where N_{sat} is the atom saturation number. The optimised 3D MOT loading curve is in Fig. 7.2 which saturates at $N_{\text{sat}} = (3.78 \pm 0.03) \cdot 10^5$ with load time constant $\tau = 1.40 \pm 0.1$ s. The upper bound of N_{sat} is limited by the MOT geometry which takes the form of a capsule shape within beam overlap due to the restricted optical access of the cryostat. The MOT loading rate could be improved by implementation of a far detuned dipole beam which could act as an atomic wave guide for enhanced atom flux [123]. The inset image in Fig. 7.2 shows an image of a 3D MOT of Cs where

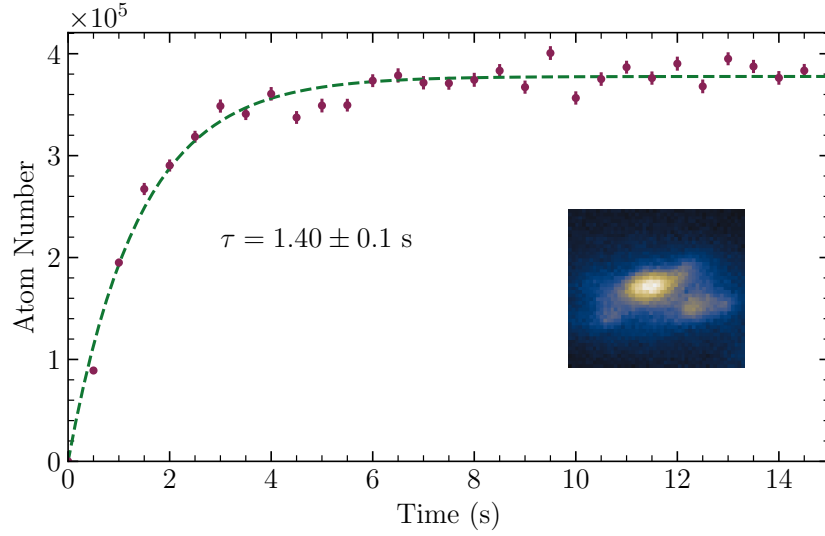


Figure 7.2: 3D MOT loading curve when receiving atoms from the 2D MOT.

most of the atoms are focused on a round lobe, but fringing is apparent which is due to the horizontal MOT beams being retro reflected and producing standing waves. This did not affect the ability to load traps. Switching to a two counter-propagating beams should remove the fringing.

7.3 Coil Induced Eddy Currents

The coil switch-on time should be sufficiently fast so that the rate of change of magnetic field is minimised otherwise the centre of the quadrupole field will shift in position as the MOT grows. There is also a risk that the current (and magnetic field) will ring if the switching time is too fast because of the coils' complex impedance. To ensure sub-millisecond switching times, without coil ring, we implemented a fast switching circuit with current feedback [204]. Despite this precaution, the 3D MOT formed above the trap region and moved between the lenses over a timescale of about a second as shown by Fig. 7.3b. This is indicative of a time dependent magnetic field with the likely cause being Eddy currents induced in the copper mount during the coil switching.

As discussed in Sect. 6.3, a cut was made through the copper mount through the

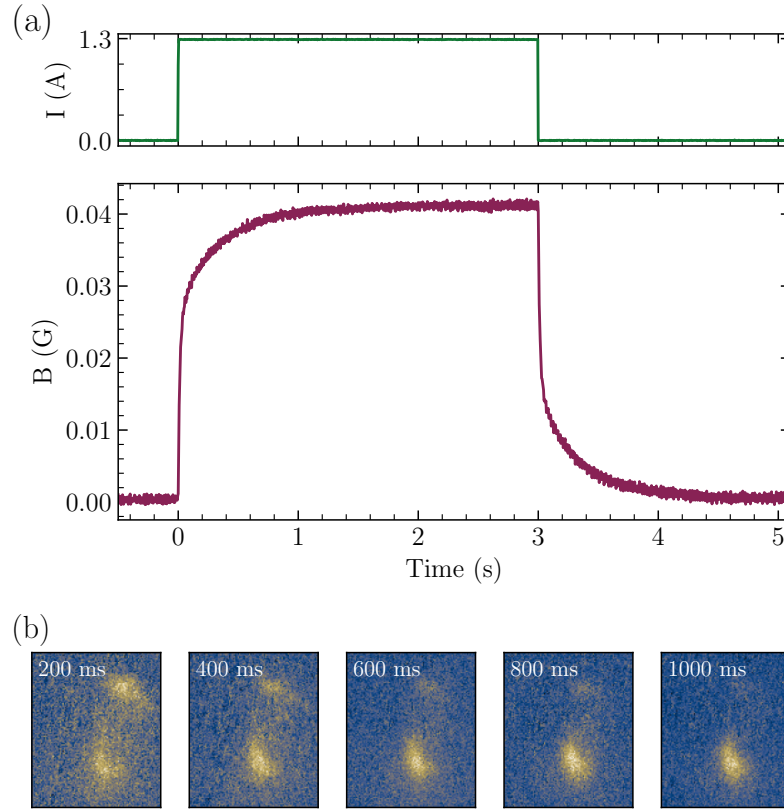


Figure 7.3: (a) Upper plot shows current passing through the superconducting MOT coils as measured with a sense resistor with switch off times of $< 100 \mu\text{s}$; the lower plot shows the corresponding magnitude of the magnetic field response measured with a 3-axis magnetic field sensor outside the chamber. (b) The trajectory of the 3D MOT as the magnetic fields relax.

coil axis to break the free path for Eddy currents to form. The cut was made on this axis as it corresponds to the largest field gradient generated by the coil pair. This does not account for the magnetic field on the radial axis, which has half the field gradient of the axial axis and is likely the source of Eddy currents in our system. The time evolution of the magnetic fields are shown in Fig. 7.3a, for the MOT coil switch-on and switch-off. For this measurement, a three-axis magnetic field sensor (Bartington Instruments: MAG-13MS100) was placed outside one of the horizontal MOT beam view ports. Measurement of the current with a sense resistor indicate that the MOT coils switch in $< 100 \mu\text{s}$, however the magnetic field relaxes with a 90:10 time of $\sim 0.5 \text{ s}$. It is likely that additional cuts in the radial axis of the mount would have reduced the

strengthen of Eddy currents in the system and hence further lowered the relaxation time of the magnetic fields.

For a future upgrade to the system, an additional consideration is the choice of material for the cryogenic mount. Particular notice should be given to the electrical conductivity which is the dominant property contributing to the strength of Eddy currents. However, the material should also have good thermal conductivity to take full advantage of the cooling power provided by a cryostat. Oxygen-free high-conductivity (OFHC) copper was chosen because it is a widely used material in cryogenics due to its superb thermal conductivity particularly below 30 K [189], however, it also has high electrical conductivity which increases by over an order of magnitude as $T \rightarrow 20$ K compared to room temperature [205]. Other common materials used in cryogenics suffer from similar issues, aluminium for example has slightly worse³ thermal conductivity compared to copper, but higher electrical conductivity below 78 K. Another potential material is silicon which has similar thermal conductivity to copper and has effectively no electrical conductivity [206]. The poor ductility of silicon would make it a challenge to machine. Copper beryllium (Cu-Be) is an alloy which has worse thermal conductivity than copper but reduced electrical conductivity and enhanced ductility [207]. Demonstration of Cu-Be in a cryogenic MOT has been shown to reduce the effect of stray fields during coil switching, but did not completely eliminate the Eddy currents [9, 121]. Using Cu-Be with additional cuts in the mount could reduce induction of Eddy currents to a manageable level. Alumina (Al_2O_3) ceramics are electrical insulators with both high thermal conductivity and good mechanical properties for operation down to 4 K [208]. However, when compared to OFHC copper, Alumina possesses lower thermal conductivity which will result higher base temperatures for a given cooling rate.

One other solution would be to move the quadrupole coils away from the cryogenic stage by making a MOT elsewhere and moving the cold atoms to the trapping region via a transport beam [209]. This is already a common method employed by quantum gas microscopes [210], however this results in low experimental repetition rates. Work has

³This is true for most grades of Al except pure Al which has a thermal conductivity an order of magnitude above that OFHC copper's [189].

already demonstrated lattice transport for tweezer arrays [193, 194] in which tweezers can be loaded continually from a reservoir [195]. Implementation in a cryogenic system would prove an engineering challenge given the reduced optical access.

For any solution, it is imperative that tests of the magnetic field response to coil switching take place before commissioning the system. For this experiment, major changes to the mount would require an invasive rebuild followed by a second bake which would take considerable time. It was decided any changes to the design would be implemented in a future planned upgrade to the system as the initial goals of the experiment can be achieved despite the shifting magnetic fields.

7.4 Polarisation Gradient Cooling

To achieve a large number of sites for a fixed laser power, it is desirable in tweezer array experiments to load into shallow traps of ≤ 1 mK. This necessitates further cooling below the Doppler limit of the MOT to ensure atoms are suitably cold. To achieve this we employ the $\sigma^+ - \sigma^-$ PGC between initial atom cooling in the MOT and the loading of tweezers. PGC relies on velocity selective Raman transitions between magnetic sublevels creating a population imbalance and hence a cooling force due to the interaction with polarisation gradient [170]. A magnetic field will break the Zeeman degeneracy of the magnetic sublevels and shift the Raman resonances leading to a skewed population imbalance across the sublevels [171]. This new imbalance will act to heat the atoms rather than cool. Therefore, it is imperative to ensure that $\sigma^+ - \sigma^-$ PGC takes place in a zero magnetic field regime.

It is very unlikely that the superconducting coils' quadrupole field zero point will be in the same location as the PGC zero field point. To ensure these locations coincide between the lenses, it is typical to perform the MOT loading and PGC stage at different shim field values. The relatively long relaxation time of the magnetic field after switching off the MOT coil discussed in Sect. 7.3 means that the PGC may require a time dependent shim field to ensure the atoms are continually cooled. To simplify the initial setup of the PGC stage, the assumption is made that the magnetic field is

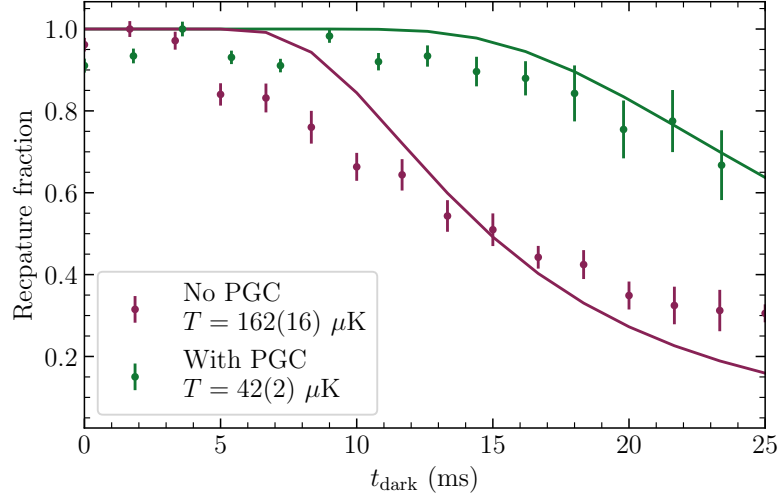


Figure 7.4: Release-recapture temperature measurements of the Cs cloud before and after PGC. The data are fit to Eq. 7.2.

roughly constant for the duration of the PGC. A shim control system which tracks the evolving magnetic field to ensure the entirety of PGC takes place at zero field will be implemented at a later date.

A release-recapture method is used to optimise the shim field values for PGC. After the MOT load stage is complete, the quadrupole coils are switched off in the dark. The PGC light is then applied for 5 ms after which the cloud is left to freely expand in the dark before being recaptured by the MOT and imaged. Cooling is optimised at the maximal recapture fraction f_r . The temperature of the atoms can be estimated by allowing the cloud to expand in the dark for time t_{dark} after the PGC is applied. The expected fractional recapture can be estimated with [211]

$$f_r = -\frac{2 \exp[-v_c^2/v_T^2] v_c}{\sqrt{\pi} v_T} + \text{Erf} \left[\frac{v_c}{v_T} \right], \quad (7.2)$$

where $v_T = \sqrt{2k_B T/m}$ is the thermal velocity of the atoms in the MOT at temperature T and $v_c = R_c/t_{\text{dark}}$ is the velocity at which the atoms reach the edge of the capture region R_c . Fig. 7.4 shows the result of the release-recapture measurements of the cloud with and without PGC, and fitted to Eq. 7.2. The atom number was determined from

the area under the Gaussian profile of the MOT fluorescence. The large error bars were due to poor signal-to-noise on the imaging axis as the recapture fraction reduced⁴. The MOT temperature was estimated to be $162(16) \mu\text{K}$; after PGC at a detuning of -4.8Γ with an intensity of $\sim 1I_{\text{sat}}$, the temperature of the cloud was estimated to be $42(2) \mu\text{K}$. The accuracy of the extracted temperatures were limited by several factors. One, Eq. 7.2 assumes a spherical cloud of atoms with a Gaussian velocity distribution which expands homogeneously; the MOT in this experiment is in fact a pill shape so will not expand as such. Secondly, the long relaxation time of the magnetic fields will affect the recapture of atoms and will indicate a higher temperature using the release-recapture method. A more accurate atom temperature could be obtained by using a time-of-flight [212] measurement, however, this was not possible due to the limited optical access provided by the cryostat. Atom temperature did not affect our ability to load traps, so further investigation was deemed unnecessary.

⁴This fitting error was hampered by the drifting magnetic field during coil switching as discussed in Sect. 7.3.

Chapter 8

Caesium Tweezer Arrays

One of the central aims of this experiment is to demonstrate enhanced trap lifetimes by utilising cryogenics in the vacuum system. This chapter will discuss the optical layout for generating tweezers arrays to trap single caesium atoms in order to characterise the vacuum limited lifetime. Initial characterisation of the traps is presented as a first step to obtaining metrics for testing the performance of the vacuum system. The final experimental results presented are a series of trap lifetime measurements which demonstrate improved survival after reducing the vacuum pressure through cryopumping and the use of a titanium sublimation pump (TSP). We finally discuss how further enhancements in trap lifetimes can be achieved by implementing additional in trap cooling for offsetting the finite heating rate of the traps. We end by discussing how design alterations would lead to further enhancements in trap lifetimes.

8.1 940 nm SLM and Imaging Optics

An spatial light modulator (SLM) is used to generate the 940 nm tweezer arrays for trapping Cs atoms [89]. This wavelength was chosen for trapping Cs as the equivalent trap depths for Rb are four times smaller effectively rendering anti-trapped at 940 nm. The SLM used for generating the arrays is a HoloEye PLUTO-2.1-NIR-118 with a sensor size of 1920 by 1080 pixels (15.36 mm by 8.64 mm) with a Raspberry Pi used as a remote display server for the holograms.

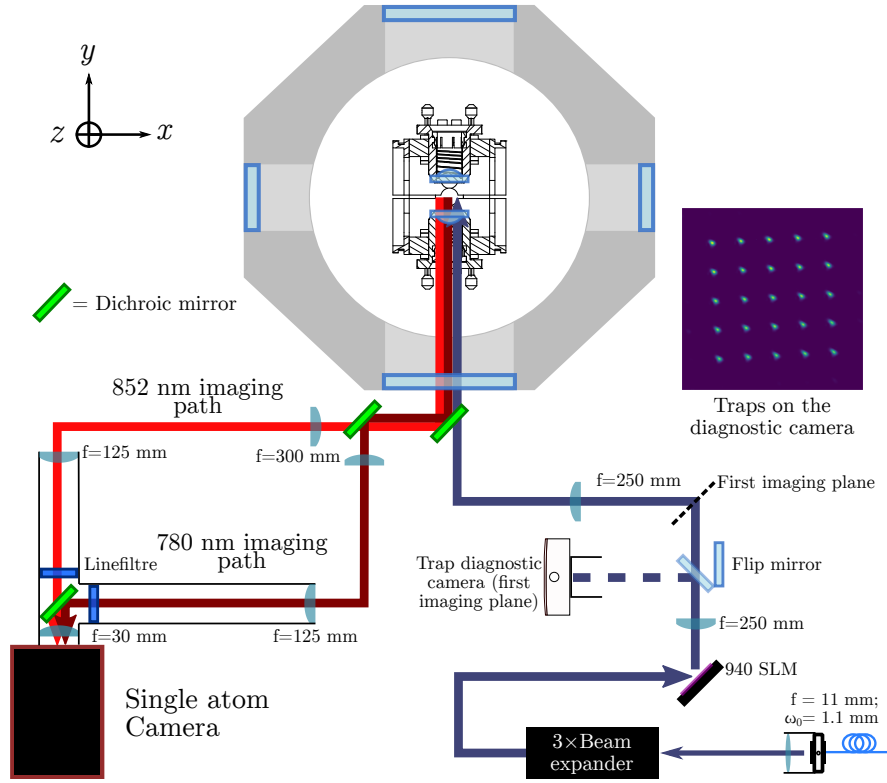


Figure 8.1: Schematic of the 940 nm trap beam path (blue arrow), and imaging path for 852 nm (light red) and 780 nm (dark red). The image shows the traps on the diagnostic camera in the first imaging plane before the chamber. The second in-vacuum lens allows for diagnostics of trap light after it passes through the chamber.

A schematic for the Cs trapping and bi-chromatic imaging layout is shown Fig. 8.1 where the 940 nm path is indicated by the blue beam trace. The 940 nm light is sourced from the laser breadboard — discussed in Sect. 5.5.1 — through a high power fibre. The beam is initially collimated with an 11 mm lens to a $1/e^2$ radius of 1.1 mm before being expanded by a 3 times beam expander to ensure the beam fills the SLM’s sensor. The spot array hologram is generated with the Gerchberg-Saxton (GS) algorithm and applied on top of a blazed diffraction grating. A lens phase mask is also added to the hologram to allow adjustments of the trap focal plane inside the chamber. The zeroth order diffraction off the SLM is blocked, and first order is used for the trap array. After the SLM, an $f = 250$ mm lens is used to create the first imaging plane outside the chamber. A flip mirror situated between the two lenses diverts the light into a camera (Thorlabs: CS165MU1/M) with its sensor located at the first imaging plane.

The camera is used to provide trap diagnostics and for feedback during calibration of the SLM. At the focus of the first imaging plane, the traps have a waist of $\sim 25 \mu\text{m}$. Zeemax modelling found that a second $f = 250 \text{ mm}$ lens after the first Fourier plane would create a trap radius $1.15 \mu\text{m}$ inside the chamber. After the second lens, the trap beam is combined with the imaging path for 780 nm and 852 nm using a series of dichroics situated next to the window of the vacuum chamber. The laser beam finally enters the chamber and the in-vacuum lens projects the hologram onto the second imaging plane for single atom trapping in the centre of the cryogenic mount.

Atom imaging of Rb and Cs is performed with a single camera. The bi-chromatic imaging paths are shown in Fig. 8.1 where the path for 852 nm is indicated by the light red arrow and the path for the 780 nm is indicated by the dark red arrow. Atom fluorescence is collected by the in-vacuum aspheric and directed out the chamber counter-propagating to the trap laser beam. The 780 nm and 852 nm light are separated from the trap light by a dichroic mirror; a second dichroic mirror then splits the 780 nm and 852 nm fluorescence along different paths. The two beams are directed towards different telescopes allowing independent control of the focus and magnification at each wavelength. The 852 nm beam requires 50 mm extra path length to ensure 1 pixel on the camera is $\sim 1 \mu\text{m}$ in the imaging plane for both species. Two optical bandpass filters at 780 nm and 852 nm ensure each arm of the imaging path only transmits its corresponding wavelength. The beams are then recombined with a second dichroic and focused onto the camera with an achromatic lens. This layout provides a magnification of 7.23 and an field of view (FOV)¹ of $126 \mu\text{m}$ for 852 nm imaging, and a magnification of 7.27 and an FOV of $128 \mu\text{m}$ for 780 nm imaging.

The camera used for single atom imaging is a Kinetix22 which has a quantum efficiency of $> 55\%$ at 852 nm and $> 70\%$ at 780 nm. Accounting for the efficiencies of the in-vacuum lens ($\text{NA} = 0.45 \implies 5.35\%$ collection efficiency), transition through the imaging optical path (77%) and the quantum efficiency of the camera; the photon collection efficiency is 2.3% for Cs assuming isotropic emission of the atoms.

¹Commonly defined as the edge-to-edge extent of the imaging plane.

8.2 Single Atom Trapping

8.2.1 Light Assisted Collisions

The number of atoms loaded into the traps is governed by Poissonian statistics which indicate that only $\sim 50\%$ of traps will load with single atoms while the rest will either have no atoms or ≥ 2 atoms. The ejection of atoms from traps with more than a single load is achieved through a stage of light assisted collisions (LAC) [213–215]. Consider a doubly occupied trap. The ground state $|\mathbf{S} + \mathbf{S}\rangle$ energy interaction is of a van der Waals (vdW) type which implies the atoms will not interact except at very short distances. The first excited state $|\mathbf{S} + \mathbf{P}\rangle$ has energy interactions of the resonant dipole-dipole (dd) type which greatly increases the range of interaction. At suitable inter-atom separation and in the presence of a red detuned laser beam, the pair state will be excited to the lower branch of the $|\mathbf{S} + \mathbf{P}\rangle$ state which is an attractive potential and provides enough kinetic energy for both the atoms to be ejected from the trap. Blue detuned LAC[216] acts to ensure only a single atom is ejected from the trap and has demonstrated loading rates of $> 90\%$ [77]. This requires more complex control of detunings and interaction timescales of ~ 0.5 s and as such we chose not to implement this into our system. Sub-Poissonian loading can also be shown with Λ -enhanced gray-molasses [217] which can achieve loading rates of 90% [93].

8.2.2 Trap Loading Procedure

The procedure for loading single Cs atoms in traps is shown in Fig. 8.2. As 940 nm is near magic for Cs, the tweezer array light is turned on for the entirety of the sequence in continuous wave (CW) operation with a trap depth of 3 mK. If instead a non-magic wavelength was selected, the traps would have to be chopped out of phase with any cooling light due to the differential stark shifts between the ground and excited states — as will be the case for the 810 nm laser for trapping Rb. Due to the long relaxation times of the magnetic field when the quadrupole coil is turned off, the MOT is loaded for a duration of 2 s. This time can be reduced once the field tracking system is implemented. After the 5 ms PGC stage, all the MOT cooling and repump light is

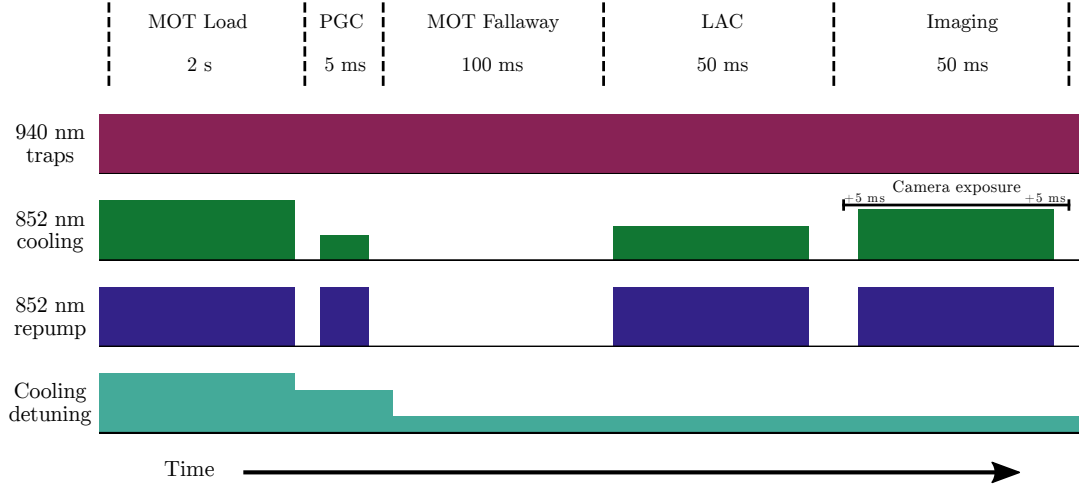


Figure 8.2: Loading sequence for single atom traps. The 940 nm trapping light remains on for the entire sequence without need for out-of-phase chopping due to the magic nature of the D_2 transition at this wavelength.

turned off and the Cs atoms are held in the tweezers for 100 ms which allows the cloud to disperse. The LAC is performed at a detuning of -4Γ and $0.7I_{\text{sat}}$ for 50 ms. The image of the trap atoms is taken following the LAC stage to obtain trap occupancy data.

8.3 Single atom imaging

Imaging of single Cs atoms is performed by illuminating the array with cooling and repump light on the horizontal MOT axis, and collecting the fluorescence signal with the in-vacuum aspheric lens and directing the signal towards the single atom camera as described in Sect. 8.1. The camera is exposed for 50 ms, but the atoms are illuminated for 40 ms allowing 5 ms of padding for initialisation of the camera. During imaging, the vertical MOT beam is blocked with a shutter to reduce scatter off the lenses. The imaging is performed at a detuning of -6Γ linewidths and a power of $17I_{\text{sat}}$ for a 5 by 5 array with $15\ \mu\text{m}$ spacing. Fig. 8.3a shows an average single atom fluorescence signal from 3000 camera images of the array. The trap beam appears to pass through a lens off centre resulting in the apparent coma present in the image. Realignment of the trap beam would correct this, but was not necessary to perform initial characterisation of

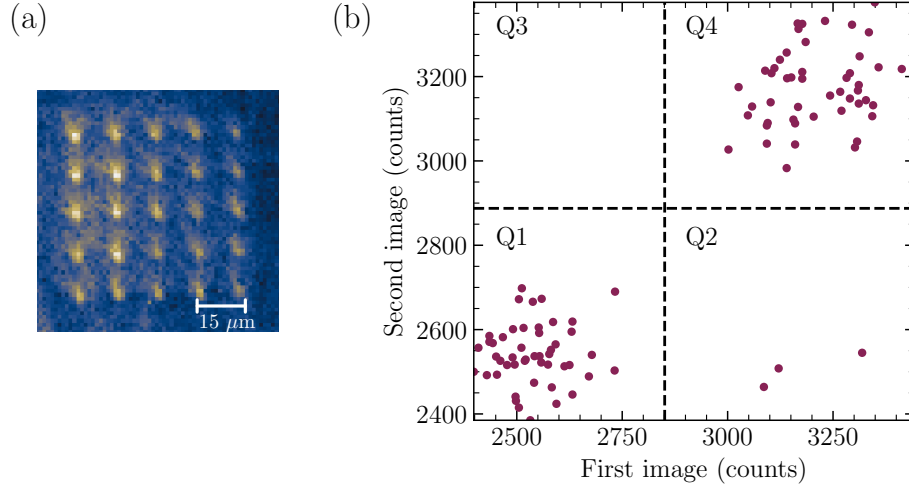


Figure 8.3: (a) Averaged image from 3000 shots of a 25 by 25 array of single Cs atoms. (b) Four quadrant plot comparing two images taken in sequence to inspect atom survival and loss. This data was recorded after optimisation imaging.

the traps. To optimise the traps, we use an region of interest (ROI) engine which uses all image shots to identify each trap position. The engine then uses the 10 brightest surrounding pixels and produces a number of counts associated with each trap on an image-by-image basis.

For initial single atom imaging optimisation, two fluorescence images separated by 100 ms are taken. Fig. 8.3b compares the ROI counts between the first and second image for a single trap across 100 experiment runs. Evidence of single atom loading is apparent by the emergence of two clusters; counts in Q1 imply empty traps in both images whereas counts in Q4 indicates an atom appears in both images. Counts in Q2 imply an atom was present in the first image but was lost by the second image.² Trap optimisation is performed by maximising the separation between the clusters in Q1 and Q4 and minimising counts appearing in Q2.

Once good separation of the clusters is achieved, single atom statistics can be investigated through the bimodal nature of the distribution. Fig. 8.4a shows a histogram of counts c for a single trap across 3000 images taken after loading the array. The bimodal distribution will follow a Poissonian statistical model, however, for a large mean c , we

²Q3 implies trap hopping.

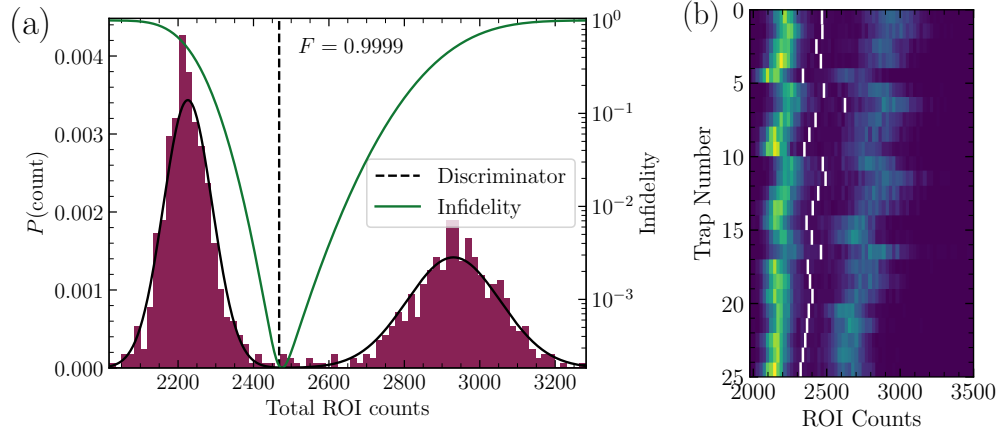


Figure 8.4: (a) Single site ROI camera counts histogram with two peaks indicating empty trap and trap loaded with a single Cs atom. (b) Histogram across all the arrays with discriminator indicated by the white line.

can instead model the imaging as a double Gaussian,

$$P(c) = \sum_{i=\{0,1\}} \frac{p_i}{\sqrt{2\pi\sigma_i^2}} \exp \left[-\frac{(c - \mu_i)^2}{2\sigma_i^2} \right], \quad (8.1)$$

where p_0 is the probability of an empty trap, p_1 is the probability of loading a trap, and σ_i and μ_i are the corresponding Gaussian width and mean respectively. Eq. 8.1 is fit to the histogram. The infidelity defines the average false positive and false negative counts,

$$(1 - F) = \frac{1}{2} \left(1 - \operatorname{erf} \left[\frac{c - \mu_0}{\sqrt{2}\sigma_0} \right] \right) + \frac{1}{2} \left(1 + \operatorname{erf} \left[\frac{c - \mu_1}{\sqrt{2}\sigma_1} \right] \right). \quad (8.2)$$

The minimum of Eq. 8.2 is the discriminator which is used to determine if a trap is loaded or not. A discriminator can be defined for each trap individually as shown in Fig. 8.4b for a 25 atom array, and can be provided to the ROI engine to allow on-the-fly trap occupancy information.

Most experiments will require multiple rounds of single atom imaging. Atom survival between images must be maximised to ensure results are not influenced by imaging losses. Across a 25 atom array, this experiment shows an image-to-image atom survival of $\sim 80\%$. The cause of this was found to be the time dependent magnetic field dis-

cussed in Sect. 7.3, which suggests imaging takes place at different fields which results in atom loss during imaging. Having each image take place at a different shim field should offset this loss. This was not done for the work presented in this thesis as it did not limit our ability to measure the trap lifetime. This will be implemented in future iterations of the experiment.

The imaging protocol discussed above will be suitable for characterisation of atoms in traps, however, it will be unsuitable for state selective readout. This is due to the repump ensuring atoms in the $F = 3$ hyperfine manifold are returned to the imaging transition. Once a system for optical pumping is implemented, clock state qubits will be readout via a blow away method in which atoms in the $F = 4$ state are ejected from the traps [218]. Blow-away readout has a couple of drawbacks. Firstly it cannot distinguish between qubits in the $F = 4$ state or whether an atom has been lost during the experimental cycle. Secondly, for multi-round error correction at scale, retention of atoms after readout is preferred, but blow-away requires the loss of some atoms. One technique for non-destructive readout (NDRO) is to use high purity circularly polarised light in the presence of a bias field which ensures atoms in the $F = 4$ manifold are imaged on a closed transition with no decay path to the $F = 3$ manifold eliminating the need for repumping [23]. This method cannot distinguish between a qubit in the $F = 3$ manifold or atom loss. An alternative method for NDRO is implementation of a spin-dependent transport lattice which moves atoms in only one of the hyperfine ground states to a specially separated tweezer where imaging is performed on atoms in both manifolds [7, 118]. This method of imaging is both non-destructive and robust for identifying atoms loss.

8.4 Trap Lifetimes

8.4.1 Cryopumping

The trap lifetimes of neutral atoms in tweezer arrays is far shorter than other atomic qubit platforms such as ions, which can achieve lifetimes of months due to the strong Coulomb force underpinning the principals of Penning traps [47–50]. The lifetimes

of neutral atoms in optical dipole traps (ODTs) in room temperature experiments is typically of order 10 – 20 s due to the weakly interacting nature of the atomic ground state. The probability of holding an atom in a trap decays exponentially with time when the system is vacuum limited and is defined by $1/e$ lifetime τ . Importantly τ scales as $1/N$ for atom number N . This means for a single atom lifetime of 10 s, scaling to 10000 atoms will result in a lifetime of 1 ms, which is far shorter than the time requirements for most experiments. If instead, the single atom lifetime can be extended to 1000 s, the array lifetime for 10000 atoms will be 100 ms, which greatly increases the time budget for experimental runs. This issue becomes more pressing when considering error correction at scale which is expected to require the scaling to systems with hundreds of thousands or millions of qubits [219, 220].

One of the main contributions to loss in ODTs is collisions with background gas in the system and it is therefore desirable to reduce the pressure of the vacuum system to extend lifetimes. The use of cryogenics has been shown to reach pressures below those typically achieved with ion pumps alone with the lowest pressure recorded in a lab environment utilising a cryostat³ [221]. Cryogenic surfaces can reduce the pressure in a vacuum system through the process of cryopumping where background gas, which would usually rebound off walls and become *frozen-out* once the temperature of cryogenic surfaces fall below the sublimation point. Fig. 8.5 shows the sublimation curves for common background gas found in-vacuum systems [183]. In a baked vacuum system, a large contribution of the background gas has sublimation temperatures between 20 K and 30 K. As our system has a base temperature of 7 K, the only impurities that cannot be cryopumped are H₂ and He. H₂ can be cryopumped in UHV with specialised pulsed tube cryocoolers which can reach temperatures of 1 K [222], however, neither H₂ nor He are heavy enough to transfer enough kinetic energy to eject an atom from a trap and instead collisions with these elements act to heat the atom [223]. Implementation of active in trap cooling can be used to counteract this heating effect.

The effectiveness of cryopumping reduces as cryogenic surfaces become saturated, as such impurities in the vacuum system no longer stick to cryogenic surfaces and

³Pressures $< 10^{-14}$ mbar [221].

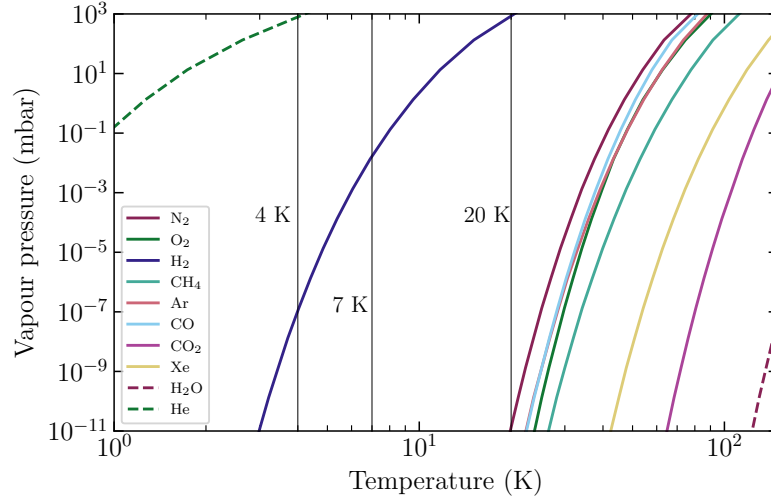


Figure 8.5: Sublimation curves of common background gases found in-vacuum systems [183] with the key temperatures of liquid He (4 K), the system’s base temperature (~ 7 K) and the sublimation point of N_2 (~ 20 K at 10^{-11} mbar) indicated by the vertical black lines.

instead bounce off. This results in a rise in background pressure and hence decreases the vacuum limited lifetime. Regeneration of cryogenic surfaces can be performed by raising the temperature of the surfaces to above the sublimation point of the impurities, allowing the ion pump to bring down the pressure and reviving the enhanced trap lifetime once cryopumping resumes [123]. Fig. 8.5 suggests raising the temperature of the cryostat to between 20 K and 30 K will suffice for many common vacuum impurities. Raising the temperature above the sublimation points of Xenon and CO_2 may further improve the regeneration of cryogenic surfaces however, the former should have minimal contribution to the background gas, so will likely limit the improvement of vacuum lifetime.

Recent development in-vacuum systems for neutral atom tweezer arrays has demonstrated vastly improved vacuum using a TSP. When a current of 48 A is passed through a TSP, the titanium sublimates and coats the walls acting as a passive pump. When carefully designing a vacuum system to maximise surface area coated with titanium, > 1000 s trap lifetimes have been achieved without the need for cryogenics [8]. The implementation of the TSP in our system was to improve lifetimes by reducing the

starting pressure before cryopumping. Additionally, the TSP may slow the saturation of cryopumping if it is fired while the cryostat's temperature is set above the sublimation point of contaminants.

Titanium is an excellent getter for pumping reactive elements and compounds such as non-saturated hydrocarbons, however, it is a poor at removing non-reactive elements and compounds such as methane [224, 225]. A TSP is notably effective at pumping H_2 which cannot be cryopumped in a 4 K UHV environment [226]. In this system, the TSP should lower the base pressure, hence reducing the saturation rate of cryogenic surfaces and ensuring the cryostat is used to pump the non-getterable impurities.

8.4.2 Measurement of Trap Lifetimes

The trap lifetime τ was obtained by recording the survival between two images separated at by time t . By comparing occupancy between the two images and scanning the value of t , the survival probability can be extracted from the function

$$P(t) = \exp\left[-\frac{t}{\tau}\right] \cdot \eta(t), \quad (8.3)$$

where $\eta(t)$ characterises the atom survival from heating due to in-trap scattering [169]. The coefficient $\eta(t)$ depends on the trap depth U_0 and the initial temperature of the atoms T_0

$$\eta(t) = 1 - \exp\left[-\frac{U_0}{k_B(T_0 + \dot{T}t)}\right] \left(1 + \frac{U_0}{k_B(T_0 + \dot{T}t)} + \frac{U_0^2}{k_B^2(T_0 + \dot{T}t)^2}\right), \quad (8.4)$$

where $\dot{T}$ is the heating rate and k_B is the Boltzmann constant. When loss is dominated by background collisions, $\eta(t)$ has negligible effect on τ and the survival will be purely exponential. Conversely, when trap heating dominates loss, the survival will no longer evolve with only an exponential trend. For measuring lifetimes in this thesis, the value for T_0 was taken to be the PGC temperature obtained in Sect. 7.4. For a more accurate value of T_0 , a release recapture measurement [227] of the atoms in traps should be used for atom temperature, however, at the time of writing this had not been performed.

Though the PGC temperature will not be as accurate as the release recapture method, variations in T_0 did not have an outsized impact on obtaining a value for τ which is the important coefficient when characterising the vacuum limited lifetime. The lower limit of $\dot{T}$ is defined as

$$\dot{T} = T_{\text{rec}}^{(D\chi)} \frac{\Gamma}{3\hbar\delta} U_0, \quad (8.5)$$

where $T_{\text{rec}}^{(D\chi)}$ is the recoil temperature correspond to the D_χ -lines where $\chi \in \{1, 2\}$. For Cs, given $T_{\text{rec}}^{(D1)} = 180$ nK and $T_{\text{rec}}^{(D2)} = 198$ nK the combined heating rate of these two transitions is $1 \mu\text{K s}^{-1}$ for a trap depth of 3 mK at 940 nm.

The first measurement of a trap lifetime was performed 2.5 months after the previous thermal cycling of the cryostat to room temperature. At room temperature, the TSP had been fired as described in Sect. 6.5; following the cooldown, the pressure on the science chamber ion-pump was $< 10^{-11}$ mbar. In the time leading up to this lifetime measurement, the pressure on the ion pump had begun to show readings above $1 \cdot 10^{-11}$ mbar, indicating the cryogenic surfaces had become saturated and the layer of titanium had become somewhat passivated. The expectation would be that due to this rise in pressure, the lifetime would be vacuum limited and likely comparable to a room temperature experiment.

Fig. 8.6 shows the survival obtained for this first lifetime measurement where $\tau = 13.2 \pm 0.5$ was measured. This is not much of an improvement over the room temperature experiment which achieved a lifetime of ~ 12 s [26]. It is likely that the lifetime is vacuum limited based of the exponential nature of survival curve implying that measures to reduce the vacuum should yield improvements.

To confirm that the trap lifetime was vacuum limited, the cryostat was warmed up for a brief period following methods from other experiments [12]. This would sublime impurities off cryogenic surfaces that are then removed from the background by the ion pump, with cryopumping reactivating once the system returns to 7 K. The process of warming up the cryostat involves slowly stepping up the temperature of the cryostat using the heater on the cold finger until the lower sensor indicates 25 K. This was chosen as a target temperature as it would encompass much of the sublimation temperatures

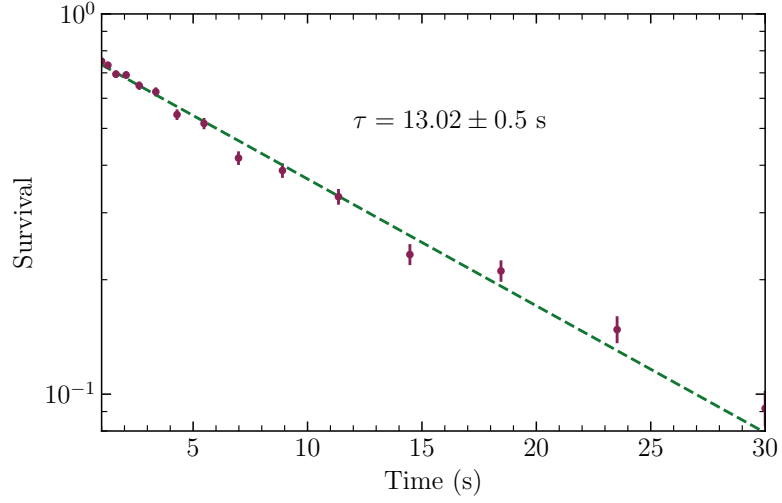


Figure 8.6: Lifetime measurement results taken when the cryogenic surfaces had saturated.

indicated in Fig. 8.5. Temperature and pressure in the science chamber are shown in Fig. 8.7. The heater was stepped slowly over the course of 3 hours to ensure any spike in pressure did not choke the ion pump. The max heater power applied was 11.25 W which had an equilibrium temperature of 25.0 K for the lower temperature sensor and 22.9 K for upper temperature sensor. The peak pressure in the science chamber was at $1.3 \cdot 10^{-8}$ mbar during this procedure. The system was left at 25 K for 25 hours. Within two hours reading on the ion-pump fell below $1 \cdot 10^{-11}$ mbar. The system returned to its base temperature within 30 minutes of switching off the heater.

The speed of heating up the cryostat was limit by the pressure reading of the ion-pump which we did not want to spike to aggressively. This was due to the level of saturation achieved since the last warm-up to room temperature. Performing more regular warm-ups will reduce the impurities being released into the vacuum system speeding up the process⁴. This was confirmed by a second regeneration of the surfaces performed a couple of days after the one described above which took 30 minutes to reach 25 K. In future, regular periodic cycling of the cryostat to 25 K will be implemented when the system is in a fully operational state.

⁴Most He is used during the cooling of the cryostat; therefore, more regular cycling will increase the rate at which the He cylinder must be replaced.

Following the cryogenic surface regeneration, a second lifetime measurement was taken shown by the red markers and green fit line in Fig. 8.8 which indicate an increased lifetime of $\tau = 33.02 \pm 3.0$ s. This result confirms that the cryopumping can be reactivated by raising the temperature of the cryostat to clean the cryogenic surfaces. The non-linear nature of the log plot in Fig. 8.8 arises from the $\eta(t)$ term in Eq. 8.3 which starts contributing to the trend and implies trap heating is starting to limit τ .

To obtain further improvements in-vacuum pressure, we decided to utilise the TSP. We chose to fire the TSP while the system was at cryogenic temperatures as there would be no line of sight between the hot element and any cryogenic surfaces. The system was warmed up with the heater to 25 K to regenerate the cryogenic surfaces. The peak pressure recorded in the science chamber was $1.8 \cdot 10^{-9}$ mbar. The ion pump was allowed to pump down the impurities for one hour before it was switched off for the TSP firing. The TSP was fired 4 times for 60 s, with 10 minutes between each firing. After 10 minutes from the last firing, the ion pump was switch on and the pressure on the ion pump spiked before falling below $1 \cdot 10^{-11}$ mbar. The system was left hot for a further 2 hours (5 hours in total) before the heater was turned off. A final lifetime measurement was then taken following this process with the results indicated by the blue markers and cyan fit line in Fig. 8.8. The results indicate a lifetime measurement of 45.32 ± 3.9 s and demonstrate that both the combination of a cryostat and TSP can be used to improve vacuum lifetimes.

8.4.3 Towards Longer Lifetimes

The series of lifetime measurements demonstrated in the previous section have shown that both a cryostat and TSP can be used to increase the trap lifetimes. Despite this, our results do not reach the same timescale achieved in state-of-the-art experiments that can reach lifetimes in the 1000s of seconds [8, 9, 123]. However, it is important to note that our design does not completely shield the atoms from room temperature surfaces and was designed instead as a proof of concept for using a cryostat to enhance lifetimes.

In this section we discuss how the lifetime could be improved from the tens of seconds

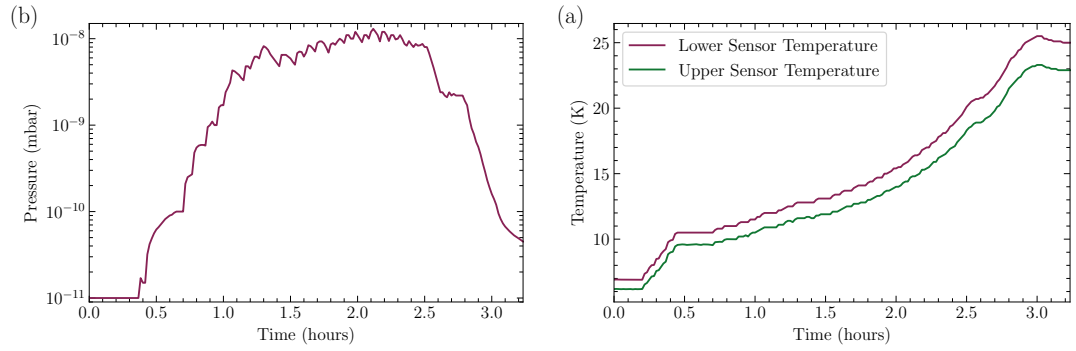


Figure 8.7: Warm up of the cryostat to 25 K. (a) Pressures of the science chamber ion pump as impurities are sublimated from cryogenic surfaces. (b) Temperature of the cryostat during the warm up.

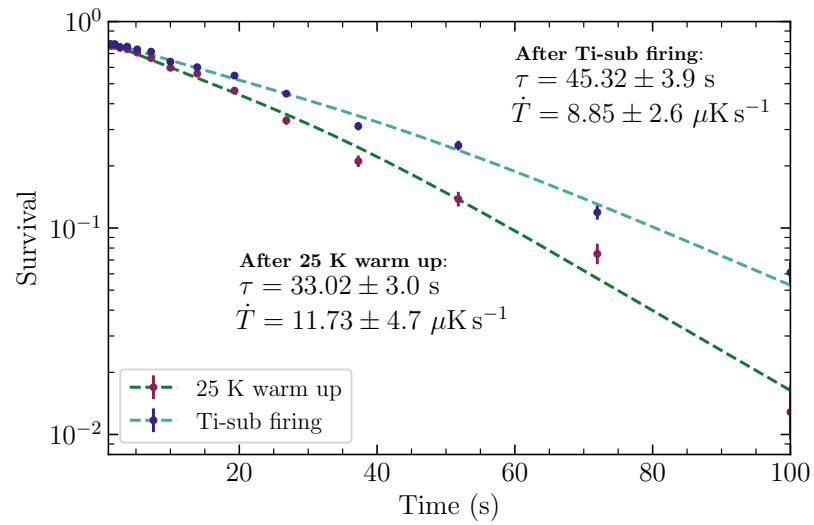


Figure 8.8: Lifetime measurements after regeneration of cryogenic surfaces by warming up to 25 K (red markers and green fit line) and lifetime measurement after firing the TSP (blue markers and cyan fit line). Each configuration improved lifetimes due to the reduction in the vacuum pressure.

range to hundreds of seconds timescale. Following this, we then discuss improvement through design and engineering.

Further Surface Regeneration

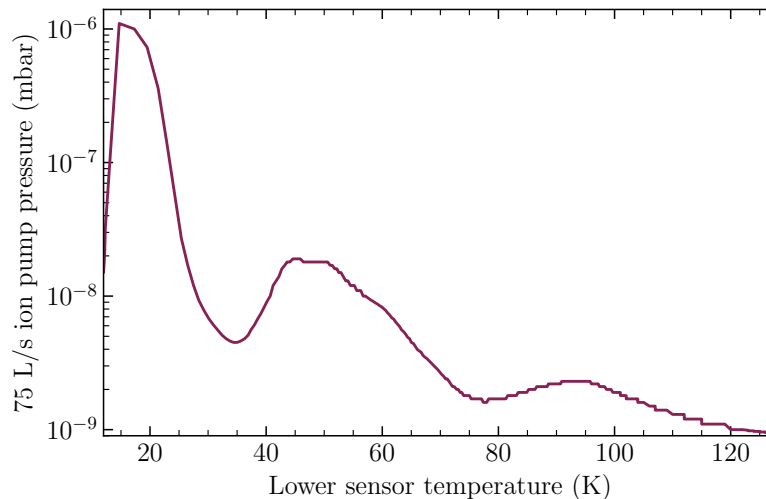


Figure 8.9: Science chamber 75 L/s ion pump pressure reading during the warm up of the cryostat. The pressure spikes correspond to spikes sublimation pressures of various contaminants.

The surface regeneration presented in Sect. 8.4.2 only involved raising the temperature of the cryostat to 25 K to encompass much of the sublimates expected to be frozen out according to Fig. 8.5. There are additional sublimation pressures for Xenon and CO₂ between 25 K and 100 K, however, the abundance of these gases should be low in the vacuum system. Fig. 8.9 shows the pressure of the vacuum system when warming up the cryostat up to ~ 130 K and indicates two additional peaks in pressure beginning at ~ 35 K and ~ 80 K. This would suggest that a more aggressive warm up towards 100 K may provide additional regeneration of the cryogenic surfaces, but improvements may be limited given the additional peaks are over an order of magnitude below that of the peak at ~ 20 K. However, as the warm-up is not linear in time⁵, the first ~ 20 K – ~ 30 K in Fig. 8.9 will represent a shorter time period for the ion-pump to bring down

⁵See Fig. 6.14 in Sect. 6.6.

the pressure. To understand the composition of these pressure spikes, a residual gas analyser could be used to indicate the make-up of the background gas, however one is not installed in this system so was not done here.

In Trap Heating

As the vacuum limited trap loss diminishes, heating within the trap becomes the dominant factor in the attainable lifetimes. This was apparent in the first cryogenic tweezer experiment with trap half-lives of 30 s [228], limited by the effects of the $\eta(t)$ term in Eq. 8.3. The application of continuous in-trap PGC extended $1/e$ lifetimes to 162 s and the application of pulsed PGC for 15 ms every 10 s resulted in $1/e$ lifetimes 355 s.⁶ Similar improvements in lifetimes have been demonstrated in subsequent cryogenic tweezer array experiments [12, 123], and room temperature experiments which utilise TSPs for improved vacuum [8].

Given that PGC for single atoms reduces the atom temperature exponentially [171] and the ODT provides a constant heating rate, it would be expected that for sufficient PGC cooling rate, the atom temperature should equilibrate far below the trap depth. Hence continuous PGC should be sufficient to extend trap lifetimes to the vacuum limit. However, the continuous application of PGC in traps creates an effective molasses in the tweezer array. Some of the residual free Cs atoms may be slowed by the PGC light which effectively raises the local pressure around the array, hence, increasing atom loss through collisions. Additionally, atoms ejected from the trap due to heating from background collisions with H₂ have an increased chance of recapture by a neighbouring occupied trap. This second doubly occupied trap will experience an effective LAC from the PGC and imaging light causing additional trap loss.

A pulsed PGC scheme can ensure the temperature of atoms in traps remain well below the trap depths while avoiding the continuous cooling of residual atoms around the array. The typical cadence of pulsed PGC is of the order 10 s [9, 123] which ensures a balance between in trap cooling and allowing the diffusion of atoms ejected from traps.

⁶The pulsed PGC lifetime was later extended to 6050 s after fixing a vacuum leak [9].

8.4.4 Vacuum and Cryogenic Design

There are a few changes to the cryogenic system that could be implemented which will improve the trap lifetimes. As described in Sect. 6.3, a large portion of the solid angle around the trapping region is cryogenic, however there are room temperature surfaces with line of sight to the arrays. This is a source of background impurities which will decrease the attainable vacuum limited lifetimes. The experimental apparatus which achieved > 6000 s trap lifetimes eliminated nearly all line of sight to room temperature surfaces by implementing windows on the cryogenic stage for MOT and ODT beams [9]. There was one hole in the cryogenic stage which allowed atoms to load from an oven, which would act as a source of background impurities⁷. An upgrade to the system which implemented a microscope objective required its placement in-vacuum but not at cryogenic temperatures [122]. This meant the last lens facing the atoms was at room temperatures, which limited the trap lifetime to ~ 100 s. Another cryogenic design implemented an out-of-vacuum microscope objective with cryogenic windows allowing optical access for tweezer arrays [123]. This system had vacuum limited trap lifetimes of 3000 s with remaining background collision attributed to residual atoms from the MOT in the trapping plane. To remove residual background atoms from the system, the 3D MOT could be moved away from the trapping region, and atoms brought in via a lattice transport as recently demonstrated [195]. This would however increase the complexity of optical design to implement.

Improving the solid angle of cryogenic surfaces facing the array would reduce the effects of black-body radiation (BBR) causing coupling to nearby Rydberg states and increase the Rydberg lifetime by several factors [124]. Recent results have demonstrated improved p -state Rydberg lifetimes in a completely enclosed cryogenic system [124]. Longer Rydberg state lifetimes are desirable as enhanced fidelities are attainable due to suppression of errors associated with spontaneous emission.

Additional vacuum pumping provided by the TSP could be enhanced with a change in design to the surfaces coated in titanium. The current design, described in Sect. 6,

⁷The > 6000 s lifetime was achieved by closing a gate valve to the oven so only 300 K surfaces would be visible through the hole [121].

has the TSP away from the trapping region with a limited surfaces area on which to coat titanium. The system that achieved trap lifetimes of > 1000 s at room temperature did this with a vacuum system in which almost all metal surfaces where coated in titanium effectively making all the walls a passive pump. Such a setup could be implemented in a cryogenic system by coating a large volume of metal surfaces near the cryostat with titanium, however, care would have to be taken to avoid hot metal gas from hitting cryogenic surfaces.

Part IV

Outlook and Conclusion

Chapter 9

Outlook

The introduction to this thesis presented three challenges facing tweezer array experiments. To reiterate in brief, they were: (1) scalability; (2) destructive readout; (3) imperfect gate operations. The work presented in this thesis has described progress towards an experimental system which hopes to address these challenges. The issues of scalability was discussed in Part II and Part III in which a cryogenic vacuum system demonstrated enhanced trap lifetimes as a route towards larger arrays. Part I presented a means of attaining higher gate fidelities utilising Förster resonances in Rydberg d -orbitals. At the time of writing, realisation of dual-species arrays as a means to perform non-destructive readout has not been fully realised. This chapter will discuss the progress towards development of a dual-species system as well as work towards a Rydberg system for studying d -orbital Förster resonances.

9.1 Dual-species Arrays

At the time of writing, only the Cs portion of the experiment had advanced towards single atom trapping and as such the results in Part III only discuss this atomic species. However, the Rb magneto-optical trap (MOT) and optical dipole trap (ODT) systems have been developed in parallel with work presented in this thesis, assembled by my fellow PhD student Ian Moodie. Finishing of the Rb system will mark the completion of one of this experiment's aims of operating a dual-species tweezer array with a cryogenic

apparatus. In this section we will discuss progress towards realising dual-species arrays.

The 852 nm cooling laser and tapered amplifier (TA) were repurposed from a previous experiment so the Cs system was developed first. The 780 nm setup on the other hand required a completely new laser system which delayed its progress. Further delays to the development of the Rb MOT were caused by issues relating to the optical access afforded by the cryostat, described in Chapter 7. Additionally, there were issues with the 780 nm TA and laser for both the cooling and repump system which required replacements of both. Despite these challenges, several key milestones have been achieved on the Rb system. A 2D MOT of Rb, shown in Fig. 9.1a, was loaded after about a month from when we exposed the high pressure side of the experiment to Rb. The Rb 2D MOT was achieved without noticeably degrading the loading of a Cs 3D MOT or the loading of Cs traps. Due to the limited 780 nm laser power, simultaneous operation of the of the Rb 2D and 3D MOT was not initially possible. However, after replacement of the 780 nm cooling TA, we have managed to achieve loading from the 2D MOT to the 3D MOT — as shown in Fig. 9.1.

Simultaneous loading of the Rb and Cs 3D MOTs has also been demonstrated as shown by the image in Fig. 9.2. The two MOTs formed spatially separated due to standing waves from the horizontal 3D MOT beams. This results in different intensity node locations at 780 nm and 852 nm and hence the position offset. The two MOTs were overlapped by making a small adjustment to the horizontal retro-mirror. Current operations of the system mean each MOT will be optimised independently for loading into arrays. Once each species can be loaded into their respective arrays, we shall move the system to operate both MOTs and load arrays simultaneously.

The 810 nm Rb tweezer array laser system was built in tandem with the Cs laser system. Both the laser and TA follow the same design as the Cs system (see Sect. 5.5). A separate spatial light modulator (SLM) is used for generating the Rb traps but the optics layout is identical to the Cs one described in Sect. 8.1. The Rb and Cs tweezer paths are combined using a dichroic mirror immediately before the chamber. Using independent paths for each array makes setup and alignment far simpler and allows for for the generation of different geometries for each atomic species.

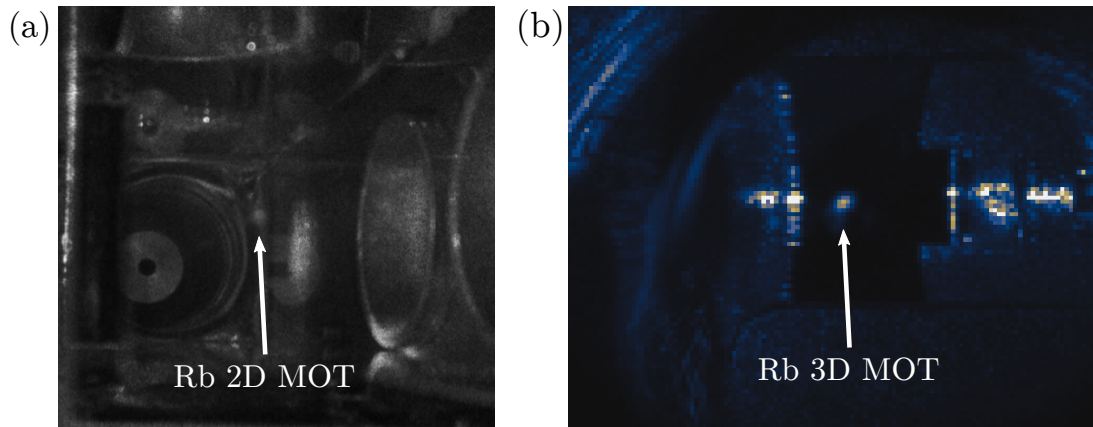


Figure 9.1: Rb (a) A Rb 2D MOT. (b) A Rb 3D MOT.

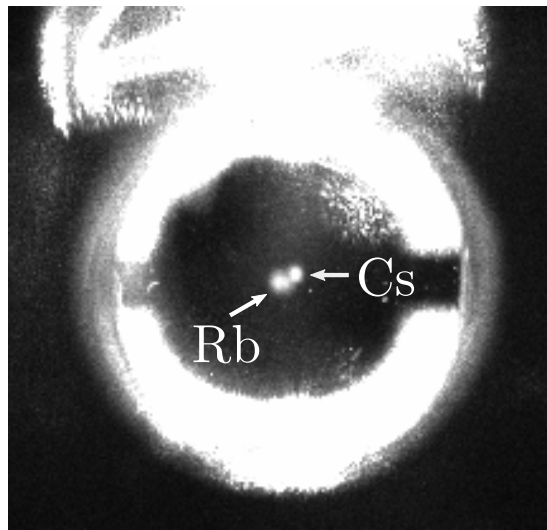


Figure 9.2: A dual-species MOT of Rb and Cs. The two clouds are not overlapped due to the retro-reflected horizontal beam creating standing waves with different node locations for 780 nm and 852 nm.

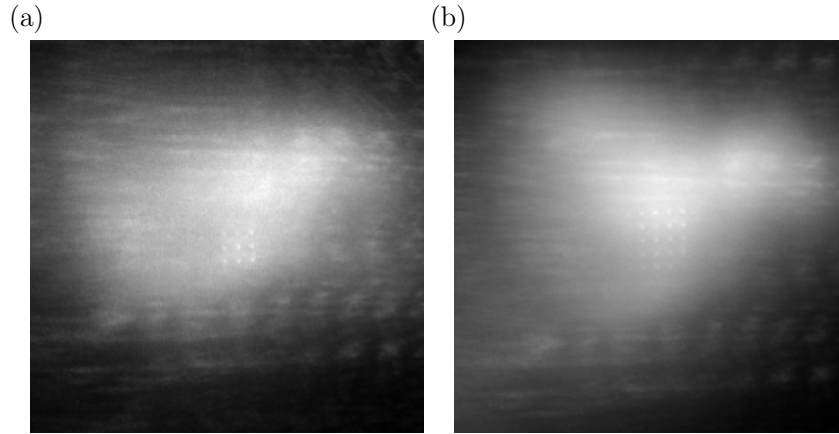


Figure 9.3: (a) 940 nm in a 3 by 3 array projected onto a Rb MOT demonstrating localisation of atoms in deep traps. (b) 810 nm in a 5 by 5 array demonstrating localisation in shallower traps.

The 940 nm system for Cs trapping was the first of the two to be brought online. This system was initially used to demonstrate atom localisation when the 940 nm traps were projected onto the Cs MOT cloud using a 5 by 5 array with 3 mK trap depths. The 940 nm ODT creates an attractive potential in the Rb $5S_{1/2}$ ground state but with trap depths four times shallower than that of Cs. This means that we should be able to trap Rb with a 940 nm ODT with deep enough traps. We first attempted to observe localisation of Rb in 940 nm in the same configuration as performed for Cs which corresponds to a trap depth of 0.75 mK for which we saw no signs of localisation. This indicates species selective trapping is possible with this wavelength of trap. Suitably cold Rb atoms can still be trapped at this trap depth, however, in future this experiment will operate at lower trap powers which should mitigate this risk. Previous work has indicated trapping of Cs in 0.2 mK trap depths which will effectively eliminate the risk of cross loading with Rb [18]. We next switched the 940 nm array to a 3 by 3 pattern which equates to a Rb trap depth of ~ 2 mK. With this array projected on a Rb MOT we were able to observe localisation when the trap light was chopped at 1 MHz with 50% duty cycle. An image of this is shown in Fig. 9.3a. The trap chopping is required

as the Stark shifts on the $5p_{3/2}$ excited states render the atom anti-trapped. Trap chopping was not required to observe localisation in Cs due to the near magic light shifts for 940 nm on the cooling transition.

After the 810 nm laser and SLM system were aligned, we performed the same localisation test for a 5 by 5 array with a trap depth of 2 mK. In this configuration we loaded a Rb MOT and projected the traps onto the cloud and were able to observe localisation when the trap laser was chopped as shown in Fig. 9.3b. The localisation of Rb at both 810 nm and 940 nm allowed us to ensure both species are trapped and imaged in the same plane.

The stochastic nature of loading tweezer arrays means that a single ODT has a 50% chance of trapping a single atom. As such, experiments require the addition of a sorting beam to ensure defect free arrays can be generated. As of writing, the laser system has been implemented on the table along with an acousto-optic deflector (AOD). As yet, this system has not been tested beyond rudimentary optimisations.

9.2 Towards Coherent Rydberg Operations

The long term ambition of this experiment is to realise error correction protocols on a dual-species platform. Towards this aim, one of the goals discussed in the introduction is to utilise enhanced Rydberg coupling strength between different atomic species for increased physical qubit gate fidelities. Chapter 4 described the emergence of zero field Förster resonances between pair states of Rb and Cs through careful choice of target Rydberg states. Building on work previously carried out in s -orbitals [132], we identified candidate Förster resonance states which possess strong Rb-Cs couplings with suppressed Rb-Rb and Cs-Cs interactions in Rydberg d -orbitals [20]. Analysis of these states using blockade leakage indicated that d -orbitals possess enhanced blockade strength when compared to s -orbitals. One pair state, the Rb $59d_{5/2}$ -Cs $68d_{3/2}$, was identified as a candidate for implementation within this experiment. We performed modelling of a gate sequence using the candidate state which indicated an upper two-qubit gate fidelity of 0.9953. Using a better optimised gate sequence should push this

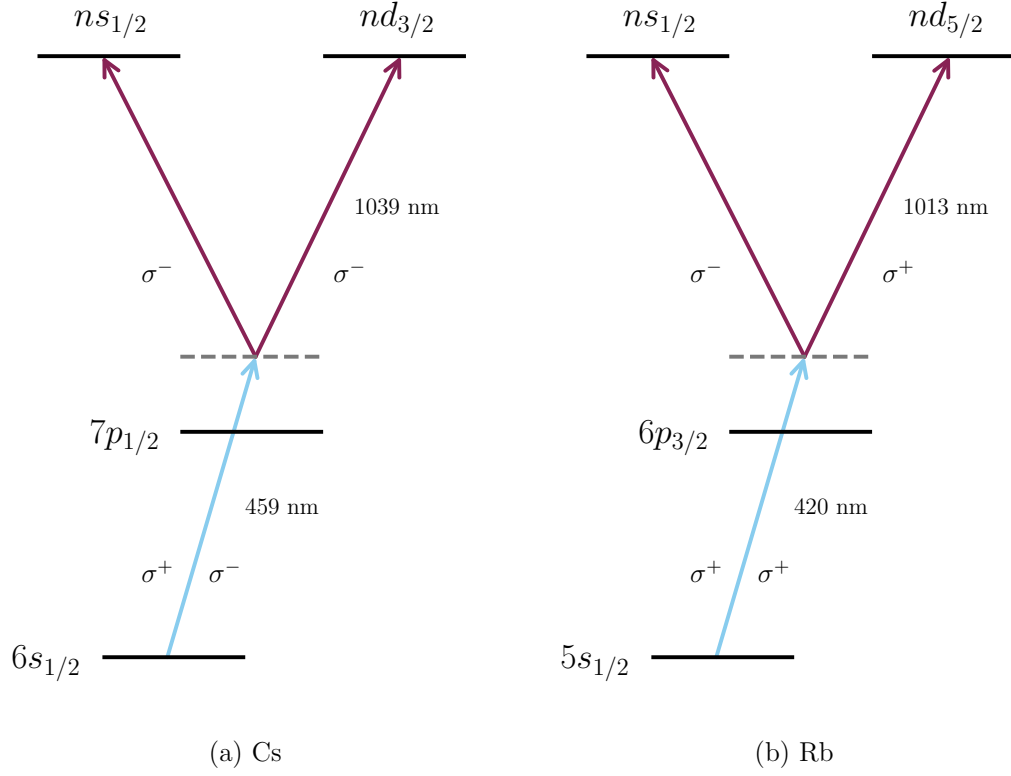


Figure 9.4: The two-photon Rydberg excitation scheme for Cs (a) and Rb (b) with wavelengths and required transitions to reach each ℓ state. For both species, s - and d -orbital Rydberg states can be addressed by changing the polarisations of one of the wavelengths.

upper bound even higher.

Following the loading of dual-species arrays, the next goal of this experiment is to realise d -orbital Rydberg excitation. To this end, work has already commenced in commissioning laser systems to address Rydberg states. Both atoms Rydberg states will be addressed via a two-photon transition for which we use two lasers per species. The target Rydberg state for Cs is the $nd_{3/2}$, $m_j = -3/2$ which we shall address via the $7p_{1/2}$, $m_j = -1/2$ with a pair of lasers at 459 nm and 1039 nm both on σ^- -transitions. The target Rydberg state for Rb is the $nd_{5/2}$, $m_j = +5/2$ which we shall address via the $6p_{3/2}$, $m_j = +3/2$ with a pair of lasers at 420 nm and 1013 nm both through a σ^+ -transitions. This pathway will also allow for excitation to Rydberg s -orbitals by

switching the blue transition of Cs to a σ^+ and the infrared transition of Rb to a σ^- .

All four Rydberg lasers are diode lasers with the infrared wavelengths amplified by a pair of fibre amplifiers. Each of the lasers are locked to an optical cavity using the Pound-Drever-Hall method [229]. We have two cavities, one for the blue lasers and one for infrared lasers. It is essential to ensure that two-photon excitation is performed using lasers with both a narrow linewidth and low phase noise. To reduce the phase noise of our Rydberg lasers, we plan to implement feed-forward system to provide real time cancellation of additional servo noise added by the optical cavity lock [97, 230–232].

As of writing, the Rydberg laser systems are being built with 1039 nm and 459 nm successfully cavity-locked. Following the development of this system, the experimental aims will be to realise Rydberg Rabi oscillations for each species, perform Rydberg spectroscopy to demonstrate the emergence of the resonant dipole-dipole (dd) interaction, and to perform a qubit gate sequence to realise the expected enhanced gate fidelities afforded by the d -orbital Förster resonances.

9.3 Future System Upgrades

To obtain lower vacuum pressure, and hence, enhanced trap lifetimes, the system presented in this thesis implemented a cryostat and titanium sublimation pump (TSP) to remove impurities from the vacuum system. Trap lifetime scales as $1/N$ for N atoms, therefore, longer trap lifetimes are a requirement to perform quantum computation on a neutral atom platform at scale. This system has implemented an aspheric lens which limits the field of view (FOV) to $\sim 120 \mu\text{m}$ which equates to a best case of ~ 550 atoms in a defect-free chequerboard array. For scaling to array sizes of thousands of atoms, the optical system will need upgrading to a microscope objective which will allow for FOVs up to the millimetre length scales [8]. This upgrade will also allow for reduced array spacing further increasing the number of trappable atoms.

Implementation of a microscope objective to this system has already been discussed in Sect. 8.4.4 where consideration must be made in how best to integrate it with a

cryogenic system. In this vain, future upgrades to the system provide an opportunity to make changes to the design of the cryostat to ensure optimal performance on larger arrays. In Sect. 6.2 we discuss the design of the low-vibration interface between the coldhead and vacuum system in which the heat exchanges are physically separated. Despite this interface, there is still vibrational coupling between the coldhead and the table which is transferred through the vacuum system leading to some optical instability on 10 –20 Hz timescales. We have mitigated these effects by making careful alignment to optics systems around the cryostat, however, future upgrades to the system should involve redesigning of a low vibration system for the cryostat. Other cryogenic tweezer arrays have opted to thermally couple their cryogenic stage to the cryostat using copper braiding [9, 121, 123], which greatly dampens the coupling of the dominant frequencies modes from the cryostat.

Copying the cryostat designs presented in Ref. [9] and Ref. [123] is not necessarily feasible as it would require a complete redesign of the coldhead and heat exchange systems. A low vibration cryostat design presented in Ref. [233] has a similar design to the one we use¹ consisting of a closed cycle coldhead with a void surrounding it which is filled with a buffer gas of He. The experiment presented is for a cryogenic cavity, but the design principles can be used in a tweezer array system. The heat exchangers consist of a pair of copper fins and heat is transferred between the coldhead and experimental stage through the buffer gas. This ensures no physical contact is made between the vibrating coldhead and the vacuum system. There are two heat exchanges, an upper one with a base temperature of 90 K which is thermally coupled to a shield; a lower heat exchange can reach a base temperature of 4 K and is used to cool a cold finger. The vacuum system is made up of two sections joined by a flexible bellow to reduce vibrational couplings between the cryostat side and experimental side of the system. The 4 K cold finder extends through the bellow and is completely surrounded by the 90 K shield. Inside the science chamber, the shield stands on a support structure. Inside this portion of the shielding, the cryogenic stage of the experiment sits on a second support structure and is thermally anchored to the cold

¹see Fig. 6.4 in Sect. 6.2.

finger via copper braids. To further reduce vibration coupling to the science chamber, an extensive support structure is build around the cryostat which sits off the optical table. This setup obtains a vibrational suppression of 25 dB.

Implementing some of the features mentioned above — namely using copper braids and moving the cryostat away from the optical table — should reduce the optical instabilities we observed in our system. In terms of cryogenic mount designs, minimising the solid angle of room temperature surfaces to the atoms will ensure maximal trap lifetimes.

A cryogenic stage that fully encloses the tweezer array should allow for the extension of trap lifetimes to the thousands of second range which allow scaling to larger arrays. Additionally, this shielding should increase the Rydberg state lifetime and hence, improvements in gate fidelities [124]. Recent work has demonstrated cryogenic tweezer arrays with 5000 s lifetimes using a microscope objective [125].

9.4 Conclusion

This thesis has presented progress towards realisation of a dual-species tweezer array in a cryogenic apparatus. As of writing only a Cs array has been realised, however, work towards a dual-species array is well underway. Although, additional design work must be carried out to take full advantage of enhanced lifetimes, this system will be able to act as a testbed for interesting physics. As such, upgrades to the system discussed above will not be implemented in the near-term. Future research with this experiment will focus on the Förster interaction between Rydberg Rb – Cs d -states discussed in Chapter 4. It is expected that these states should allow for the realisation of high fidelity multi-qubit gates, which opens an avenue of future research for implementing primitive error correction code patches on this system. However, scaling to more than a few logical qubits will require the implementation of a microscope objective. Beyond gate based quantum computing, this platform will also open up opportunities to study interesting many-body physics [16, 104] as we gain as the Rydberg interaction Hamiltonian for a dual-species system has access to two interaction terms.

Appendices

Appendix A

d -orbital Channel Coupling Table

Rydberg pair states coupling channels for all $n_{\text{Rb}}d_{m_j^{\text{Rb}}} \otimes n_{\text{Cs}}d_{m_j^{\text{Cs}}}$ combinations. Each configuration has 9 coupling channels k listed across four tables below.

Target state	Channel k	Channel coupling state
$n_{\text{Rb}}d_{5/2} \otimes n_{\text{Cs}}d_{5/2}$	1	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}p_{3/2}$
	2	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}p_{3/2}$
	3	$n'_{\text{Rb}}f_{7/2} \otimes n'_{\text{Cs}}p_{3/2}$
	4	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}f_{5/2}$
	5	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}f_{7/2}$
	6	$n'_{\text{Rb}}f_{3/2} \otimes n'_{\text{Cs}}f_{3/2}$
	7	$n'_{\text{Rb}}f_{7/2} \otimes n'_{\text{Cs}}f_{3/2}$
	8	$n'_{\text{Rb}}f_{3/2} \otimes n'_{\text{Cs}}f_{7/2}$
	9	$n'_{\text{Rb}}f_{7/2} \otimes n'_{\text{Cs}}f_{7/2}$

Table A.1

Appendix A. *d*-orbital Channel Coupling Table

Target state	Channel k	Channel coupling state
$n_{\text{Rb}}d_{3/2} \otimes n_{\text{Cs}}d_{5/2}$	1	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}p_{3/2}$
	2	$n'_{\text{Rb}}p_{1/2} \otimes n'_{\text{Cs}}p_{3/2}$
	3	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}p_{3/2}$
	4	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}f_{5/2}$
	5	$n'_{\text{Rb}}p_{1/2} \otimes n'_{\text{Cs}}f_{5/2}$
	6	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}f_{7/2}$
	7	$n'_{\text{Rb}}p_{1/2} \otimes n'_{\text{Cs}}f_{7/2}$
	8	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}f_{5/2}$
	9	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}f_{7/2}$

Table A.2

Target state	Channel k	Channel coupling state
$n_{\text{Rb}}d_{5/2} \otimes n_{\text{Cs}}d_{3/2}$	1	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}p_{3/2}$
	2	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}p_{1/2}$
	3	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}f_{5/2}$
	4	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}p_{3/2}$
	5	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}p_{1/2}$
	6	$n'_{\text{Rb}}f_{7/2} \otimes n'_{\text{Cs}}p_{3/2}$
	7	$n'_{\text{Rb}}f_{7/2} \otimes n'_{\text{Cs}}p_{1/2}$
	8	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}f_{5/2}$
	9	$n'_{\text{Rb}}f_{7/2} \otimes n'_{\text{Cs}}f_{5/2}$

Table A.3

Target state	Channel k	Channel coupling state
$n_{\text{Rb}}d_{3/2} \otimes n_{\text{Cs}}d_{3/2}$	1	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}p_{3/2}$
	2	$n'_{\text{Rb}}p_{1/2} \otimes n'_{\text{Cs}}p_{3/2}$
	3	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}p_{1/2}$
	4	$n'_{\text{Rb}}p_{1/2} \otimes n'_{\text{Cs}}p_{1/2}$
	5	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}p_{3/2}$
	6	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}p_{1/2}$
	7	$n'_{\text{Rb}}p_{3/2} \otimes n'_{\text{Cs}}f_{5/2}$
	8	$n'_{\text{Rb}}p_{1/2} \otimes n'_{\text{Cs}}f_{5/2}$
	9	$n'_{\text{Rb}}f_{5/2} \otimes n'_{\text{Cs}}f_{5/2}$

Table A.4

Appendix B

Candidate Förster Resonant Rb-Cs Pair States

The following tables list candidate Förster resonant pair states identified through asymptotic approximation of the $C_{3,k}$ dispersion coefficient and smallest energy defect Δ (see Chapter 4 for more details). The target states listed have the conditions of $50 \leq n \leq 90$, $|C_{3,k}| > 1 \text{ GHz } \mu\text{m}^3$ and $|\Delta| < 0.005$ of the level spacing. Each identified pair state is presented with the blockade leakage calculated as described in Sect. 4.2.2 using eigenstate values obtained through direct matrix diagonalization performed with PairInteraction [164]. For Rb-Cs we calculate the blockade leakage for the $|1r\rangle$ two-atom basis for a resonant π -pulse ($P_{1r}(t_\pi)$) at an interatomic spacing of $R = 6 \mu\text{m}$. For the Rb-Rb and Cs-Cs we calculate blockade leakage for the $|rr\rangle$ two-atom basis for a resonant π -pulse ($P_{rr}(t_\pi)$) at $R = 8.5 \mu\text{m}$. Each interspecies interaction is presented with the fitted $\tilde{C}_6$ coefficient for the range $4 < R < 15 \mu\text{m}$.

Rb - Cs Förster Resonances		$\Delta/2\pi$	$C_{3,k}$	Rb-Cs at $R = 6 \mu\text{m}$		Rb-Rb at $R = 8.5 \mu\text{m}$		Cs-Cs at $R = 8.5 \mu\text{m}$	
Target Pair	Resonant Pair	(MHz)	(GHz μm^3)	$P_{1r}(t_\pi)$	$\tilde{C}_3$ (GHz μm^3)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)
$55d_{5/2} - 53d_{5/2}$	$57p_{3/2} - 50f_{7/2}$	18.71	-2.73	0.9860	-3.13	0.9989	21.4	0.9999	-7.4
$55d_{5/2} - 63d_{5/2}$	$56p_{3/2} - 61f_{5/2}$	12.66	-2.24	0.9969	-11.30	0.9989	21.5	0.9840	-98.8
$55d_{5/2} - 61d_{5/2}$	$53f_{7/2} - 63p_{3/2}$	-12.19	-1.78	0.9747	2.08	0.9989	21.4	0.9923	-62.1
$62d_{5/2} - 71d_{5/2}$	$63p_{3/2} - 69f_{5/2}$	-9.53	-3.66	0.9997	15.79	0.9855	64.1	0.7747	-416.8
$69d_{5/2} - 79d_{5/2}$	$70p_{3/2} - 77f_{7/2}$	-19.14	-21.97	0.9998	26.34	0.7855	192.8	0.0013	-1448.9
$69d_{5/2} - 66d_{5/2}$	$71p_{3/2} - 63f_{5/2}$	-4.29	-1.79	0.9933	7.49	0.7855	192.8	0.9538	-164.5
$77d_{5/2} - 88d_{5/2}$	$78p_{3/2} - 86f_{7/2}$	1.91	-34.24	0.9998	-47.42	0.2966	393.5	0.0282	-4303.1
$74d_{5/2} - 82d_{5/2}$	$72f_{7/2} - 84p_{3/2}$	5.63	-6.15	0.9967	-16.85	0.3269	367.5	0.1115	-2258.4
$70d_{5/2} - 80d_{5/2}$	$71p_{3/2} - 78f_{7/2}$	13.69	-23.21	0.9990	-29.03	0.7442	223.9	0.0238	-1812.1
$76d_{5/2} - 87d_{5/2}$	$77p_{3/2} - 85f_{5/2}$	-19.30	-8.41	0.9997	40.58	0.1156	450.0	0.0388	-4482.5
$84d_{5/2} - 80d_{5/2}$	$86p_{3/2} - 77f_{7/2}$	11.74	-15.58	0.9996	-24.33	0.0514	1318.7	0.0148	-1692.0
$80d_{5/2} - 70d_{5/2}$	$78f_{7/2} - 68f_{5/2}$	-18.70	3.05	0.9932	2.38	0.0229	698.0	0.8301	-356.6
$83d_{5/2} - 79d_{5/2}$	$85p_{3/2} - 76f_{7/2}$	-10.78	-14.80	0.9975	12.46	0.0957	1108.5	0.0013	-1448.9

Table B.1: Rb-Cs Förster resonances from Rb $d_{5/2}$ - Cs $d_{5/2}$ states with $|\Delta|/2\pi < 20$ MHz and $|C_{3,k}| > 1$ GHz μm^3 . $P_{1r}(t_\pi)$ and $P_{rr}(t_\pi)$ along with fitted $\tilde{C}_3$ and $\tilde{C}_6$ coefficients are calculated for $\theta = 90^\circ$ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (5/2, -5/2)$ assuming $\Omega/2\pi = 2$ MHz.

Rb - Cs Förster Resonances		$\Delta/2\pi$	$C_{3,k}$	Rb-Cs at $R = 6 \mu\text{m}$		Rb-Rb at $R = 8.5 \mu\text{m}$		Cs-Cs at $R = 8.5 \mu\text{m}$	
Target Pair	Resonant Pair	(MHz)	(GHz μm^3)	$P_{1r}(t_\pi)$	$\tilde{C}_3$ (GHz μm^3)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)
51 $d_{3/2}$ - 59 $d_{3/2}$	52 $p_{3/2}$ - 57 $f_{5/2}$	-18.12	2.08	0.9912	0.83	0.9995	13.8	0.9995	-15.8
50 $d_{3/2}$ - 57 $d_{3/2}$	51 $p_{1/2}$ - 55 $f_{5/2}$	16.13	-5.82	0.9958	-6.33	0.9997	11.4	0.9998	-10.9
51 $d_{3/2}$ - 62 $d_{3/2}$	50 $f_{5/2}$ - 59 $f_{5/2}$	-15.98	4.52	0.9814	3.15	0.9995	13.8	0.9995	-13.2
65 $d_{3/2}$ - 75 $d_{3/2}$	66 $p_{3/2}$ - 73 $f_{5/2}$	8.32	5.58	0.9955	-36.55	0.7416	234.6	0.7704	-440.2
69 $d_{3/2}$ - 66 $d_{3/2}$	71 $p_{1/2}$ - 63 $f_{5/2}$	-6.16	-7.35	0.9979	6.81	0.4719	224.4	0.9706	-141.0
64 $d_{3/2}$ - 61 $d_{3/2}$	66 $p_{3/2}$ - 58 $f_{5/2}$	3.96	1.64	0.8913	-0.35	0.7767	211.9	0.9939	-63.6
65 $d_{3/2}$ - 74 $d_{3/2}$	66 $p_{1/2}$ - 72 $f_{5/2}$	-13.07	-17.02	0.9974	16.67	0.7416	234.6	0.8297	-375.1
76 $d_{3/2}$ - 67 $d_{3/2}$	74 $f_{5/2}$ - 65 $f_{5/2}$	2.04	8.99	0.9868	-12.99	0.0129	723.9	0.9587	-163.3
77 $d_{3/2}$ - 73 $d_{3/2}$	79 $p_{3/2}$ - 70 $f_{5/2}$	-6.12	3.50	0.9789	3.82	0.0692	820.6	0.8729	-308.5
71 $d_{3/2}$ - 79 $d_{3/2}$	69 $f_{5/2}$ - 81 $p_{3/2}$	9.09	1.63	0.9944	-11.82	0.2858	329.2	0.3238	-545.4
70 $d_{3/2}$ - 77 $d_{3/2}$	68 $f_{5/2}$ - 79 $p_{1/2}$	7.97	-5.30	0.9961	-10.84	0.3789	287.4	0.6138	-629.1
72 $d_{3/2}$ - 83 $d_{3/2}$	73 $p_{3/2}$ - 81 $f_{5/2}$	10.33	8.46	0.9985	-42.76	0.2049	390.5	0.1082	-2156.1
79 $d_{3/2}$ - 91 $d_{3/2}$	80 $p_{3/2}$ - 89 $f_{5/2}$	10.38	12.33	0.9998	-51.08	0.1280	1112.0	0.0234	-3178.2
78 $d_{3/2}$ - 90 $d_{3/2}$	79 $p_{3/2}$ - 88 $f_{5/2}$	-10.66	11.75	0.9998	21.39	0.1607	957.1	0.0262	-2568.1
73 $d_{3/2}$ - 83 $d_{3/2}$	74 $p_{1/2}$ - 81 $f_{5/2}$	-0.70	-27.26	0.9995	28.54	0.1360	452.0	0.1082	-2156.1
81 $d_{3/2}$ - 92 $d_{3/2}$	82 $p_{1/2}$ - 90 $f_{5/2}$	4.39	-41.55	0.9997	-47.30	0.0590	1411.2	0.0153	-3836.9
84 $d_{3/2}$ - 74 $d_{3/2}$	82 $f_{5/2}$ - 72 $f_{5/2}$	-8.89	13.56	0.9906	3.10	0.0368	1937.6	0.6502	-484.5
84 $d_{3/2}$ - 80 $d_{3/2}$	86 $p_{1/2}$ - 77 $f_{5/2}$	11.47	-16.47	0.9953	-24.02	0.0368	1937.6	0.7637	-388.1
80 $d_{3/2}$ - 89 $d_{3/2}$	78 $f_{5/2}$ - 91 $p_{3/2}$	4.98	2.68	0.9966	-21.96	0.0440	1263.6	0.0135	-2640.7
80 $d_{3/2}$ - 88 $d_{3/2}$	78 $f_{5/2}$ - 90 $p_{1/2}$	-3.06	-9.23	0.9357	9.12	0.0440	1263.6	0.0007	-2226.4
80 $d_{3/2}$ - 91 $d_{3/2}$	81 $p_{1/2}$ - 89 $f_{5/2}$	-13.66	-39.62	0.9998	48.51	0.0440	1263.6	0.0234	-3178.2
83 $d_{3/2}$ - 79 $d_{3/2}$	85 $p_{1/2}$ - 76 $f_{5/2}$	-10.65	-15.65	0.9983	10.75	0.0603	1610.2	0.0904	-1035.9

Table B.2: Rb-Cs Förster resonances from Rb $d_{3/2}$ - Cs $d_{3/2}$ states with $|\Delta|/2\pi < 20$ MHz and $|C_{3,k}| > 1$ GHz μm^3 . $P_{1r}(t_\pi)$ and $P_{rr}(t_\pi)$ along with fitted $\tilde{C}_3$ and $\tilde{C}_6$ coefficients are calculated for $\theta = 90^\circ$ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (3/2, -3/2)$ assuming $\Omega/2\pi = 2$ MHz.

Rb - Cs Förster Resonances		$\Delta/2\pi$	$C_{3,k}$	Rb-Cs at $R = 6 \mu\text{m}$		Rb-Rb at $R = 8.5 \mu\text{m}$		Cs-Cs at $R = 8.5 \mu\text{m}$	
Target Pair	Resonant Pair	(MHz)	(GHz μm^3)	$P_{1r}(t_\pi)$	$\tilde{C}_3$ (GHz μm^3)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)
$59d_{5/2} - 68d_{3/2}$	$60p_{3/2} - 66f_{5/2}$	13.40	-11.23	0.9992	-14.36	0.9951	41.6	0.9410	-206.1
$50d_{5/2} - 37d_{3/2}$	$49f_{7/2} - 38p_{1/2}$	-5.91	-2.93	0.9864	3.87	0.9999	8.2	1.0000	-0.2
$52d_{5/2} - 60d_{3/2}$	$53p_{3/2} - 58f_{5/2}$	15.19	-6.70	0.9964	-8.37	0.9997	12.2	0.9958	-45.0
$66d_{5/2} - 76d_{3/2}$	$67p_{3/2} - 74f_{5/2}$	11.12	-17.74	0.9997	-22.84	0.9176	114.4	0.4467	-692.9
$68d_{5/2} - 60d_{3/2}$	$66f_{5/2} - 58f_{5/2}$	-8.50	1.52	0.9272	3.71	0.8373	170.4	0.9993	-18.9
$60d_{5/2} - 66d_{3/2}$	$58f_{7/2} - 68p_{1/2}$	-14.00	-2.87	0.9906	2.65	0.9929	48.3	0.9698	-144.1
$67d_{5/2} - 49d_{3/2}$	$66f_{7/2} - 50p_{1/2}$	-5.80	-9.64	0.9932	11.22	0.8889	142.5	1.0000	-2.7
$78d_{5/2} - 74d_{3/2}$	$80p_{3/2} - 71f_{5/2}$	2.73	-11.14	0.9977	-16.52	0.0056	557.1	0.8297	-375.1
$79d_{5/2} - 91d_{3/2}$	$80p_{3/2} - 89f_{5/2}$	-12.53	-36.95	0.9999	49.07	0.0070	619.2	0.0052	-3596.5
$73d_{5/2} - 84d_{3/2}$	$74p_{3/2} - 82f_{5/2}$	9.09	-26.72	0.9949	-36.07	0.4852	314.6	0.1085	-2230.7
$77d_{5/2} - 56d_{3/2}$	$76f_{7/2} - 57p_{1/2}$	-20.99	-16.90	0.9912	17.86	0.2966	393.5	0.9996	-14.4
$81d_{5/2} - 90d_{3/2}$	$79f_{7/2} - 92p_{3/2}$	16.57	2.90	0.9953	-19.25	0.0649	765.4	0.0024	-3282.6
$83d_{5/2} - 73d_{3/2}$	$81f_{5/2} - 71f_{5/2}$	15.70	3.43	0.9939	-12.68	0.0957	1108.5	0.8729	-308.5
$81d_{5/2} - 89d_{3/2}$	$79f_{5/2} - 91p_{1/2}$	5.56	-2.58	0.9802	-23.45	0.0858	839.1	0.0261	-2866.1
$84d_{5/2} - 61d_{3/2}$	$86f_{7/2} - 61p_{3/2}$	-6.70	-0.00	0.9991	26.62	0.0774	1227.1	0.9939	-63.6
$80d_{5/2} - 92d_{3/2}$	$81p_{3/2} - 90f_{5/2}$	7.42	-38.76	0.9999	-51.12	0.0233	705.2	0.0144	-5756.9

Table B.3: Rb-Cs Förster resonances from Rb $d_{5/2}$ - Cs $d_{3/2}$ states with $|\Delta|/2\pi \lesssim 20$ MHz and $|C_{3,k}| > 1$ GHz μm^3 . $P_{1r}(t_\pi)$ and $P_{rr}(t_\pi)$ along with fitted $\tilde{C}_3$ and $\tilde{C}_6$ coefficients are calculated for $\theta = 90^\circ$ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (5/2, -3/2)$ assuming $\Omega/2\pi = 2$ MHz.

Rb - Cs Förster Resonances		$\Delta/2\pi$	$C_{3,k}$	Rb-Cs at $R = 6 \mu\text{m}$		Rb-Rb at $R = 8.5 \mu\text{m}$		Cs-Cs at $R = 8.5 \mu\text{m}$	
Target Pair	Resonant Pair	(MHz)	(GHz μm^3)	$P_{1r}(t_\pi)$	$\tilde{C}_3$ (GHz μm^3)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)
$59d_{3/2} - 57d_{5/2}$	$61p_{1/2} - 54f_{5/2}$	2.93	-1.03	0.9823	-5.56	0.6608	187.8	0.9981	-32.2
$53d_{3/2} - 60d_{5/2}$	$54p_{1/2} - 58f_{5/2}$	-4.41	-1.97	0.9968	7.12	0.9989	19.7	0.9938	-57.0
$64d_{3/2} - 71d_{5/2}$	$62f_{5/2} - 73p_{3/2}$	4.08	-3.26	0.9909	-7.31	0.7767	211.9	0.7383	-435.9
$78d_{3/2} - 88d_{5/2}$	$79p_{1/2} - 86f_{5/2}$	6.51	-9.49	0.9915	-39.56	0.1607	957.1	0.0085	-5492.7
$76d_{3/2} - 87d_{5/2}$	$77p_{3/2} - 85f_{5/2}$	6.49	2.81	0.9986	-52.91	0.0129	723.9	0.0186	-5038.3
$80d_{3/2} - 70d_{5/2}$	$78f_{5/2} - 68f_{5/2}$	3.69	2.95	0.9920	-11.62	0.0440	1263.6	0.8048	-377.9
$83d_{3/2} - 79d_{5/2}$	$85p_{3/2} - 76f_{7/2}$	8.93	4.92	0.9918	-32.33	0.0522	1813.4	0.0000	-1466.9
$90d_{3/2} - 86d_{5/2}$	$92p_{1/2} - 83f_{5/2}$	2.55	-5.81	0.9961	-40.04	0.0066	3762.9	0.0201	-4212.1

Table B.4: Rb-Cs Förster resonances from Rb $d_{3/2}$ - Cs $d_{5/2}$ states with $|\Delta|/2\pi \lesssim 20$ MHz and $|C_{3,k}| > 1$ GHz μm^3 . $P_{1r}(t_\pi)$ and $P_{rr}(t_\pi)$ along with fitted $\tilde{C}_3$ and $\tilde{C}_6$ coefficients are calculated for $\theta = 90^\circ$ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (3/2, -5/2)$ assuming $\Omega/2\pi = 2$ MHz.

Rb - Cs Förster Resonances		$\Delta/2\pi$	$C_{3,k}$	Rb-Cs at $R = 6 \mu\text{m}$		Rb-Rb at $R = 8.5 \mu\text{m}$		Cs-Cs at $R = 8.5 \mu\text{m}$	
Target Pair	Resonant Pair	(MHz)	(GHz μm^3)	$P_{1r}(t_\pi)$	$\tilde{C}_3$ (GHz μm^3)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)	$P_{rr}(t_\pi)$	$\tilde{C}_6$ (GHz μm^6)
$48s_{1/2} - 51s_{1/2}$	$48p_{3/2} - 50p_{1/2}$	-5.53	1.69	0.6281	3.00	0.9998	-9.5	0.9996	-14.5
$59s_{1/2} - 57s_{1/2}$	$58p_{1/2} - 57p_{1/2}$	-16.45	3.54	0.3752	1.35	0.9767	-114.2	0.9943	-56.6
$69s_{1/2} - 68s_{1/2}$	$68p_{1/2} - 68p_{3/2}$	-9.89	6.92	0.9959	12.69	0.3230	-701.2	0.6503	-430.2
$68s_{1/2} - 67s_{1/2}$	$67p_{1/2} - 67p_{3/2}$	2.63	6.50	0.9727	-9.12	0.4660	-594.3	0.7432	-367.5
$61s_{1/2} - 65s_{1/2}$	$61p_{1/2} - 64p_{1/2}$	2.77	4.80	0.9735	-2.62	0.9487	-168.9	0.8689	-263.5
$72s_{1/2} - 75s_{1/2}$	$72p_{1/2} - 74p_{3/2}$	2.70	9.65	0.9835	-12.40	0.0040	-1106.6	0.0359	-1187.3
$77s_{1/2} - 81s_{1/2}$	$77p_{3/2} - 80p_{1/2}$	-2.10	12.28	0.9969	16.42	0.0100	-2103.4	0.0344	-2151.8
$72s_{1/2} - 70s_{1/2}$	$71p_{3/2} - 70p_{1/2}$	-7.90	8.51	0.9992	9.82	0.0040	-1106.6	0.4005	-605.3
$71s_{1/2} - 69s_{1/2}$	$70p_{3/2} - 69p_{1/2}$	9.45	8.01	0.9987	-10.66	0.0643	-955.7	0.5348	-518.5

Table B.5: Rb-Cs Förster resonances from Rb $s_{1/2}$ - Cs $s_{1/2}$ states previously identified [132]. $P_{1r}(t_\pi)$ and $P_{rr}(t_\pi)$ along with fitted $\tilde{C}_3$ and $\tilde{C}_6$ coefficients are calculated for $\theta = 90^\circ$ and $(m_j^{\text{Rb}}, m_j^{\text{Cs}}) = (1/2, -1/2)$ assuming $\Omega/2\pi = 2$ MHz.

Appendix C

2D MOT Velocity Selection

As discussed in Sect. 6.1.1 and Chapter 7, the 2D magneto-optical trap (MOT) was unable to load the 3D MOT when the experiment was initial built. This was due in part to atoms not reaching the 3D MOT at suitably low velocities to be trapped. Below we discuss how we addressed this issue through introducing a tilt to the 2D MOT vacuum chamber.

Fig. C.1 shows a cut view of the vacuum system with the 2D MOT chamber on the left and science chamber on the right. When the 2D MOT chamber sits flat on the table, atoms have to travel along the red arrow through the differential pumping tube to reach the 3D MOT. The total distance between the front 2D MOT region and the 3D MOT is 592 mm, over which atoms will follow a parabolic trajectory.

We modelled the ballistic trajectories of atom disseminating from the 2D MOT to identify which velocity classes reach the 3D MOT. This is shown in Fig. C.2 where the green lines indicate the differential pumping tube, the horizontal black lines are the vacuum walls and the vertical mot lines are is the copper mount. The trajectory of the push beam is represented by the dashed grey line which is set to pass through the centre of the 2D MOT and 3D MOT. The red ballistic trajectories show that all velocity classes pass through the differential pumping tube unimpeded, however, only the atoms with a velocity of $> 20 \text{ ms}^{-1}$ reach the 3D MOT while the rest are blocked by the vacuum walls or copper mount. Analysis described in Ref. [180] indicates that most atoms in a 2D MOT have a velocity of $7 \text{ ms}^{-1} - 20 \text{ ms}^{-1}$, which implies our MOT

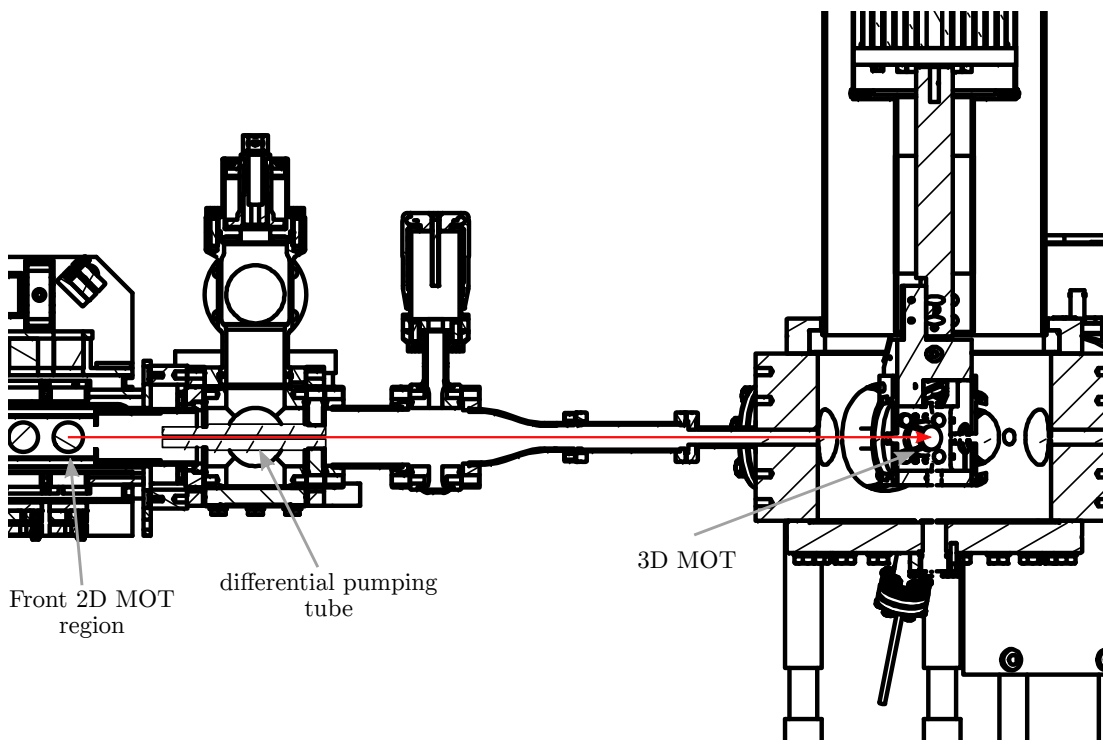


Figure C.1: Cut view of the vacuum system shown the trajectory of atoms from the front region of the 2D MOT which travel along the red line to the 3D MOT at the centre of the cryogenic stage.

setup is not suitable for loading.

As discussed in Sect. 6.1.1, we chose to put the 2D MOT vacuum chamber on a tilt to ensure slower atoms can reach the 3D MOT. Ref. [181] and Ref. [182] suggest that for a 10 G/cm quadrupole field gradient and a cooling detuning of $\delta = -3\Gamma$, a Cs 3D MOT should have capture velocities of $< 18\text{m/s}$. As such, we chose to optimise the 2D MOT's angle to ensure velocity classes $< 20\text{m/s}$ would reach the 3D MOT. Fig. C.3 shows the ballistic trajectory of atoms from the 2D MOT for a chamber angle of 2° . The results indicate that atoms with a trajectory $\geq 15\text{ m/s}$ should reach the 3D MOT region unimpeded.

We found that the 2° angle allowed us to load the MOT. Details of how we implemented this are provided in Sect. 6.1.1.

Appendix C. 2D MOT Velocity Selection

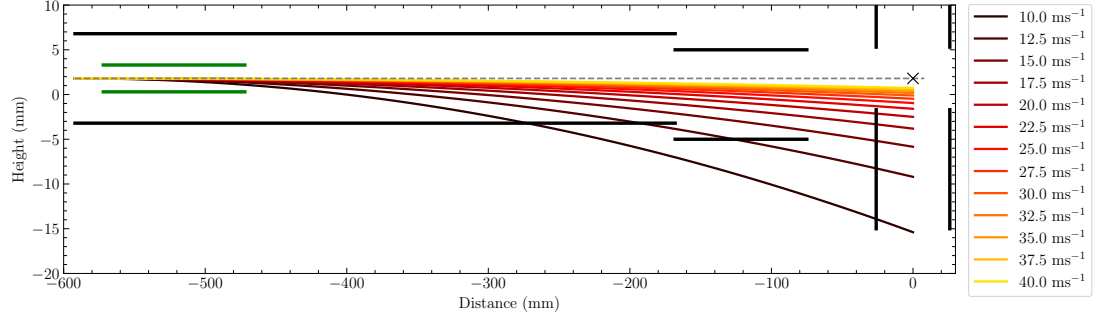


Figure C.2: Modelled parabolic trajectories of atoms from the 2D MOT at different velocities with the 2D MOT vacuum chamber flat on the table. The grey dashed line indicates the path of the push beam, the black horizontal black lines indicate the vacuum chamber walls and vertical black lines indicate the copper mount.

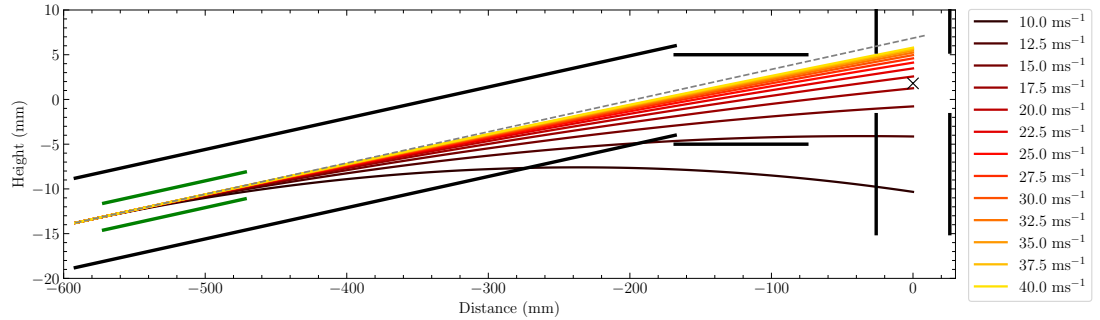


Figure C.3: Modelled parabolic trajectories of atoms from the 2D MOT at different velocities with the 2D MOT vacuum chamber at a 2° tilt on the table. The grey dashed line indicates the path of the push beam, the black horizontal black lines indicate the vacuum chamber walls and vertical black lines indicate the copper mount.

Appendix D

Cryogenic Mirror Polarisation

The cryogenic 3D MOT design is described in detail in Sect. 6.3 and Sect. 6.4. Because of the restricted optical access due to the cryostat, the vertical MOT beam enters and exits the vacuum chamber below by making two bounces off attached to the cryogenic stage. The MOT beam is retro reflected back through the chamber where it makes a further two reflections off cryogenic mirrors for a total of 4 reflections (not including the retro-reflection). The two beam axes are separated by an angle of 50° . Fig. D.1 shows the reflection path of the vertical MOT beam.

A MOT requires good circular polarisation to correctly address Zeeman shifted sub-levels and create the trapping force. Each reflection off a mirror will degrade the degree of circularity of the laser due to differences in reflectivity in the s-polarisations r_s and p-polarisations r_p . The mirrors used on the cryogenic stage are silver coated glass (Thorlabs: ME05S-P01) which have reflectivities at an incidence of 12° shown in Table D.1. To quantify the degree of circularity we use the third stokes parameter

$$\overline{S}_3 = \frac{r_s}{r_p} = \sin 2\eta, \quad (\text{D.1})$$

where 2η defines the angle from the linear equatorial line of the Poincaré sphere. A $|\overline{S}_3| = 1$ implies pure circularity while $\overline{S}_3 = 0$ implies pure linear polarisation.

We calculated the expected degradation of circular polarisation assuming a beam with $\overline{S}_3 = 1$ using reflectivities in Table D.1 at 850 nm and Eq. D.1. The result after

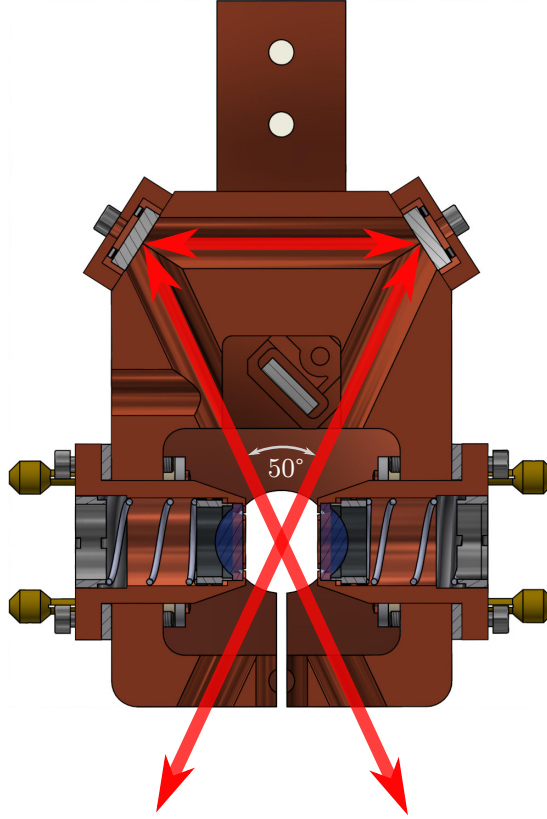


Figure D.1: Cut view of the cryogenic mount with red arrows indicating the paths for the vertical MOT beams. The beams are directed by two mirrors situated at the top of the mount.

Wavelength (nm)	Reflectivity (%)		
	p-polarisation	s-polarisation	unpolarised
780	98.40	98.03	98.21
850	97.25	96.25	96.76

Table D.1: Reflectivity of a Thorlabs ME05S-P01 silver coated mirror.

Appendix D. Cryogenic Mirror Polarisation

No. Reflections	Calculated $\bar{S}_3$
1	0.99
2	0.98
3	0.97
4	0.96

Table D.2: Calculated $\bar{S}_3$ parameter using data from Table D.1.

each of the four reflections of the internal mirror are shown in Table D.2. These values for $\bar{S}_3$ are only true for angle of incidences of 12° , but in the cryogenic system, it is really 32.5° . To get an accurate measurement of $\bar{S}_3$, we rebuild the mirror layout out of vacuum and used a polarisation analyser (Schäfter+Kirchoff: Polarization Analyser SK010PA-NIR) to obtain a value for η after two bounces. We measured $\eta = 36.7^\circ$ which corresponds to $\bar{S}_3 = 0.96$ which extrapolated this to $\bar{S}_3 = 0.92$ after 4 bounces.

To test whether a MOT can be created with this much polarisation degradation, we set one of the vertical beams of the Square experiments MOT [26] to have an $\bar{S}_3 = 0.92$ using the input waveplate. This will reduce the polarisation quality on one axis as the vertical MOT beams are retro-reflected and only make two passes through the trapping region. Under these conditions, a MOT was successfully loaded in the square experiment. As the Square experiment has identical beam geometry to this experiment this result indicates that our system should be able to load a MOT with the current degradation in polarisation.

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