

Scaling Neutral Atom Tweezer Arrays to 6,100 Qubits

Thesis by
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In Partial Fulfillment of the Requirements for the
Degree of
Doctor of Philosophy



CALIFORNIA INSTITUTE OF TECHNOLOGY
Pasadena, California

2026
Defended May 22, 2026

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ACKNOWLEDGEMENTS

Life through graduate school has been a journey, in many ways expected and in other ways unknowable. I feel incredibly grateful to my advisor Manuel Endres for his guidance and insight throughout. Through his perspective and example, I have learned to not be satisfied until ensuring that I am asking the right question, that rigor is worth the necessary patience (and enthusiasm!), and to take all things step by step. Manuel's passion for science combined with his care to cultivate the team has created a unique environment in which I have been able to grow in confidence and competence as a scientist, which I doubt I would have found anywhere else. I am enduringly grateful.

I would additionally like to thank Andrei Faraon, David Hsieh, and Austin Minnich for serving on my committee and for the contributions they have made to my growth. Particular thanks goes to Andrei for his group's collaboration in developing future scaling experiments, and to Dave for his investment in programs such as FUTURE for Women in Physics, which brought me to Caltech's campus for the first time in 2019. I would like to thank John Dabiri for his graciousness and generosity in extending mentorship on the level of integrating faith and practice. I would like to also extend deep gratitude to Bryce Gadway and Richard Scalettar for their invaluable mentorship in my college years as I started out in physics research. Their encouragement and patience allowed me room to enjoy science, run into problems, and grow.

I would next like to thank my fellow graduate student labmates with whom I journeyed most closely through grad school and in science: Gyohei Nomura and Elie Bataille. The work that is contained in this thesis would not exist without them—from the first chamber bake to the first atoms, we have undertaken the vast majority of the journey as a team. I thank Gyohei particularly for teaching me to be methodical about the details that matter, and how to 'take it to a 9'. The MOT optics that he built allowed us to make a 3D MOT the same day that we activated the 2D MOT— which I think is a perfect example of his work ethic and forethought. I thank Elie for his technical prowess, for his energy and persistence that spills over, and for his willingness to bring others along (even when it comes to his 'French math'). His ability to always find a way around any challenge—experimental or theoretical—is astounding. I have benefited beyond words from Elie and Gyohei's scientific instincts, from their humor, from their humility, and from their friendship.

I would additionally like to acknowledge and thank others who have been a part of the team over the years, or who are just joining: Kon H. Leung, Xudong Lv, Nadine Meister, and Kanghuen Kim. I am incredibly grateful to Kon for his willingness to explain and teach technical details as well as life lessons, and for his significant contributions to the work in this thesis in his tenure as a postdoc. His expertise and thoroughness have been invaluable to the team. I am additionally grateful to have overlapped with Nadine who already knows how to ask good questions and find good answers. I look forward to the future of the experiment. *Zug zum Torr!*

The thanks extends beyond to the strontium lab. Thanks to Adam Shaw, Ivaylo Madjarov, Joonhee Choi, and Xin Xie for being the group that invited me in and taught me many early lessons of doing physics. Adam deserves special thanks and credit for recruiting me to the lab, and for his willingness to go on an emotional-support run when Elie and I broke the 2D MOT glass cell, among other contributions of friendship. I thank the next generation of strontium folks who each have taught me invaluable technical and pedagogical lessons: Pascal Scholl, Ran Finkelstein, Richard Tsai, and Xiangkai Sun. I thank additionally the other members of the Endres lab and collaborators with whom I have worked: Aadithya Shankar, Alice Jeffery, Andrew Winnicki, Andrew Ivanov, Dennis Naughton, Franz Pfanner, Jasen Zion, Lewis Picard, Phillippe Pearson, Ria Rosenauer, Shrishti Kulkarni, Su Direkci, Yuan Le, and Zunqi Li. Thanks especially to Shrishti for helping me grow as mentor, and to Richard and Ria for the runs and the hikes.

I extend thanks to those who have worked across Caltech facilities and in the physics and applied physics departments for their immense contributions to the lab running smoothly in aspects spanning from the administrative side to the HVAC, and for their responses in the various emergency situations that have arisen. I am particularly grateful to Loly Ekmekjian for her stamina and encouragement through friendship and to Alan Rice for his diligence and enthusiasm even in retirement.

I am also grateful to Chandler Conn, Phelan Yu, and Yuiki Takahashi for their AMO companionship over the years. I additionally would like to acknowledge our intramural softball team that brought few wins but much sweaty, froggie-cheer joy. I credit this endeavor with bringing the beginning of my friendship with Ana Duarte that I cherish.

The gift of friendship across the broader Caltech community was first cultivated by living in the Cats, rooming with Sarah Habib whose humor, thoughtfulness and cocoa contributed to the growth of many friendships and a uniquely special group.

This group includes: David Abramovitch, Isaac Legred, Ivy Bennet-Ford, Kiran Shila, Kyle Nelli (who also is to thank for his friendship since undergrad physics at Illinois), Nicholas Rui, Quinn Morgan, Rhiannon Udall, Sasha Alabugin, and Yoonsoo Kim. The combination of the intellect, humor, and kindness of this group brightened many days of grad school, and I have many words of thanks I wish I could write for each one.

I am grateful additionally for the Christian community that I found at Caltech and the friendships therein. The stability of this community, its encouragement in faith, and its sense of humor, was truly a gift. I would like to thank Alessandra Flaherty, Albert Wandui, Caroline Paules, Catherine Felce, George Daghljan, Grace Chen, Grigory Heaton, Luis Caldera, Maria Camarca, Nathan Wei, Patill Daghljan, Rachel Humann, Raymond Zhang, Rebekah Loving Ngo, Saharai Caldera, Sami Chang, Saren Daghljan, Tanner Harms, Tara Leininger, and Widi Moestopo. I thank Rachel for her kindness in reaching out to me early on, and for investing in friendship. I thank Maria for her humor, wisdom, and faithful companionship. I thank Caroline, Cat, George, Grigory, Patill, Saren, and Tara for their consistent commitment and encouragement through morning prayer (and the occasional theological throw-down).

I additionally thank my orchestra second-violin friends, for making music-making an endeavor of friendship. I am particularly indebted to Leah Stevenson, and Julie Cho in this regard.

I would like to express deep gratitude to the friends from college who have accompanied me throughout grad school through visits, calls, prayers, and laughter: Angela Kim, Carla Stelsel, and Emily Delacruz. You are each truly a gift.

I would like to thank my family from the bottom of my heart. I cannot imagine coming to the end of these years without their presence and without the support they have given. I thank them for seeing my love for science even when I found it growing blurry. I thank my sister, Melissa Manetsch, for being my best friend. Her passion, science, humor, and thoughtfulness each inspire me. I am incredibly grateful also to my parents— Scott and Cathy Manetsch who have both steadfastly supported me through all seasons of life. I thank them for their wisdom, prayers, love, and for picking up the phone to share in both the joys and sorrows. I thank my grandparents Tom and Lee McKee (Daddah and Bammy), and Tom, and Edna Manetsch (Grandpa and Grandma) for instilling within me and teaching through example, a love for science, for God, and for people.

At the foundation of every tree are its roots, and at the base of every building is a cornerstone. These are the source of both growth, and steadiness. I thank God for His presence, faithfulness, and goodness throughout graduate school. It is through the nourishment of these Roots and the stability of this Cornerstone that I stand at the end of this journey seeing both the joy and fruit of scientific endeavor, enabled by grace.

"...thus says the Lord God, 'Behold, I am the one who has laid as a foundation... a precious cornerstone, of a sure foundation: Whoever believes will not be in haste' "

Isaiah 28:16 (ESV)

ABSTRACT

Optical tweezer arrays of neutral atoms have been demonstrated in recent years to be a useful platform for quantum simulation, computation, and metrology. Realizing the full potential of these applications requires increasing qubit numbers while simultaneously maintaining high-fidelity control, which has proved a major challenge across quantum platforms. To date, quantum science experiments have typically been operated with at most hundreds of qubits.

In this thesis, we review the relevant physics of the optical tweezer array architecture and present the technical advances involved in scaling the platform to 6,100 atomic qubits. In tandem with scalability, we demonstrate robust vacuum design, state-of-the-art coherence times, and high single-qubit gate fidelities. To capitalize on the intrinsic connectivity of the architecture at large scale, we investigate long-distance coherent atom transport and briefly describe progress towards two-qubit gate implementation. Finally, we discuss the prospects for scaling the optical tweezer array platform further by an order of magnitude.

PUBLISHED CONTENT AND CONTRIBUTIONS

- [1] Hannah J. Manetsch, Gyohei Nomura, Elie Bataille, Xudong Lv, Kon H. Leung, and Manuel Endres. “A tweezer array with 6,100 highly coherent atomic qubits”. en. In: *Nature* 647.8088 (Nov. 2025), pp. 60–67. issn: 0028-0836, 1476-4687. DOI: 10.1038/s41586-025-09641-4.

H.J.M. and M.E. conceived the idea and experiment. H.J.M., G.N., E.B., K.H.L. and X.L. performed numerical simulations and contributed to the experimental set-up. H.J.M., G.N., E.B. and K.H.L. performed the experiments. H.J.M., G.N. and E.B. analysed the data. All authors contributed to the manuscript. M.E. supervised this project. Excerpts and figures from this paper (or adapted versions) are reproduced in the following thesis.

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Chapter 1

INTRODUCTION

This thesis introduces a new optical tweezer array experiment built using neutral cesium atoms trapped in far-detuned traps. The experimental apparatus was designed for scaling up qubit numbers in optical tweezer arrays by orders of magnitude, and has thus demonstrated the scalability of neutral atoms for quantum science. In this chapter, we begin with a broad overview of the field of quantum science (1.1) and then give a motivation for the necessity of scaling quantum systems for useful applications (1.2). We go on to provide an overview of pertinent work that has preceded this one across platforms in quantum computing and simulation (1.3), before narrowing in on advancements in neutral atoms in optical tweezers (1.4). Finally, we briefly introduce the experimental approach taken for trapping atoms (1.5), before outlining the contributions contained in this thesis (1.6), couched within the ever-exciting new developments that optical tweezer arrays of neutral atoms have presented.

1.1 The Landscape of Quantum Science

During the tenure of the work described in this thesis, the 100th anniversary of quantum mechanics was celebrated, commemorating the Heisenberg paper that turned the focus of the field towards observable quantities [1]. Rather than attempting to picture electrons like planets orbiting atoms with defined trajectories that could not be observed, this result pivoted to base the development of quantum theory on measurable quantities, such as energy levels in an atom. This turning-point led to the development of the matrix formulation of quantum mechanics, which has a natural starting point in the two-level system. With matrices, one can describe for instance, a spin-1/2 electron that can occupy either a spin up $|\uparrow\rangle = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$ or spin down $|\downarrow\rangle = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$ state.

The matrix formalism of quantum mechanics led subsequently to an understanding of the role of *superposition*, enabling a spin to be in multiple states at once, with a probability of being measured in each state. For a two-level system in state $|\psi\rangle$, this can be written as $|\psi\rangle = \alpha|\uparrow\rangle + \beta|\downarrow\rangle$, where the probability of the spin being measured in the $|\uparrow\rangle$ state is given by the normalized $p_{\uparrow} = \frac{\alpha^2}{\alpha^2 + \beta^2}$.

Extending this matrix-level understanding beyond a single spin occupying a particular Hilbert space results in a state that spans the tensor product of the Hilbert spaces occupied by each spin. A natural consequence of this fact is the possibility of *entanglement*, where the probability for measuring one spin in a particular state is dependent on the measurement outcome of another spin. For two spins labeled $\{1, 2\}$, one can write the state of the system $|\psi_{system}\rangle = \gamma|\uparrow\rangle_1 \otimes |\downarrow\rangle_2 + \delta|\downarrow\rangle_1 \otimes |\uparrow\rangle_2 + \epsilon|\uparrow\rangle_1 \otimes |\uparrow\rangle_2 + \zeta|\downarrow\rangle_1 \otimes |\downarrow\rangle_2$, where, depending on the values of γ, δ, ϵ and ζ , measuring spin 1 may change the probabilistic outcome of a measurement of spin 2. While we can write this state down for two spins quite easily, when the system size increases to N spins, the number of complex coefficients that we must use to describe the system scales as d^N (d being the number of states, where $d = 2$ for a two-level system), in contrast with the linear scaling $O(N)$ necessary to describe N classical particles.

This scaling with system size poses simultaneously a challenge and an opportunity. The challenge being that simulating quantum systems using classical computing resources quickly becomes impossible with growing numbers of particles. However, an opportunity can also arise, in that quantum particles inherently store information quantum mechanically, and thus, as Feynman [2] and others [3, 4] discussed in the 1980's, quantum particles themselves can be used to simulate quantum systems.

There are a few modalities that attempt to capitalize on this opportunity. One is quantum simulation, in which quantum systems can be simulated that have particular structures or dynamics that resemble the experimental platform being used. Another is *quantum computation*, in which quantum particles can be manipulated to implement gate-based computing with quantum properties. In this scenario, two states of a quantum particle can be chosen to represent a quantum bit or *qubit*, and universal quantum computing can be achieved to solve arbitrary quantum problems. Beyond simulating quantum systems, properties of quantum particles such as entanglement can be used to enhance sensing and metrology, and enable encrypted information transfer that is more secure. The goal of using collections of quantum particles to realize something that cannot be done using a classical computer: this is the proposition of *quantum advantage*.

1.2 The Mountain Range: Quantum Advantage

Scientists have identified many applications for which achieving quantum advantage would provide significant advancement, as laid out in the previous section and illus-

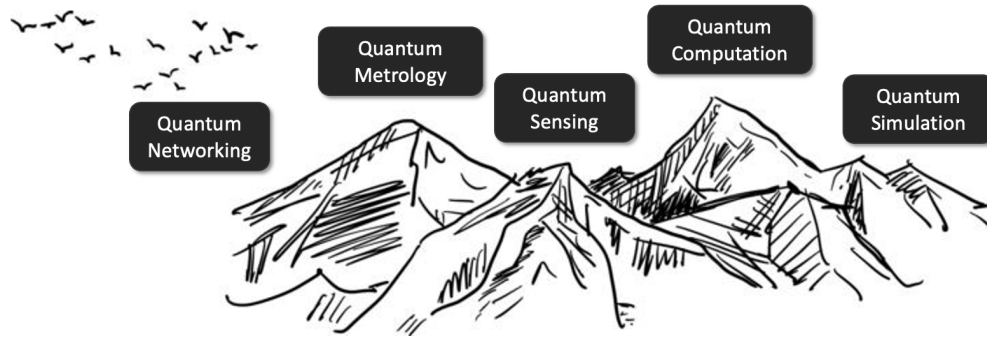


Figure 1.1: **Mountain Range for Quantum Advantage.** Within the landscape of quantum science, there's something of a mountain range for quantum advantage. There are various peaks or goals that scientists want to achieve, as well as many paths up these peaks, represented by various platforms for quantum science. In this thesis, we climb and bushwhack along the path of neutral atoms in optical tweezer arrays.

trated in Fig. 1.1. Experimentally reaching these goals, however, is a monumental challenge for each application— a peak as it were, that the community is attempting to summit. Various quantum platforms— including trapped ions, superconducting qubits, solid state qubits, and neutral atoms – have begun to carve out pathways towards these various peaks. However, across all of these platforms as well as for each of these summits, there is a similar challenge, which is to scale up the system size or qubit number, while simultaneously maintaining controllability and fidelity of the quantum system.

The necessity of scaling up system size becomes readily apparent as one looks more closely at each application. For instance, in quantum metrology, optical atomic clocks gain in stability and sensitivity as atom numbers increase [5]. This is true both in the standard quantum limit as well as in the Heisenberg-limited case. For quantum simulations at scale, increased system size can allow for probing boundary versus bulk effects and can enable parallel implementation of large-scale simulations. One could additionally use a large quantum simulator to study emergent collective phenomena [6, 7] or demonstrate verifiable quantum advantage [8, 9].

Finally, and perhaps most strikingly, for realistic applications of quantum computing, quantum error correction (QEC) is required to overcome the tension between increasing system size and required operational fidelity [10, 11]. Such tension may at first glance seem to be prohibitive to realizing useful computation, due to the fact

that the fidelity of preparing a quantum state drops exponentially with system size, and so maintaining the same fidelity at higher system size requires vast improvements in operational fidelity. However, encoding qubit information across many qubits (collectively dubbed a ‘logical qubit’) can allow errors to be identified and even corrected. This creates the possibility for overcoming limiting fidelities to implement computations at relevant scales for solving problems such as factoring. However, the cost comes in qubit overhead, and once again the answer is simple, but challenging: we need more qubits. In particular, QEC necessitates tens to hundreds of physical qubits per logical qubit even with encodings that capitalize on arbitrary connectivity [12, 13], and problems of interest require thousands of logical qubits. While there have been recent results that have demonstrated increasingly high encoding rates [14, 15] this still adds up to paint a picture in which scaling up system size of a quantum platform is paramount for reaching any of the summits.

1.3 Picking a path: Neutral atoms in optical tweezer arrays

Much recent work has been done to move forward the field of quantum science towards the goals discussed in the previous section, and several platforms have emerged as contenders to reach these peaks, particularly in the area of gate-based fault tolerant quantum computing. They each have their benefits and drawbacks, as we introduce here. However for each platform, system sizes and fidelities have been limited for quantum computing applications, and small demonstrations of QEC have not been utility scale. Thus, the era of Noisy Intermediate-Scale Quantum (NISQ) technology has developed [16].

Superconducting qubits are built on chip-based electrical circuits cooled to very cold temperatures in cryostats [17]. Transmon qubits, whose development in 1985 has received the Nobel Prize in 2025 [18], demonstrated the first two-qubit interactions with fidelity high enough to surpass the surface code threshold in 2014 [19] and have been a leading technology for quantum computing during the NISQ era. Recent results include demonstrating quantum error correction below the surface code threshold on a 105 qubit processor [20]. Superconducting qubits operate with significantly faster clock times than ion qubits or neutral atom qubits, but they suffer from deleterious effects when scaling. In particular, among other challenges with fabrication, qubits require radio-frequency connections that add heat to the cryostat and thereby limit the number of qubits that can be in one cryostat [21].

Single ions have been trapped using electric and magnetic fields since the 1980s [22],

and early on, trapped ions were proposed as a platform for implementing quantum computing [23]. Entangling gates were demonstrated with trapped ions [24] in 1995, and the platform has been a leading candidate for quantum computing ever since [25]. Long trap lifetimes and coherence times of ions have led to higher gate fidelities compared with other platforms and recent results demonstrate repeated error correction on 30 qubits with error rates that improve compared with physical error rates [26]. Ions however, are both slower than superconducting qubits, and present significant technical challenges for maintaining fidelities and controllability when scaling [25].

Single neutral atoms were trapped in optical tweezers (or tightly focused laser beams) for the first time in 2001 [27] providing the basis for the neutral atom quantum computing platform based on optical tweezer arrays. Proposals to utilize Rydberg excitations for entangling gates in neutral atoms were developed around the same time [28], and such entangling gates were first demonstrated in 2009 [29]. The controllability of the platform became clear with the development of arbitrary atom rearrangement in 2016 [30–32] and fidelities of two-qubit gates reached levels necessary for surface code thresholds only in the early 2020s [33, 34]. While the platform has only become a serious contender for realizing utility-scale quantum computing in the last few years, the progress has been rapid and encouraging, with recent results including operation of logical qubits with below threshold performance [35, 36]. Neutral atoms exhibit slower clock times than do superconducting qubits, limited by motion and readout. However, they have the advantage of the possibility for all-to-all connectivity [37]. Until the work presented in this thesis, it has not been clear whether scaling optical tweezer arrays of neutral atoms for quantum computing while maintaining fidelities would be possible.

This thesis will focus on the pathway being forged by neutral atom optical tweezer arrays and will present novel results for scaling this quantum platform, particularly geared towards summing the peaks of quantum metrology, quantum simulation, and quantum computation.

1.4 The neutral atom path: recent trailmarkers

Much work has already been done in blazing the pathways for neutral atoms towards the peaks of quantum metrology, quantum simulation, and quantum computation. In the following, we discuss a few relevant recent ‘trailmarkers’ for neutral atoms.

Towards quantum metrology, work has been done using optical tweezer arrays of

neutral atoms to combine the benefits of ion clocks that capitalize on single-particle control, and lattice clocks which benefit from high atom numbers [38–40].

Quantum simulation has been utilized by neutral atom optical tweezer arrays to explore a myriad of topics, wherein a Hamiltonian native to the platform is realized and mapped onto questions of interest [41]. The Hamiltonians that have been studied are constrained by the type of interactions that are implemented and the particular trapping architecture that is realized. In particular, one can implement Van-der Waals interactions [42] or dipole-dipole interactions [43], and trap using various tweezer geometries with different interaction distances, or even choose to superpose an optical lattice on top of the tweezers to enable hopping [44]. These systems have been used to study quantum phase transitions [45–48] as well as many-body quench dynamics [49, 50].

As already alluded to in the previous section, neutral atoms have in recent years developed quickly as a relevant platform for quantum computation. Thus, it has been more recently that gate sets for universal quantum computation have been realized on neutral atom quantum processors [51–54]. Steps towards QEC have been subsequently taken, preparing and executing computation using logical qubits [35, 36]. However, all of these scientific results were performed with 10s-100s of atoms. Recently there have been a few experiments that have reached around a thousand atoms trapped in optical tweezers via a non-contiguous array [55] or via continuous loading [56, 57], but they have done so without simultaneously demonstrating long coherence times and coherence-preserving transport. In parallel to this thesis work, there has also been development of continuous loading into about three thousand optical tweezers [58].

It is following this array of foundational work that we work to scale the platform further such that advancement can be made towards quantum advantage in each of these areas using optical tweezer arrays. We show here almost an order of magnitude scaling for the platform of neutral atoms up to 6,100 qubits trapped in 12,000 tweezers while maintaining high coherences and fidelities [59]. This represents the largest number of qubits across any platform, and paves the way for fault tolerant quantum computing using 10,000 neutral atoms.

1.5 Introduction to the experimental sequence

We will provide here a bird’s-eye view of the experimental sequence used to trap single neutral atoms for quantum science. This is additionally pictured in a schematic

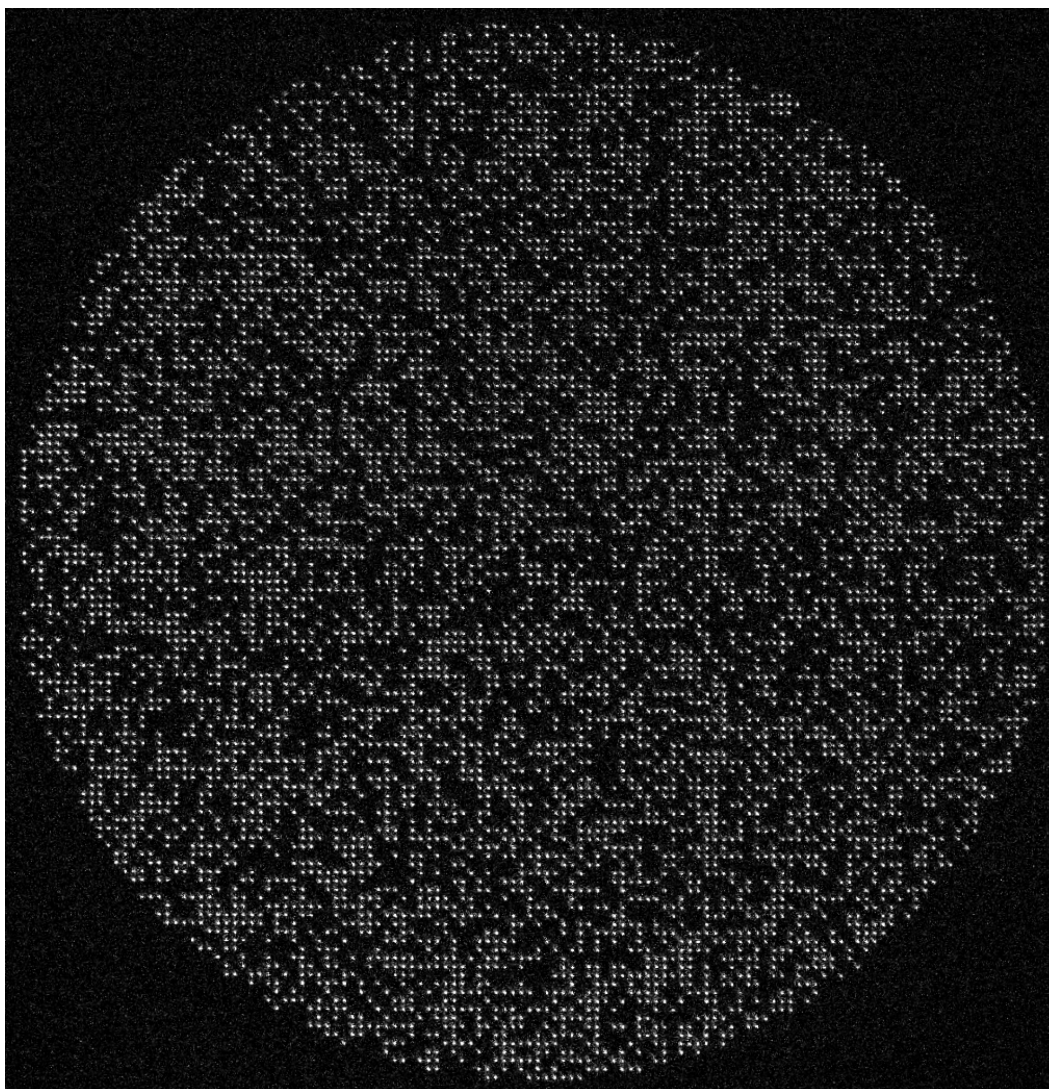


Figure 1.2: **6,100 atomic qubits.** Pictured is a single shot image of 6,100 atomic qubits in 12,000 optical tweezers. This thesis discusses the approach used to reach this order of magnitude improvement in scaling, and presents data demonstrating that these atomic qubits have state-of-the-art trap lifetimes, long coherence times, and record single-qubit gate fidelities. We additionally demonstrate long-distance coherent motion of qubits across the array, and preliminary two-qubit gate fidelities.

in Fig. 1.3 for visual clarity.

Neutral atom experiments for quantum science are performed in vacuum chambers under ultra-high vacuum (ultra-low pressure) conditions (See Fig. 2.37). The atom of interest must first be released from a source within the chamber before it can be trapped for experimental use. In our experiment, we use dispensers to

output cesium (Cs) atoms into our commercial source cell, where we make a 2D MOT (Magneto-Optical Trap) before pushing atoms through a pinhole into the main vacuum chamber (see Appendix A for details of chamber design). Atoms travel through this chamber to a glass cell, where they are trapped using magnetic and optical fields into a 3D MOT with about 10 million atoms at temperatures of around $100\mu K$. For experiments using directly-loaded single atoms in optical tweezers, laser beams (dubbed optical tweezers) are focused into this cloud of cold atoms, and the oscillating electric field of these beams induces dipole moments on surrounding atoms, attracting them to the center—the most intense part of the beam (see Ch.2.1.3 for more details). Once there, atoms in each tweezer undergo light assisted collisions, such that pairs of atoms are lost from the trap, parity projecting on the initial number of atoms in the trap to leave single atoms randomly filling $\sim 50\%$ of the optical tweezers. Subsequently, atoms can be located via imaging, and then science can be done – for instance, atoms can be rearranged, prepared in a particular state, or subjected to single or two-qubit gate operations. Finally, a second image can be taken in order to see where atoms remain. A given atomic state can be mapped to atom-presence or location in order to perform readout for quantum science experiments.

1.6 Outline

In this dissertation, we begin by discussing in more depth the trapping architecture for neutral atom quantum computing, and the choice of cesium for fulfilling the requirements for scalability while maintaining high fidelity and controllability (Chapter 2.1- 2.3). We then discuss large array tweezer generation (Chapter 2.4) and present the experimental platform scaled to 6,100 atoms and our measured high imaging fidelity, long imaging survival and long vacuum lifetime: metrics necessary on the level of controlling single atoms with high fidelity (Chapter 2.5- 2.7). We turn to implementing qubit operations on 6,100 atoms and present our experimental results of state-of-the-art coherence times and single-qubit gate fidelities at scale (Chapter 3). We go on to look more closely at the connectivity present in the neutral atom platform and present experimental data for the long distance transport necessary for a large-scale quantum computer, as well as discussion towards other zone-based operations including entanglement (Chapter 4). We conclude with a discussion of the considerations necessary for scaling beyond 6,100 qubits toward several tens or hundreds of thousands of qubits necessary for achieving a utility-scale machine (Chapter 5).

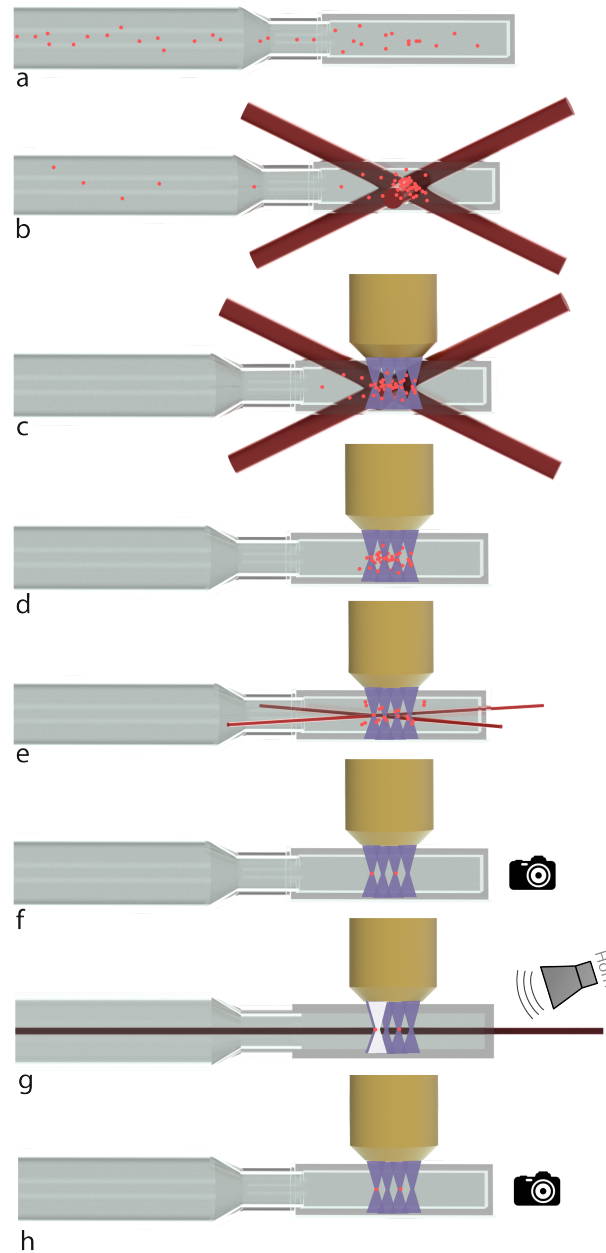


Figure 1.3: **Overview of the Experimental Sequence.** Atoms are first **a**, loaded into the vacuum chamber from the 2D MOT source cell and **b**, trapped in a 3D MOT using optical and magnetic fields (coils not pictured). **c,d**, Then tweezers are focused into the MOT cloud, and atoms are attracted to the highest intensity focal point. **e**, The atoms are then parity-projected and cooled using PGC (polarization gradient cooling) beams, and **f**, single atoms are left in the traps, which are imaged. After single atoms are located **g**, science can be done (pictured is a beam used for qubit state preparation, microwaves used for single-qubits gates, and another tweezer used for moving atoms), after which **h**, atoms are imaged again, mapping presence to an atomic state.

Chapter 2

TRAPPING ATOMS AT SCALE

As we come to the question of scaling up the platform of optical tweezer arrays, we first focus on the atom-light interactions that allow for atoms to be trapped with optical tweezers (2.1). We then discuss considerations for scaling this trapping architecture while simultaneously maintaining coherence (2.2), followed by our choice of cesium along with its prominent atomic transitions (2.3). We go on to discuss tweezer array generation techniques (2.4), and then finally our results for trapping a record number of 6100 atoms (2.5), image characterization with state-of-the-art high fidelity and survival (2.6) and record vacuum lifetimes allowing for atoms to be held for 23 minutes (2.7). These results in sections 2.6-2.7 represent important characterization to make the atoms useful for quantum computation or simulation, though they thus far represent characterization of classical properties of the atoms and the apparatus. Quantum properties of the qubits are examined in Chapter 3.

2.1 Optical Tweezers for Atom Trapping

In examining the experimental sequence as pictured in Fig. 1.3, it becomes quickly clear that there is a resource of millions of atoms in the Magneto-optical trap (MOT) that are not being used in typical optical tweezer experiments that trap only tens to hundreds of atoms. Thus, as we consider scaling up optical tweezer arrays, we must focus first on increasing the number of optical tweezers that have enough depth to trap an atom. In order to understand the use of optical tweezers as potential-wells in which atoms can be trapped, we look first at the basics of light-atom interactions that can be applied to this context. We begin by treating a two-level system, and then discuss how this is extended to a multi-level atom including hyperfine structure.

2.1.1 Light-Atom Interactions: The two-level system

The Hamiltonian for a two-level system with energy difference between the states of ω_{eg} can be written as

$$H_0 = \sum_n E_n |n\rangle\langle n| = \frac{1}{2} \hbar \omega_{eg} \hat{\sigma}_z \quad (2.1)$$

The Hamiltonian describing light-atom interaction in the dipole approximation can

be written as $H = H_0 + H_{int}$ with

$$H_{int} = -\vec{d} \cdot \vec{E}(t) \quad (2.2)$$

where $\vec{d} = e\vec{r}$ is the dipole operator describing the position and charge of the valence electron of the atom, and we can take $E(t) = E_0 \cos(\omega_L t) \hat{e}$ as the time-varying electric field of the light with frequency ω_L , polarization $\hat{e}$, and complex amplitude E_0 . The atom, via the dipole approximation, is assumed to exist within a time-varying field that is constant in space relative to the size of the atom. The interested reader is pointed to [60] for a derivation of this starting point.

By inserting a complete set of states, $\vec{d} = \sum_{n,m} |n\rangle \langle n| \vec{d} |m\rangle \langle m|$ and summing over an excited and ground state in a two-level system labeled as e and g , one can write

$$\vec{d} \cdot \vec{E}(t) = -\hbar\Omega \cos(\omega_L t) (|g\rangle \langle e| + |e\rangle \langle g|) \quad (2.3)$$

where we have defined $\hbar\Omega = E_0 \langle e | \vec{d} \cdot \hat{e} | g \rangle$ as the Rabi frequency for the two-level system. We can move into the interaction frame through a unitary transformation and then can apply the rotating wave approximation (RWA) where quickly oscillating components are dropped. Then the interaction Hamiltonian can then be written as

$$H = \frac{-\hbar\Omega}{2} (e^{i\Delta t} |g\rangle \langle e| + e^{-i\Delta t} |e\rangle \langle g|), \quad (2.4)$$

$$H = \frac{\hbar\Delta}{2} \hat{\sigma}_z - \frac{\hbar\Omega}{2} \hat{\sigma}_x. \quad (2.5)$$

The energy eigenvalues of this Hamiltonian are found to be

$$E_{+,-} = \pm \frac{\hbar\Omega_R}{2} \quad (2.6)$$

where $\Omega_R = \sqrt{\Omega^2 + \Delta^2}$ is the generalized Rabi frequency. In the case that the laser is far-detuned from the transition and $\Delta \gg \Omega$, one can Taylor expand Ω_R to find that the eigenvalues are

$$E_{\Delta \gg \Omega} = \pm \frac{\hbar\Delta}{2} \pm \frac{\hbar\Omega^2}{4|\Delta|} \quad (2.7)$$

which can be seen clearly to be a correction second order in perturbation theory to the two-level system which is not interacting with light. This energy shift is called the *AC Stark shift*, and gives rise to the trapping forces that we will use.

If we remind ourselves that that the Rabi frequency is related to the magnitude of the electric field and therefore the intensity of the light ($I = \frac{1}{2} c \epsilon_0 |E_0|^2$), we can see that the energy potential created by the light incident on an atom is

$$U(I) = \pm I \frac{|\langle e | \vec{d} \cdot \hat{e} | g \rangle|^2}{2c\epsilon_0 \hbar\Delta} \quad (2.8)$$

where c is the speed of light, and ϵ_0 is the vacuum permittivity.

To describe the dynamics of a closed level system under unitary evolution, we can choose to use the ‘density matrix’ $\rho = \sum_i p_i |\psi_i\rangle\langle\psi_i|$, which is a sum over the outer product of the states in the Hilbert space, weighted by their probability. The expectation value of any operator A can be calculated as $\langle A \rangle = \text{Tr}(\rho A)$.

One can derive dynamics by writing the Schrödinger equation in the von Neumann form

$$i\hbar \frac{d\rho}{dt} = [H, \rho]. \quad (2.9)$$

The two-level solutions under the Hamiltonian in equation 2.5 are given by the so-called *Optical Bloch equations*.

2.1.2 Light-Atom Interactions: Many-level system

Oftentimes, one transition can be viewed as dominant in an atom relative to an incident frequency of light, and so the two-level system is a useful conceptual grounding for the origin of the optical tweezer trapping potential, as discussed in the previous section. However, to calculate accurate energy potentials experienced by an atom, one must include all of the states that can play a role in the light-atom interaction.

In order to include the many states that will be relevant within an atom, we can return to equation 2.3, and instead of just summing over ground and excited states, we can sum over all relevant states.

$$\vec{d} \cdot \vec{E}(t) = -\frac{1}{2} \sum_{n,m} \langle n | \vec{d} \cdot \hat{\epsilon} | m \rangle E_0 (e^{i\omega_L t} + e^{-i\omega_L t}) |n\rangle \langle m| \quad (2.10)$$

In the far-detuned limit, the AC Stark shift is a small correction to the bare energy caused from mixing the state with other states. Thus, we can implement second-order time-dependent perturbation theory on this periodic potential. We write down the energy shift on state g from multiple states to be the following:

$$\Delta E_g = \frac{1}{4} \sum_e |\langle e | \vec{d} \cdot E_0 \hat{\epsilon} | g \rangle|^2 \left(\frac{1}{E_g - E_e + \hbar\omega_L} + \frac{1}{E_g - E_e - \hbar\omega_L} \right) \quad (2.11)$$

$$= -\frac{E_0^2}{2\hbar} \sum_e |\langle e | \vec{d} \cdot \hat{\epsilon} | g \rangle|^2 \frac{\omega_{eg}}{\omega_{eg}^2 - \omega_L^2} \quad (2.12)$$

We will now choose to work in the hyperfine basis because for atoms with non-zero nuclear spin in low-lying energy states— including for cesium, the main character of this thesis— F and m_F are ‘good’ quantum numbers.

What is a ‘good’ quantum number?

A quantum number is ‘good’ for a particular situation, when the eigenstates of the system being described are also eigenstates (or almost eigenstates) of the quantum number’s associated operator which is true when the associated operator commutes with the relevant Hamiltonian.

The total spin of the atom $\vec{S} = \sum_i \vec{s}_i$, when coupled to the total orbital angular momentum of the atom $\vec{L} = \sum_i \vec{l}_i$, gives a quantum number $\vec{J} = \vec{S} + \vec{L}$ which results in values ranging from $|S - L|$ to $|S + L|$. In an atom with a nuclear spin $\vec{I} = 0$, where spin-orbit coupling dominates, J would be a ‘good’ quantum number and the energy splitting is described by the *fine structure*. However in the case of an atom that has a nuclear spin of finite magnitude, one must define $\vec{F} = \vec{I} + \vec{J}$ as a ‘good’ quantum number with values ranging from $|I - J|$ to $|I + J|$, and the associated energy splitting is termed the *hyperfine structure*.

When atomic transitions are additionally shifted by magnetic fields of small magnitude compared with the fine or hyperfine splittings, magnetic Zeeman sublevels can be additionally labeled by m_J or m_F respectively, with values ranging from $-J$ to $+J$ or $-F$ to F .

The *Wigner-Eckart theorem* will allow us to rewrite the matrix element in the equation above, separating an angular part that depends on the light polarization and experimental quantization axis (or magnetic field orientation) from a radial part that depends only on wavefunction overlap. The theorem, applied to the transition $|F, m_F\rangle \rightarrow |F', m'_F\rangle$ can be stated:

$$\langle F' m'_F | d_q | F m_F \rangle = \langle F' || d || F \rangle \langle F' m'_F | F m_F; 1q \rangle \quad (2.13)$$

where d_q is the q^{th} spherical component of d with $q \in \{-1, 0, 1\}$, corresponding to $\{\sigma^-, \pi, \sigma^+\}$ polarization respectively relative to a $\hat{z}$ -directional quantization axis. We thus take a decomposition of $\vec{d}$ where $d_0 = d_z$ and $d_{\pm 1} = \mp \frac{1}{\sqrt{2}}(d_x \pm id_y)$. $\langle F' m'_F | F m_F; 1q \rangle$ is a Clebsch-Gordan coefficient which takes into account angular

momentum addition for the orbitals in question, including the impact of polarization. The selection rule of $m'_F = m_F + q$ arises from this coefficient. $\langle F' || d || F \rangle$ is the hyperfine *reduced dipole matrix element (RDME)*. We can choose to further reduce $\langle F' || d || F \rangle$, which is already independent of polarization and magnetic sublevels because the dipole operator only couples to the electron and not the nuclear spin. So we can write:

$$\langle F' || d || F \rangle = \langle J' || d || J \rangle (-1)^{F+J'+1+I} \sqrt{(2F+1)(2J'+1)} \begin{Bmatrix} J' & J & 1 \\ F & F' & I \end{Bmatrix} \quad (2.14)$$

where $\langle J' || d || J \rangle$ is the fine structure RDME that can be calculated from wavefunction overlap or measured. These values have been calculated and collected in this database [61] that is used for the relevant calculations included in this thesis.

The Wigner 6-j symbol above in brackets represents angular momentum combination for three angular momenta and can be related to a combination of Clebsch-Gordan coefficients as is explained in-depth in section 7.1.4 of reference [62].

In section 7.7 of the same reference, a detailed derivation can be found for the following total change in energy due to light incident on an atom, assuming that the external quantization axis is pointing in the $\hat{z}$ direction.

$$\Delta E = -\alpha^{(0)} |E_0|^2 - \alpha^{(1)} |E_0|^2 (i\hat{\epsilon}^* \times \hat{\epsilon})_z \frac{m_F}{F} - \alpha^{(2)} \frac{(3|E_{0z}|^2 - |E_0|^2)}{2} \left(\frac{3m_F^2 - F(F+1)}{F(2F-1)} \right) \quad (2.15)$$

This total energy shift on an atomic state from incident light is dubbed the *light shift*. Here, $\alpha^{(0)}$, $\alpha^{(1)}$, and $\alpha^{(2)}$ are functions of ω and F and are components of the real part of the polarizability tensor $\text{Re}(\alpha)$. The polarizability tensor α can be defined as $d = \alpha E$ where α relates the amplitude of the dipole moment d to the amplitude of the field E . $\alpha^{(0)}$, $\alpha^{(1)}$, and $\alpha^{(2)}$ are called the ‘static’, ‘vector’, and ‘tensor’ polarizabilities respectively and are given as

$$\alpha^{(0)} = \sum_{F'} \frac{2\omega_{F'F} |\langle F' || d || F \rangle|^2}{3\hbar(\omega_{F'F}^2 - \omega_L^2)} \quad (2.16)$$

$$\alpha^{(1)} = \sum_{F'} (-1)^{F+F'+1} \sqrt{\frac{6F(2F+1)}{F+1}} \begin{Bmatrix} 1 & 1 & 1 \\ F & F & F' \end{Bmatrix} \frac{\omega_L |\langle F' || d || F \rangle|^2}{\hbar(\omega_{F'F}^2 - \omega_L^2)} \quad (2.17)$$

$$\alpha^{(2)} = \sum_{F'} (-1)^{F+F'} \sqrt{\frac{40F(2F+1)(2F-1)}{3(F+1)(2F+3)}} \begin{Bmatrix} 1 & 1 & 2 \\ F & F & F' \end{Bmatrix} \frac{\omega_{F'F} |\langle F' || d || F \rangle|^2}{\hbar(\omega_{F'F}^2 - \omega_L^2)}. \quad (2.18)$$

Each term scales as we expected with respect to the frequency of the light and of the atomic transition in equation 2.12. We note that we do not apply the RWA here, as there are contributions from different states with different detunings relative to the laser, and for far-enough detuned transitions, the co-rotating and counter-rotating terms may both impact the contribution. The vector light shift can be seen to scale linearly with the m_F state, and as a result it is sometimes called a ‘fictitious magnetic field’. It is noteworthy that for linearly polarized light there is no vector light shift term. Further, for ground states of alkali atoms such as cesium (with ground state $6S_{1/2}$) the tensor light shift is also almost exactly zero. This can be seen clearly when the tensor polarizability is rewritten in the fine structure basis. Thus, to determine the trap potential for a ground state alkali atom trapped in linearly polarized optical tweezers, one must only consider the static light shift¹.

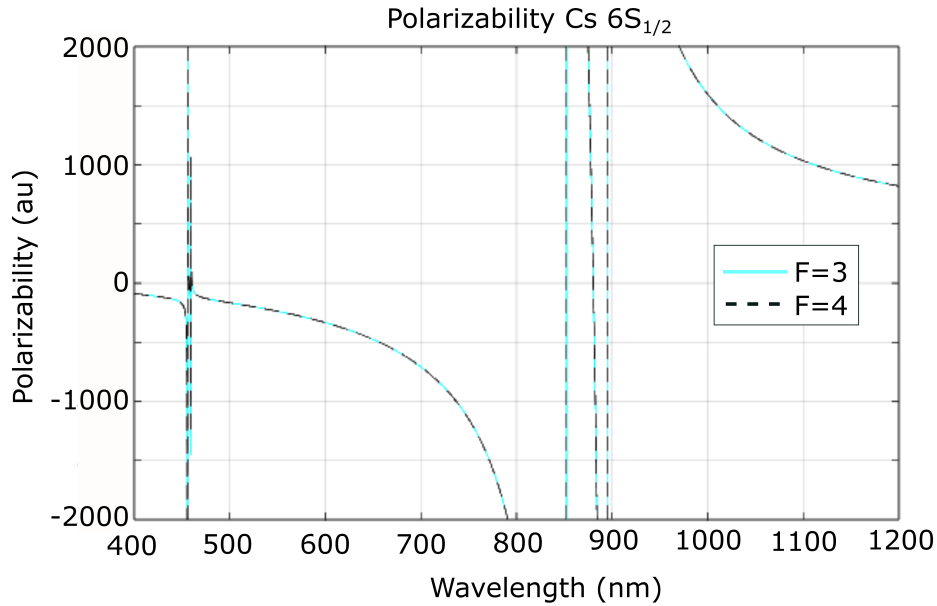


Figure 2.1: **Cs $6S_{1/2}$ Polarizability** The ground state dynamic polarizability for cesium hyperfine states is shown for light that is linearly polarized. The static light shift dominates the ground state polarizability. The vector light shift is exactly zero, and the tensor light shift is very small. All states in the NIST database are included in the calculation.

¹We note that linearly polarized optical tweezers exhibit regions of circular polarization (off-center) due to being strongly focused. This can have various experimental impacts. [63]

2.1.3 The Trapping Potential

As was true in the case of a two-level system (equation 2.8), the light shift above is proportional to the intensity of the incident light $|E_0|^2 \propto I$. Thus, the total energy shift or trap depth, can also be written in terms of the polarizability α as

$$\Delta E = U = \frac{-1}{2c\epsilon_0} \text{Re}(\alpha) I \quad (2.19)$$

where the real part of the polarizability describes the component of the dipole oscillation that is in-phase with the external field. This gives a restoring force, $F = -\vec{\nabla}U \propto -\vec{\nabla}I$, towards the highest intensity of the light when the detuning is negative.

A Gaussian beam has an intensity profile in space given by

$$I(r, z) = I_0 \left(\frac{w_0}{w(z)} \right)^2 e^{-\frac{2r^2}{w(z)^2}} \quad (2.20)$$

where w_0 is the waist radius, r is the distance from the beam center radially, z is the distance from the beam center axially, $w(z)$ is the beam $1/e^2$ waist at position z , and I_0 is the maximum beam intensity at the focus. The spatial derivative of this intensity profile thus gives rise to the restorative force that enables atom trapping with red-detuned optical tweezers. In order to trap and control single atoms rather than ensembles of atoms, these optical tweezers must be tightly focused [27].

Equation 2.19 makes clear that in order to increase the trap depth of a tweezer— or in fact, in order to make more tweezers of a given trap depth— we have two options. We can either increase the intensity that we are using linearly with the number of tweezers we would like to make, or we can trap using wavelengths that are closer to strong atomic transitions in order to increase the polarizability which scales inversely with the detuning.

2.1.4 Focusing light: Diffraction limited tweezers

In practice, high numerical-aperture (NA) objectives are used to focus laser light in a Gaussian (TEM₀₀ mode) into tightly focused optical tweezers. The NA of a focusing lens is defined as²

$$\text{NA} = n \sin \left(\arctan \left(\frac{D}{2f} \right) \right) \quad (2.21)$$

²While this is the standard definition of NA, note that objective manufacturers may use different definitions when quoting specifications. In particular the objective that was used in this work for generating tweezers used the following definition: $\text{NA} = n \frac{D}{2f}$.

where n is the index of refraction of the medium, D is the aperture size of the lens, and f is the effective focal length. This gives the angle between the principal plane and the maximal ray collected by the lens.

We solve for the field near the focal plane of the lens $E(x, y, z)$, given an input beam with field $E_0(x', y')$ in the back Fourier plane of the lens (often located at or inside of the back aperture of an objective). We start with the diffraction integral in the Fresnel approximation, derived in section 4.2 of [64] and propagate the field by a distance f to find the input field to the lens $E_1(x_l, y_l)$

$$E_1(x_l, y_l) = \frac{e^{ikf}}{i\lambda f} \iint E_0(x', y') e^{i\frac{k}{2f}((x_l-x')^2+(y_l-y')^2)} dx' dy' \quad (2.22)$$

We then multiply our input field by the lens function that gives a quadratic phase delay ($e^{-i\frac{k}{2f}(x_l^2+y_l^2)}$) and then again propagate the field a distance z using Fresnel diffraction.

$$E(x, y, z) = \frac{e^{ikz}}{i\lambda z} \iint E_1(x_l, y_l) e^{-i\frac{k}{2f}(x_l^2+y_l^2)} e^{i\frac{k}{2z}((x-x_l)^2+(y-y_l)^2)} dx_l dy_l \quad (2.23)$$

After integrating the outer integral over $dx_l dy_l$, we find:

$$E(x, y, z) = \frac{e^{ik(f+z)}}{i\lambda f} \iint E_0(x', y') e^{i\frac{k}{2f}(1-\frac{z}{f})(x'^2+y'^2)} e^{-i\frac{k}{f}(xx'+yy')} dx' dy'. \quad (2.24)$$

When $z = f$, the output field is just proportional to the Fourier transform of the input field, which is also the far-field Fraunhofer approximation for diffraction (section 4.3 of [64]).

$$E(x, y, f) = \frac{e^{i2kf}}{i\lambda f} \iint E_0(x', y') e^{-i\frac{k}{f}(xx'+yy')} dx' dy' \quad (2.25)$$

$$= \frac{e^{i2kf}}{i\lambda f} \mathcal{F}(E_0(x', y')) \quad (2.26)$$

For an input Gaussian beam, this is just another Gaussian beam with the intensity profile given in equation 2.20 with a waist radius given by

$$w_0 = \frac{\lambda f}{\pi w_{in}} \quad (2.27)$$

where w_{in} is the waist radius of the input beam.

The intensity of an AC electric field is defined in the following fashion

$$I = \frac{c\epsilon_0}{2}|E|^2, \quad (2.28)$$

although often we drop the prefactor and write this as $I = |E|^2$, when the relative intensity distribution is the relevant metric.

For our Gaussian beam here, the peak intensity I_0 of the focused beam is given as

$$I_0 = \frac{2P}{\pi w_0^2} \quad (2.29)$$

where P is the total power of the integrated beam, as derived clearly in chapter 3.1 of [65] along with other useful properties of Gaussian beams. This peak intensity then gives the depth of the optical trap $U_0 = U(I_0)$ when inserted into equation 2.8 which takes into account also the strength of the transition and detuning of the laser from the transition.

Not only does a small waist increase the peak intensity and therefore trap depth, minimizing the trap waist also impacts the trap frequencies that are determined in the harmonic trap approximation. Trap frequencies often determine the timescale for which operations can be performed without heating the atoms. This approximation is valid when the atom is cold compared to the trap depth. The radial (ω_r) and axial (ω_a) trap frequencies are given as:

$$\omega_r = \sqrt{\frac{4U_0}{mw_0^2}} \quad (2.30)$$

$$\omega_a = \sqrt{\frac{2U_0}{mz_R^2}} \quad (2.31)$$

where the trap depth is U_0 , and m is the atomic mass. z_R is the Rayleigh range of the beam, defined as $z_R = \frac{\pi w_0^2}{\lambda}$, and denotes the z -distance along the axial length of the Gaussian beam, where the waist radius reaches $w(z_R) = \sqrt{2}w_0$.

For an input beam whose diameter is large in comparison to the size of the input aperture for the lens or objective being used for generating tweezers, however, we must consider the impact of the input beam being clipped by this aperture. To do so, we can multiply the input field by an aperture function A in equation 2.26

$$E(x, y, f) = \frac{e^{i2kf}}{i\lambda f} \iint A e^{-\frac{x'^2+y'^2}{w_{in}^2}} e^{-\frac{ik}{f}(xx'+yy')} dx' dy' \quad (2.32)$$

where A has value unity when $x'^2 + y'^2 \leq (\frac{D}{2})^2$ and zero otherwise.

If the beam is so large that it can be approximated as uniformly filling the aperture, the result in the focal plane is the so-called ‘Airy disk’ intensity distribution, which appears Gaussian towards the center, with $w_0 \approx 1.22 \frac{\lambda f}{D}$, but has rings surrounding the center peak.

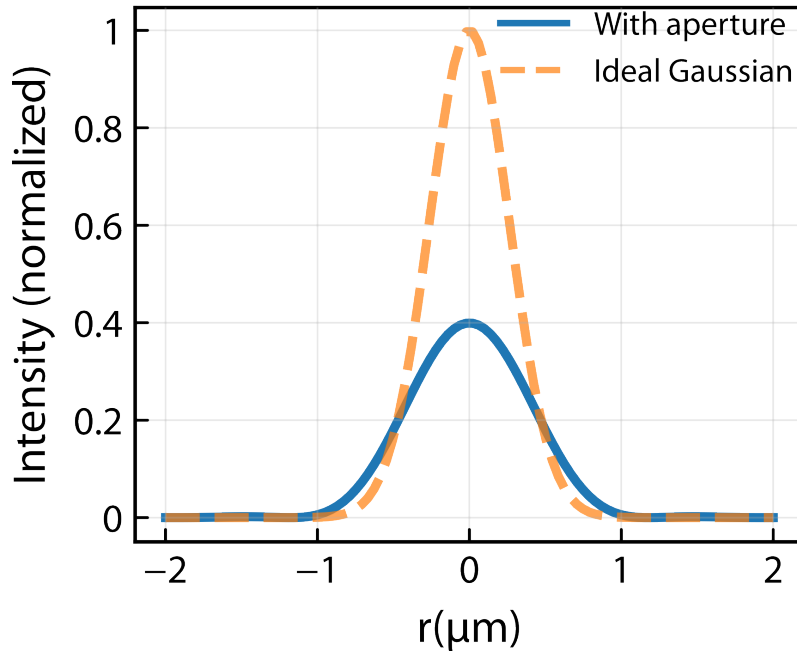


Figure 2.2: Gaussian Approximation vs Diameter-matched Clipped beam Pictured is the comparison of the intensity profile of a Gaussian beam that is focused by an objective for two scenarios. In blue, the objective aperture radius is matched to the waist of the beam and thus clips the beam, causing some power loss and additional diffraction. In orange is a beam with the same parameters that is not clipped with an aperture, and thus satisfies the Gaussian approximation. The intensities are normalized by the peak Gaussian beam intensity. For both cases, we choose a wavelength of $\lambda = 1060$ nm and an objective NA of 0.65 in accordance with our physical experimental parameters. For these parameters, the ideal Gaussian beam waist is 527 nm and the clipped waist (as determined by the waist at $1/e^2$ the peak intensity) is 759 nm.

In practice, we will work with beams that fill the back aperture of an objective such that $\frac{2w_{in}}{D} \sim 1$, which maximizes the trap frequencies and closely approaches the

maximum trap depth for a beam input through the aperture, as can be seen in [33]. Thus, this is as close as we can get to reaching *diffraction-limited performance*.

However, even for this optimal case, it is notable, that though the power loss from the $1/e^2$ tail being clipped on the objective is only about 14%, the diffraction pattern increases the beam waist, and lowers the trap depth to only about 40% of what one would expect in the Gaussian approximation as can be seen in Fig. 2.2. Making the beam much smaller causes the beam waist at the focal plane to grow significantly, and decreases both the trap depth and trap frequencies.

2.2 Scattering rates

While at first glance, it may appear beneficial for power efficiency to trap as close as possible to atomic resonances when scaling up trap numbers, an additional consideration arises. In bringing the trapping laser closer to resonance with a transition, one also increases the scattering that the atom experiences, as induced by the trapping light. Such increased scattering can heat the atom, or cause decoherence if the atom has been prepared in a particular qubit state. We will explore this consideration further in the following section.

2.2.1 Raman and Rayleigh Scattering

We now pivot to take a look at the imaginary part of the polarizability which describes the out-of-phase component of the dipole oscillation with respect to the field. One can obtain this part of the polarizability by no longer assuming that the states we are summing over have infinite lifetimes, but rather incorporating the decay rate by taking $E_e \rightarrow E_e - i\hbar\gamma_e/2$ in equation 2.12 for instance, where γ_e is the transition's natural linewidth [66]:

$$\gamma_e = \frac{1}{3\pi\epsilon_0\hbar c^3} \sum_g \omega_{eg}^3 |\langle e|\vec{d}|g\rangle|^2 \quad (2.33)$$

This out-of-phase component in the polarizability gives rise to absorption by the atom from the field, which is subsequently re-emitted. This may result in a two-photon process that returns the atom to its original state, which is called *Rayleigh scattering* or it may result in the atom being transferred into another state. If the atom has been prepared in a particular qubit ground state, this other state may be either the other qubit state, or another magnetic sublevel outside of the qubit subspace. This process is called *Raman scattering* [60].

The total scattering rate (Γ_{sc}) for an atom in an optical trap in state $|i\rangle$, including both Raman and Rayleigh scattering, is related to the imaginary part of the polarizability in the following way [67]:

$$\Gamma_{sc}^i = \frac{1}{\hbar\epsilon_0 c} \text{Im}(\alpha_i) I \quad (2.34)$$

where the polarizability depends on the frequency of the dipole trap.

In the rotating wave approximation and far from saturation, the Kramers-Heisenberg equation can describe the scattering rate of the process where the atom begins in state $|i\rangle$ and returns to state $|f\rangle$, after a two-photon scattering process through an intermediate state with hyperfine magnetic sublevel $e = \{F, m_F\}$ in k , where k represents an excited manifold with good quantum number J and $\Delta = \omega_L - \omega_{ei}$, is the difference between the laser frequency and the transition frequency between the initial state and the excited intermediate state [68, 69].

$$\Gamma_{sc}^{i \rightarrow f} = \frac{I}{2c\epsilon_0 \hbar^2} \frac{(\omega_L - \omega_{fi})^3}{3\pi\epsilon_0 \hbar c^3} \sum_{q \in \{-1, 0, 1\}} \left| \sum_k \sum_{e \in k} \frac{\langle f | \vec{d} \cdot \hat{\epsilon}_q^* | e \rangle \langle e | \vec{d} \cdot \hat{\epsilon}_L | i \rangle}{\Delta} \right|^2 \quad (2.35)$$

Here, to find the total scattering rate, we sum over the possible polarizations of the scattered light which has frequency $(\omega_L - \omega_{fi})$ where ω_{fi} is the frequency difference between the final and initial states, and we sum over all the states in the excited manifold and all excited manifolds.

To then find the total scattering out of the state $|i\rangle$, we sum over all possible final states $|f\rangle$ to find that the imaginary polarizability is

$$\text{Im}(\alpha_i) = \frac{1}{2\hbar} \sum_f \frac{(\omega_L - \omega_{fi})^3}{3\pi\epsilon_0 \hbar c^3} \sum_{q \in \{-1, 0, 1\}} \left| \sum_k \sum_{e \in k} \frac{\langle f | \vec{d} \cdot \hat{\epsilon}_q^* | e \rangle \langle e | \vec{d} \cdot \hat{\epsilon}_L | i \rangle}{\Delta} \right|^2 \quad (2.36)$$

If we sum only over the cases where $f = i$, we find the Rayleigh scattering rate and if we sum over $f \neq i$, we find the Raman scattering rate. It becomes apparent in implementing the sums that in the case of Raman scattering, there is some interference that amounts to scattering that scales as $\sim \frac{1}{\Delta^4}$, whereas in the case of Rayleigh scattering, the scaling remains as $\sim \frac{1}{\Delta^2}$. This scaling is clearly favorable for trapping, since we found earlier in equation 2.8 that the depth of the trapping potential scales as $\sim \frac{1}{\Delta}$, so the scattering falls off more quickly with detuning than does the trapping potential. However, we would still like to greatly minimize any coherence loss through our choice of trapping wavelength, and quantify the effect of

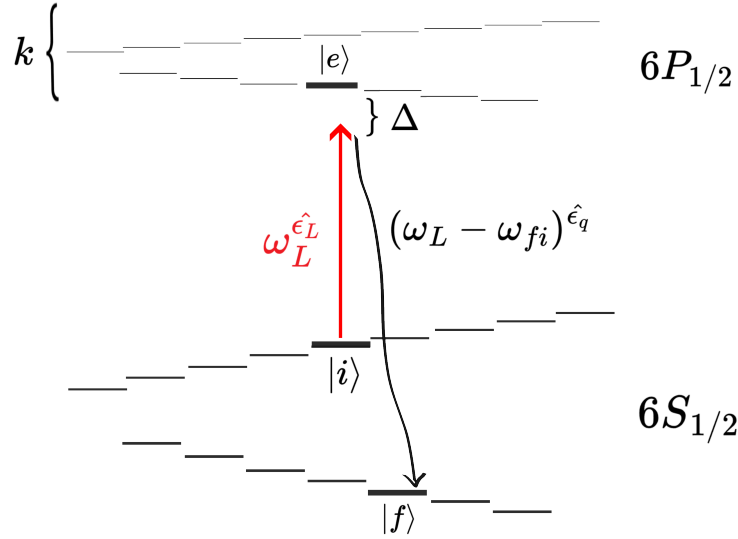


Figure 2.3: **Schematic depicting scattering on Cs D₁ line.** The approach to calculating scattering rates is described in equation 2.35, and we depict here the process of a laser with frequency ω_L and polarization $\hat{\epsilon}_L$ scattering off of excited state $|e\rangle \in k$. In this example we choose to depict scattering from the $F = 4, m_F = 0$ hyperfine ground state of Cs through $6P_{1/2}$. To find the total scattering rate of the particular state, one needs to sum over all possible excited states $|e\rangle \in k$ for all excited states k , and all possible final states $|f\rangle$.

trap-induced scattering. To do so, we can fix the trap depth and study the scattering and coherence as a function of wavelength. We rewrite the scattering rate from state $|i\rangle$ in equation 2.34, using equation 2.19, to be only a function of the laser frequency ω_L with constant trap depth U

$$\Gamma_{sc}^i(\omega_L) = -\frac{2 \operatorname{Im}(\alpha_i(\omega_L))}{\hbar \operatorname{Re}(\alpha_i(\omega_L))} U \quad (2.37)$$

and in Fig. 2.4 we plot scattering rates as a function of wavelength for Cs and Rb, summing over possible scattering off of all states in the D₁ and D₂ manifolds (see Fig. 2.7).

Typical experiments that use alkali atoms as qubits, encode the qubit states in the $m_F = 0$ magnetic sublevels in the hyperfine ground state since these levels are minimally sensitive to magnetic fields. Thus, we examine the scattering out of these states in order to understand the impact that scattering will have on our qubit coherence.

In Fig. 2.4, we plot Raman scattering for Cs (a,b,d) and Rb (c) at fixed trap depths

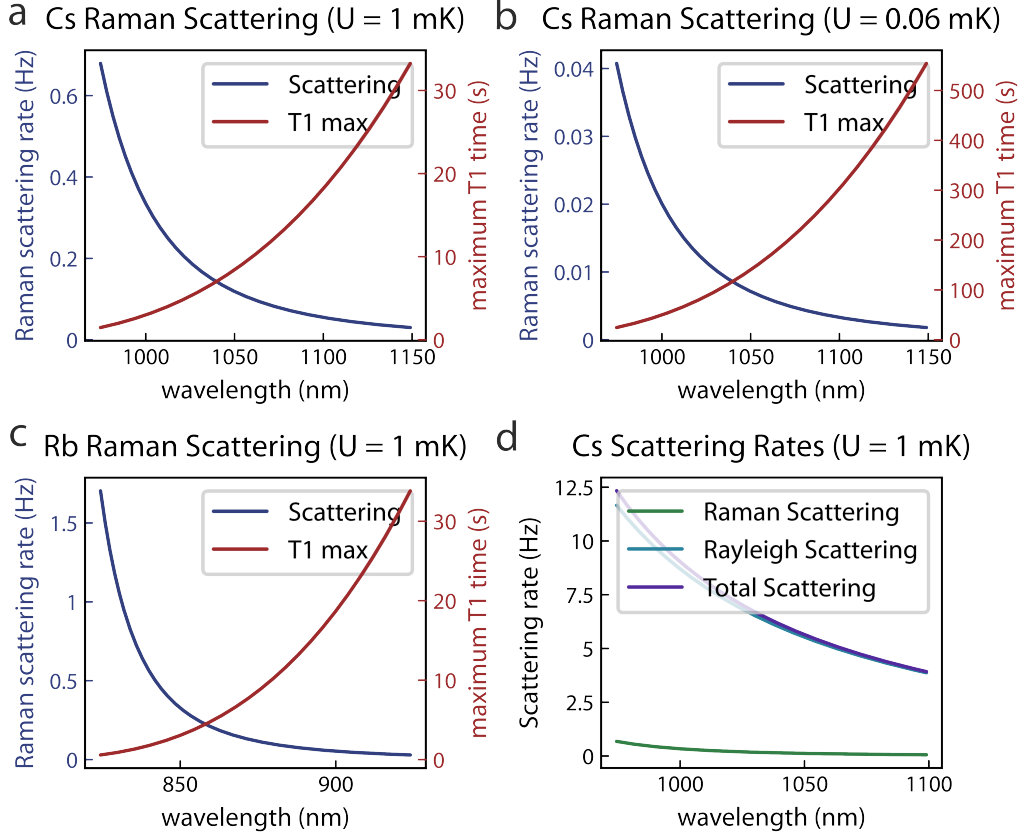


Figure 2.4: **Scattering rates from trapping potentials** We look more closely at Raman scattering rates as a function of wavelength at fixed trap depths. For each calculation we include only the hyperfine sublevels in the D_1 and D_2 lines in the calculation of $\text{Im}(\alpha)$, but we take all states into account in our calculation of $\text{Re}(\alpha)$. In particular, we look at scattering from $F = 4, m_F = 0$ in **a**, Cs at $U = 1$ mK trap depth, **b**, in Cs at $U = 0.06$ mK trap depth (the depth that we use during microwave operation in chapter 3), and **c**, from $F = 2, m_F = 0$ in Rb at $U = 1$ mK. We then also show in **d**, total scattering rates and Rayleigh scattering rates along with Raman scattering in the same conditions as in plot a, for a sense of scale.

out of the $F = 4, m_F = 0$ and $F = 2, m_F = 0$ states respectively. We show these calculations in Cs for both 1mK (Fig. 2.4a) as well as 0.06mK (Fig. 2.4b), given that in section 3.2 we will measure our T_1 time at 0.06mK. We look at ranges of wavelengths that are commonly used for trapping the respective atomic species. In Fig. 2.4d, we show clearly the dominance of Rayleigh scattering as a proportion of the total scattering rate which comes from the difference in scaling with Δ .

2.2.2 Impact on Coherence

Given that information can be lost from the qubit subspace via scattering, the von Neumann equation 2.9 can no longer be used to describe the dynamics. Open-system dynamics can be incorporated by tracing out the scattered photons, as described by the ‘Lindblad Master equation’:

$$\dot{\rho} = -\frac{i}{\hbar}[H, \rho] + \sum_n \left(C_n \rho C_n^\dagger - \frac{1}{2} \{C_n^\dagger C_n, \rho\} \right). \quad (2.38)$$

This describes dynamics $\dot{\rho} = -\frac{i}{\hbar}[H_{eff}, \rho]$ averaged over an ensemble of atoms evolving under the effective Hamiltonian

$$H_{eff} = H - i\frac{\hbar}{2} \sum_n C_n^\dagger C_n \quad (2.39)$$

and subject to jumps $C_{ij} = \sqrt{\Gamma_{ij}}|j\rangle\langle i|$ from state $i \rightarrow j$.

In the case of a two-level qubit system then, such jumps C_{00} or C_{11} represent Raman scattering and C_{01} or C_{10} represent instances of Rayleigh scattering. To characterize coherence times, we break up information loss into two categories: population loss (de-polarizing), and phase loss (de-phasing). The decay of population out of a qubit state $|0\rangle$ into $|1\rangle$ or into other non-qubit states via Raman scattering can be calculated by looking at the diagonal term ρ_{00} which contains the population of state $|0\rangle$. The $1/e$ time of this decay is called the T_1 time of the qubit. The phase incoherence between the qubit states on the other hand is contained in ρ_{01} , and the $1/e$ time of this decay is called the T_2 time of the qubit.

Evaluating the diagonal matrix element results in the following:

$$\dot{\rho}_{00} = -\Gamma_{Raman}^{(0)} \rho_{00} \quad (2.40)$$

where the solution $\rho_{00}(t) = \rho_{00}(0)e^{-\Gamma_{Raman}^{(0)}t}$ makes clear that $T_1 = 1/\Gamma_{Raman}^{(0)}$ as limited by this particular light-atom interaction. Evaluating the off-diagonal matrix element, we find

$$\dot{\rho}_{01} = -\left(\frac{1}{2T_1} + \frac{1}{T_\phi}\right)\rho_{01} \quad (2.41)$$

where now $\frac{1}{T_2} = \frac{1}{2T_1} + \frac{1}{T_\phi}$, where T_ϕ is dephasing due to the scattering rates from $|0\rangle$ and $|1\rangle$ being potentially different.

Raman scattering then clearly has a direct impact on the T_1 time of our qubits because it represents a direct de-polarization channel and full loss of information.

We then can also bound the maximum T_2 time, $T_2 \leq 2T_1$ as seen above. Given the scaling of Raman scattering being $1/\Delta^4$, one might simply assume that it will be negligible. However, since our primary motivation is to scale up the number of qubits while maintaining high fidelities, we should determine how closely one can detune to major atomic transitions while still maintaining this assumption. In figure 2.4, we look at scattering rates from Rubidium and cesium as a function of detuning for fixed trap depths. If we wish to not be limited to less than 10s [20s] of T_2 coherence time (5s [10s] of T_1 coherence time) at 1mK trap depth, we need to detune to at least ~ 860 nm [880 nm] for trapping Rb and ~ 1020 nm [1060 nm] for Cs (Fig. 2.4 a,c).

Rayleigh scattering, while contributing a significantly higher portion of the total scattering rate (seen in Fig. 2.4 d)– and certainly contributing to heating via momentum kicks from absorption and decay– returns the qubit to its original state. In doing so, and since the photons that are both absorbed and emitted are of frequency ω_L whether the qubit is in state $|0\rangle$ or $|1\rangle$, the quantum information is protected from decoherence [68].

Given these results, and our goal of maintaining long coherence times and high fidelities as we scale up the number of atoms, we choose to trap in far-detuned tweezers at wavelengths where high power is readily accessible, rather than trapping closer to a transition.

2.3 Cesium Transitions

In this section, we lay out the choice of atomic species that we have made for this project particularly aimed at scaling up qubit numbers in neutral atom quantum processors while maintaining fidelities. In particular, we have chosen to work with cesium (Cs) trapped in far-detuned optical tweezers.

Cs is an alkali atom with one valence electron in the first group of the periodic table, and has atomic number $Z=55$. Cs has only one stably occurring isotope: ^{133}Cs , which will be used in this work. Cs was notably used for the first two-qubit gate implemented using Rydberg atoms [29].

We describe first considerations of laser power availability for trapping and how this assists in the decision process. Then we go on to describe the transitions in Cs and how these are used throughout the remainder of this work.

2.3.1 Laser power availability

As a result of the considerations laid out in the previous sections, we now look at the available laser powers at trapping wavelengths that are far-detuned from the major internal transitions in atoms that have been trapped before in optical tweezers. In table 2.1, one can see for different atomic species and trapping wavelengths what the atomic polarizabilities and available powers are for the given wavelength. Cs, because of its heavier mass, has transitions that are closest to 1064 nm out of the stable alkali elements, providing therefore the highest ground state polarizability. At the time of the experimental design in this work, high power Yb-doped fiber amplifiers provided $\sim 120\text{W}$ of laser power from 1050 – 1080 nm, and by the writing of this thesis, the technology has been extended down to 1013 nm with 200W available, and the power at 1064 nm has been increased to 300W. Additionally, in the course of this work, the expected power at 532 nm has increased to have 80W available for continuous wave sources. Thus, the table reflects the currently available laser power at the time of writing. We thus estimate the number of 1mK deep traps we can make, assuming Gaussian tweezers with diffraction limited beam waists and the relationship between beam intensity and power as described in section 2.1. Note that this does not incorporate the loss of depth due to clipping (shown in Fig. 2.2), which reduces the depth to $\sim 40\%$ of the original.

Given the available laser powers, and the polarizabilities of available atoms at far-detuned trap wavelengths, we choose Cs for the purpose of scaling up the optical tweezer platform. Cs and the approximate wavelength that we use in this work for trapping is highlighted in gray in table 2.1, although the power that we use, as available at the time of experimental design, is 120 W. The assumption of 85% loss accounts for 80% AOM diffraction efficiency, $\sim 45\%$ SLM diffraction efficiency at high power, and then 85% transmission from the collective effect of other optics, and then 60% transmission through the objective. For the final column in the table, we assume trap depth reduction due solely to power loss of 85% though the effects caused by aperture clipping will further reduce the depth (shown in Fig. 2.2).

As discussed throughout the previous sections, the atomic level structure and detunings relative to the trapping wavelength along with angular factors determine the atomic polarizability. These factors, in addition to the natural lifetimes of states contribute also to scattering rates. Thus, we include in Fig. 2.5 a schematic of some relevant levels in the level structure of Cs along with their lifetimes. We additionally

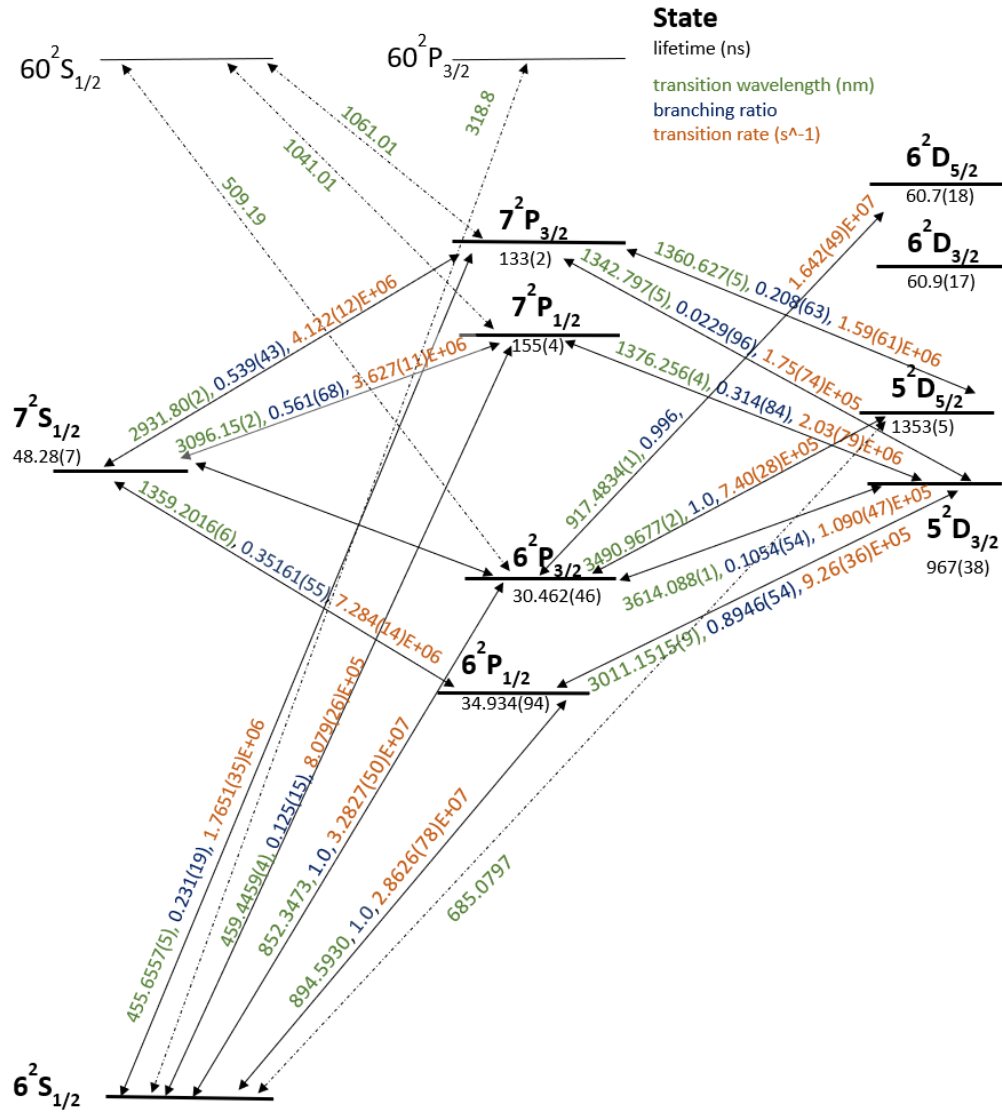


Figure 2.5: **Cs level structure.** Some lower-lying states in Cs that will be referenced throughout this thesis are depicted along with representative Rydberg states. Note that not all lines are included for all states. Notable are the 895 nm D_1 and 852 nm D_2 transitions from $6S_{1/2} \rightarrow 6P_{1/2}$ and $6S_{1/2} \rightarrow 6P_{3/2}$ respectively used for trapping, cooling, and single-qubit operations as described in this section. Additionally, we will use the two-photon Rydberg transition (455 nm and 1061 nm) through $7P_{3/2}$ for entanglement, as introduced in section 4.5. The 685 nm quadrupole transition $6S_{1/2} \rightarrow 5D_{5/2}$, and the 917 nm transition from $6P_{3/2} \rightarrow 6D_{5/2}$ will also be discussed alongside approaches for imaging and readout in the context of quantum computation (section 4.4).

Element	Wavelength (nm)	Laser Power (W)*	Polarizability (au)	Possible tweezer number (no loss)	Possible tweezer number (85% loss)
K	1064	300	599	94,834	14,225
Rb	1064	300	688	108,924	16,338
Cs	1064	300*	1168	184,918	27,738
Yb	532	80	258	43,569	6,535
Sr	813	25	286	6,462	969
K	850	30	1551	38,476	5,771
Rb	860	30	1879	45,535	6,830

Table 2.1: Atomic Species Ground State Polarizability and Trapping Wavelengths. We approximate the expected available laser power for various wavelengths, and the atomic polarizability at these wavelengths for various atomic species. Note that this table has been updated to reflect the laser power that is commercially available at the time of writing rather than at the time of experimental design. From these, we calculate a possible number of tweezers that can be created with 1mK trap depth, assuming an objective with 0.65 Numerical Aperture (NA) and an unclipped Gaussian beam waist beam input. We include the assumption of no loss (i.e. all the power transmits as tweezers), and then include also a category for loss, assuming that only 15% of the light reaches the atomic plane, after loss from an AOM, SLM, many mirrors and cubes, and an objective. Highlighted in gray is the choice of atom and wavelength that we have made for this thesis. *We note however that the power available at the time of building the apparatus described in this thesis was 120 W rather than 300 W, and this is the power utilized in the thesis. The polarizabilities here are taken from [61] and [70].

picture transition wavelengths, rates, and branching ratios which may be referenced throughout.

2.3.2 Cesium transitions for tweezer loading

Within the level structure of Cs, there are several prominent transitions that have been used consistently over the years for atomic cooling and manipulation. The two strongest transitions are the ‘D₁’ and ‘D₂’ lines, from $6^2S_{1/2} \rightarrow 6^2P_{1/2}$ and $6^2S_{1/2} \rightarrow 6^2P_{3/2}$ respectively. In this work, the D₂ line is used for the initial cooling stages of the MOT, and also for the subsequent molasses and polarization-gradient cooling (PGC) stages that cool and parity project atoms into tweezers. The D₁ line is used in optical pumping for qubit state preparation as well as for Raman transitions between the qubit states. For single-qubit gate control, however, we primarily directly utilize the 9.1926 GHz microwave transition. We picture the transitions that we use prominently in this work in Fig. 2.7, along with the Cs hyperfine structure parameters, as taken from [71].

We include a depiction of a typical experimental sequence in Fig. 2.6 and walk

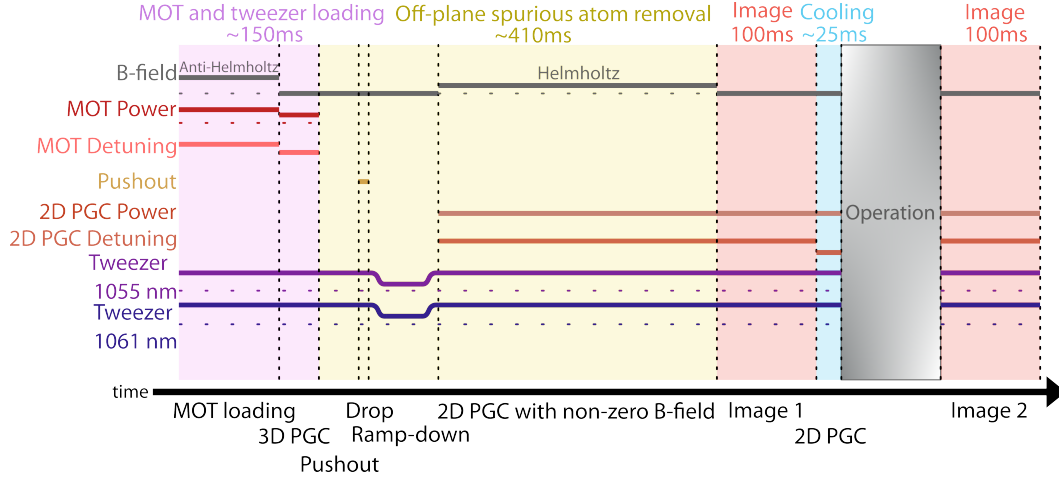


Figure 2.6: Experimental Sequence A typical experimental sequence is pictured, with increasing time from left to right, and relevant laser parameters indicated as on or off or ramped throughout. We typically load the MOT for 90 ms after which we implement 3D PGC cooling. We then implement various techniques to expel atoms unintentionally trapped in off-axis traps. Then we image for 100 ms, re-cool the atoms using 2D PGC, and do our science operations. We finally take another image, mapping atom loss onto the $F = 4, m_F = 0$ state that we push out. This figure is adapted from Ext. Data Fig. 1c in [59]

through in the following, the various stages of that process, highlighting also the relevant transitions in Fig. 2.7 and the lasers used to address them.

D₂ lasers

To address the 852 nm D₂ transition (from $6S_{1/2} \rightarrow 6P_{3/2}$), we use two phase-locked Distributed Bragg Reflector (DBR) lasers. We lock the repump directly to absorption spectroscopy from a Cs source cell on the $F = 3 \rightarrow F' = 3$ to $F = 3 \rightarrow F' = 4$ crossover of the D₂ line and then use AOMs to tune to the resonant $F = 3 \rightarrow F' = 4$ frequency for repumping. We offset lock the other laser on the $F = 4 \rightarrow F' = 5$ transition using a heterodyne beat detector referenced to a DDS (Direct Digital Synthesis) signal generator input that is referenced to a Rb clock. By switching the DDS, we can change the frequency detuning between the two lasers in $\sim 1 \mu\text{s}$.

D₁ lasers

We have two lasers that we use to address the 895 nm D₁ transition (from $6S_{1/2} \rightarrow 6P_{3/2}$). One is also a DBR laser which is offset locked (via a double-pass AOM)

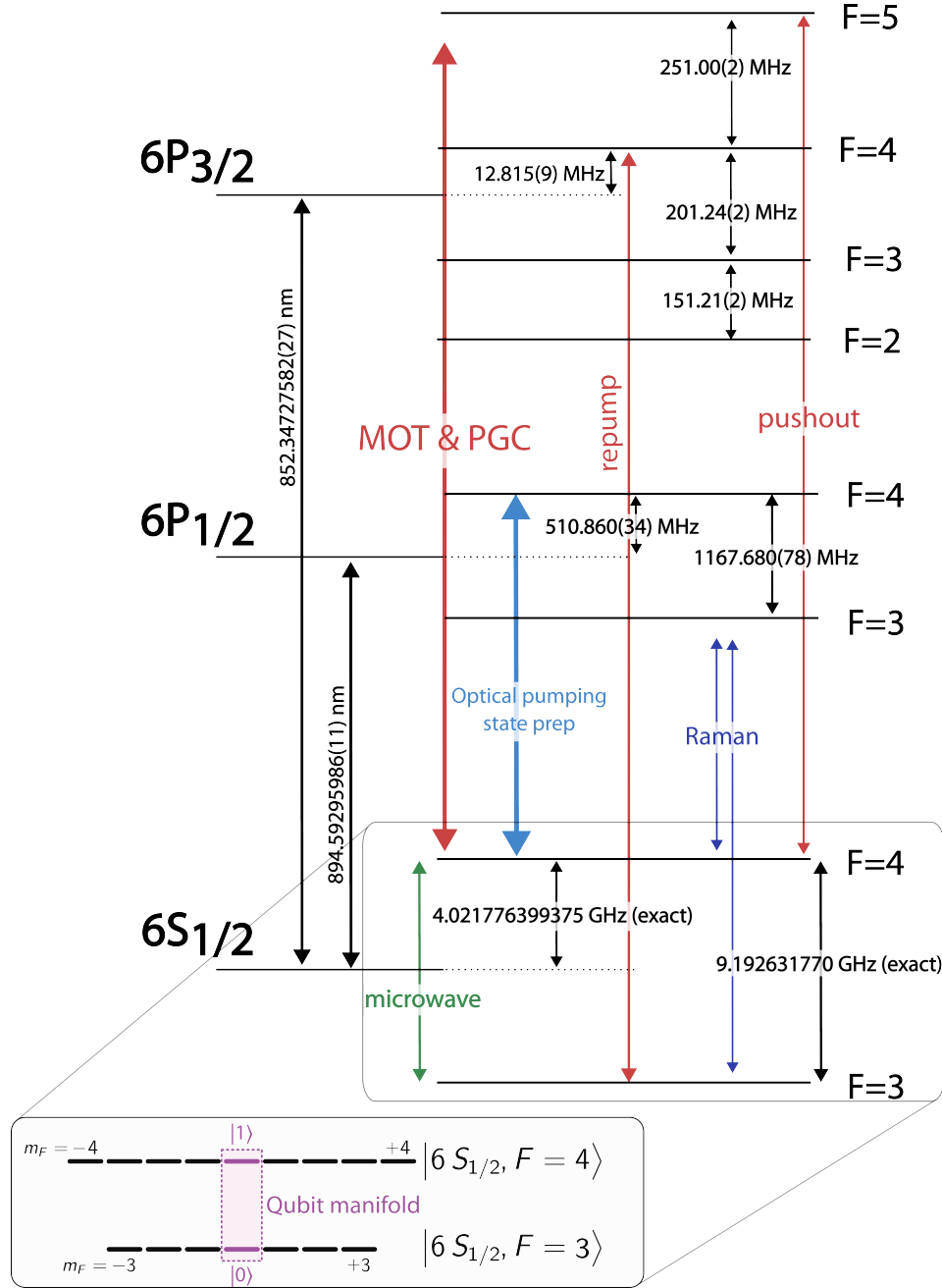


Figure 2.7: Cs hyperfine structure. We picture the Cs hyperfine structure for the ground state $6S_{1/2}$ and the excited states $6P_{1/2}$ and $6P_{3/2}$. The frequency spacings are labeled from [71] and we include in color the lasers that are used in this work. On the D₂ line (red), MOT&PGC beams are used for cooling and imaging the atoms respectively, alongside the repump. The pushout is used for readout mapped to loss from the $F = 4$ ground state. On the D₁ line (blue), optical pumping is used for state preparation of the qubits in $|F = 4, m_F = 0\rangle$ and Raman is used for single-qubit rotation. Primarily in this work, however, we implement single-qubit gates using microwave transitions (green).

to the $F = 4 \rightarrow F' = 4$ resonance via frequency modulation spectroscopy of a Cs vapor cell³. The laser is then brought back to resonance with another double-pass AOM for optical pumping state preparation. We additionally have another DBR laser which can be phase locked to the former using a similar configuration as on the D_2 line. Future schemes for utilizing these lasers for lambda grey molasses and enhanced loading are detailed in chapter 5.

We additionally have a Matisse TiSaph laser which is tuned to 895 nm and being used for generating Raman transitions on the D_1 line for single-qubit rotations, utilizing an amplitude-modulation approach for driving a two-photon transition with one beam [72].

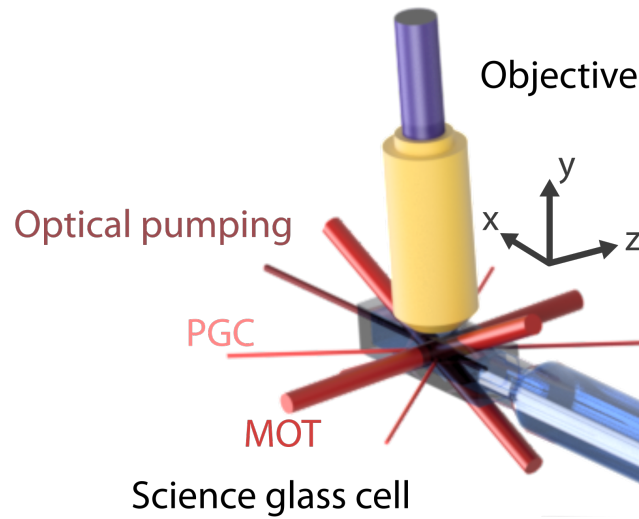


Figure 2.8: **Glass Cell.** We include a CAD rendering of our cooling and state-preparation beams at the glass cell. The 2.5 cm diameter retro-reflected MOT beams (and co-propagating repumping beams) are mutually orthogonal, with the anti-Helmholtz field axis being in the $\hat{z}$ direction. The 3.5 mm diameter 2D PGC beams (and co-propagating repumping beams) are located in the x-z plane. The optical pumping D_1 beam (and co-propagating D_2 repumping beam) used for state preparation is propagating along $\hat{x}$ and is linearly polarized.

MOT loading

Our atoms begin in a cesium source cell from Inflection which generates a 2D MOT with input light that has the same detuning as our 3D MOT, and are pushed through the vacuum chamber to the science cell with a push beam. Subsequently, atoms

³One cannot lock to a crossover on the D_1 line because the Doppler broadening isn't large enough to cover the D_1 hyperfine structure splitting.

are trapped in a 3D MOT using three mutually-orthogonal, retro-reflected beams pictured in Fig. 2.8. The beams are each 2.5cm in diameter, and are detuned from the $F = 4 \rightarrow F' = 5$ transition on the D₂ line by $\Delta = -3.17\gamma$, where γ is the linewidth of the $6P_{3/2}$ state ($\gamma_{6P_{3/2}} = 2\pi \times 5.2227(66)$ MHz). Resonant re-pumping beams ($F = 3 \rightarrow F' = 4$) of the same size are overlapped with the MOT beams. The total intensity of the MOT beams is $\sim 10I_{sat}$ and the intensity of the repumping beams is $\sim 1.6I_{sat}$, where $I_{sat} = 1.1\text{mW/cm}^2$ and is the saturation intensity of the stretched transition. We generate a magnetic field gradient for the MOT of around 20 G/cm using coils in an anti-Helmholtz configuration (see Appendix B for coil design), and typically load for 95ms. The external magnetic field from the earth or other lab equipment is also zeroed with offset fields created by coils in all three directions. The resulting MOT is roughly 1.6-mm $1/e^2$ diameter.

What is saturation intensity?

In addressing a state with finite lifetime, the Heisenberg uncertainty principle reminds us that

$$\Delta E \Delta t \approx \hbar. \quad (2.42)$$

The lifetime of the state $\tau = \Delta t$ comes also with an energy spread $\Delta E = \hbar \Delta \nu$, giving an inherent frequency linewidth $\Delta \nu = \frac{1}{2\pi\tau}$ or angular decay rate $\gamma = \frac{1}{\tau}$ which is the spontaneous emission rate given in equation 2.33.

The saturation intensity I_{sat} is reached when stimulated emission from the excitation laser is balanced with spontaneous emission. This is given:

$$I_{sat} = \frac{hc\gamma}{2\lambda\sigma_{eff}}, \quad (2.43)$$

where hc/λ is the energy of a scattered photon, and $\sigma_{eff} = \xi^2 \frac{3\lambda^2}{2\pi}$ is the scattering cross-section. If the atomic dipole is aligned with the field polarization, $\xi = 1$; otherwise the effective cross-section is weighted by the angular component of the dipole matrix element.

3D PGC

We then switch off the anti-Helmholtz field gradient which takes $\sim 8\text{ms}$ and simultaneously decrease the MOT beam intensity to $7I_{sat}$ and further detune to -19.5γ to

implement sub-Doppler cooling via polarization-gradient cooling (PGC). We apply this 3D PGC for about $\sim 40\text{ms}$, during which time atoms are loaded into our optical tweezers which are of typical depth $\sim 0.18\text{mK}$, parity projected into 0 or 1 atoms in each tweezer, and further cooled. Our detuning and power and offset fields during this stage are optimized on the atomic loading. After the tweezers have been loaded and cooled, we drop the 3D MOT/PGC beams.

Off-plane spurious atom removal

After dropping the MOT, the hope would be that we are just left with individual atoms trapped in the Gaussian tweezers that we are loading on the focal plane. However, given the way in which we generate our optical tweezers (see section 2.4.3), there are additional spots with low potential energy off the focal plane that can also trap atoms, due to the light for different tweezers interfering constructively depending on their phases which we do not control (see Fig. 2.19). These out-of-plane traps create a strong background that severely deteriorates imaging fidelity. Therefore, we go to much effort to remove these atoms. Before all of the out-of-plane heating techniques, we first implement on-plane cooling via 2D PGC as described in the following subsection.

To rid the array of out-of-plane atoms, we first send a resonant pushout pulse for $2\text{ }\mu\text{s}$, calibrated such that loading rate is not hindered to below 50%. Secondly, after re-cooling for 30ms using 2D PGC beams, we ramp down the tweezers quasi-adiabatically to one-fifth of the power, and wait for around 70ms , before which we again ramp up quasi-adiabatically.

Finally, we implement what we dub ‘bad PGC’, which is simply applying our 2D PGC beams with an added offset magnetic field of 0.19G , which applies a heating configuration rather than a cooling one. We do this for 180ms to obtain the imaging fidelities that are shown in section 2.6.2, but there is room for optimization of the field strength and detunings for this step in order to minimize the necessary time.

Why does a quasi-adiabatic ramp heat out-of-plane atoms?

Central to the reasoning for utilizing a quasi-adiabatic ramp for spurious atom removal, is the fact that the out-of-plane traps have lower trap depths than the on-plane traps. As a result, their trap frequencies are lower, since both radial (ω_r) and axial (ω_a) trap frequencies scale as $\sqrt{U}$.

The *Quantum Adiabatic Theorem* tells us that, given a particle in an eigenstate of an initial Hamiltonian, if the rate of change of the Hamiltonian is much smaller than the energy gap between the eigenstate and any nearby eigenstate, then the atom will remain in the corresponding instantaneous eigenstate throughout the evolution. This can be written:

$$\left| \frac{\langle p | \frac{\partial H}{\partial t} | n \rangle}{E_p - E_n} \right| \ll \frac{|E_p - E_n|}{\hbar} \quad (2.44)$$

where for a one-dimensional harmonic oscillator potential, $\langle p | \frac{\partial H}{\partial t} | n \rangle = m\omega(t)\dot{\omega}(t)\langle p | x^2 | n \rangle$, and $x = \sqrt{\frac{\hbar}{2m\omega}}(a + a^\dagger)$. The closest non-zero matrix element is when $p = n + 2$, and thus, this condition becomes

$$\left| \frac{\frac{\hbar\dot{\omega}}{2}\sqrt{n+1}\sqrt{n+2}}{2\hbar\omega} \right| \ll \frac{2\hbar\omega}{\hbar} \quad (2.45)$$

which simplifies to:

$$\left| \frac{\dot{\omega}}{\omega^2} \right| \ll \frac{8}{\sqrt{n+1}\sqrt{n+2}}. \quad (2.46)$$

Here now, it can clearly be seen that if we barely satisfy adiabaticity for the ω of our traps, then the off-axis traps which are more shallow and have smaller trap frequencies will not experience an adiabatic trajectory, thus heating up and having a higher likelihood of loss.

2D PGC

After having applied various heating to our atoms through the spurious atom removal, we re-apply cooling to the atoms in our tweezers. We now use two pairs of counter-propagating beams with a 3.5 mm $1/e^2$ diameter which are in the atom

plane as seen in Fig. 2.8. We use the same beams for imaging as well as for the ‘bad PGC’ described in the previous section. For optimal cooling, we utilize the detuning of $\Delta = -19.5\gamma$, as was used for 3D PGC.

We are able to achieve high imaging fidelity as well as high imaging survival in this configuration as detailed in section 2.6.3. While the 2D PGC cooling is only in the radial direction, we believe the axial direction of the atom is also being cooled as a result either of the plane of the 2D PGC beams having a slight angle relative to the plane of the atoms, or because of axial-radial temperature coupling due to the optical tweezers being imperfect Gaussian beams.

Imaging

To image our atoms, we use the same beams as used for 2D PGC. For our typical 80ms imaging, we lessen the detuning to $\Delta = -15.5\gamma$, as optimized on imaging survival. We also implement faster imaging in 20ms, where we detune even closer to $\Delta = -9.5\gamma$ in order to scatter more photons.

We collect photons for imaging on a qCMOS camera (Hamamatsu ORCA-Quest C1555020UP), which we choose for its fast readout time and its high resolution. We use the same objective to collect imaging light as we use to create the tweezers, which are combined on a dichroic mirror just above the objective (the alignment procedure for these objectives can be found in Appendix C). This allows thermalization of the objective to be shared as common-mode, as pictured in the schematic in Fig. 2.9. At our 852 nm imaging wavelength the objective has NA=0.55. After collecting only the photons within the cone half-angle determined by the NA, and due to the optical losses in the imaging system on the way to the camera, only around 2.7% of the photons scattered by the atoms enter the camera. The quantum efficiency of the camera at 852 nm is 44%. We focus the image onto the camera using an objective with $f = 40$ mm, resulting in a total magnification factor of the imaging system of 5.1. In section 2.6, we detail our single atom imaging results for our large atom array.

We first image prior to any quantum operation in order to locate where our atoms are present. We image a second time after all operations in order to map atomic presence/absence to the atomic state. We do this by using the resonant pushout beam as labeled in Fig. 2.7, to expel atoms that are in the $F = 4$ ground state, leaving only atoms remaining in $F = 3$.

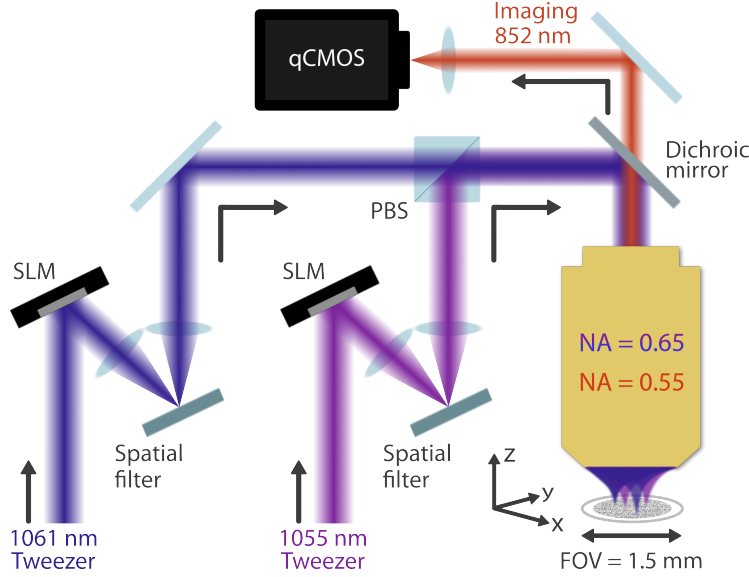


Figure 2.9: **Tweezer and Imaging pathways.** We depict the initial iteration of our combination of tweezer pathways with the imaging pathway on a dichroic. The objective both forms the optical tweezers from the SLM phase patterns which are imaged to the back focal plane, and also collects scattered photons from atoms when they are cycled on the imaging transition. These photons are then imaged onto a qCMOS camera and photon histograms are mapped to the presence or absence of an atom via thresholding (see section 2.6.1). This figure is reproduced from [59].

2.3.3 Cesium transitions for single-qubit operation

In the previous section, we described the transitions in Cs that are useful for preparing atoms into tweezers and imaging them. Now we proceed to discuss transitions that are used for quantum operation after atoms are loaded and imaged as pictured in Fig. 2.6. During single-qubit quantum operation, we set a small static magnetic field of 2.70G unless otherwise specified, in order to set a quantization axis and split the hyperfine Zeeman sublevels.

State Preparation

To begin with, the atomic states of our loaded atoms are evenly distributed amongst the Zeeman hyperfine sublevels of the $F = 4$ hyperfine ground state, as verified by microwave spectroscopy in our experiment. In order to implement quantum operations for simulation or computation, we must start with our atoms in a known state within the qubit subspace.

In order to prepare all the atoms into the qubit state $F = 4, m_F = 0$ or qubit state $|1\rangle$, we implement optical pumping with linearly polarized 895 nm light from $6S_{1/2}, F =$

$4 \rightarrow 6P_{1/2}, F = 4$ on the D_1 line⁴. This light is focused to $3.3\text{mm} \times 73\text{ }\mu\text{m}$ ($1/e^2$ waists) at the tweezer array and is incident on the atoms perpendicular to the quantization axis, along the direction of the cell, as can be seen in Fig. 2.8, and the polarization direction is aligned with the quantization axis. We simultaneously repump out of $|F = 3\rangle$ state using the repump on the $F = 3 \rightarrow F = 4'$ D_2 transition. Both beams are coaligned and linearly polarized using a Glan-Thompson prism. The transition $F = 4, m_F = 0 \rightarrow F' = 4, m'_F = 0$ is forbidden by selection rules because this transition would require an atom with no angular momentum to absorb a photon that has spin 1 without gaining momentum. From the Wigner-Eckhart theorem in 2.13, this results in the Clebsch-Gordan coefficient evaluating to zero. Since this transition is forbidden, and all other states are being pumped out of, population collects in the $F = 4, m_F = 0$ state after several spontaneous emissions. We are able to prepare atoms in this state with 99.2(1)% fidelity after 5ms of optical pumping, as inferred from the early-time contrast of the Rabi oscillations in Fig. 3.4. After preparing the atoms in $|1\rangle$, the trap depth is adiabatically lowered to $k_B \times 55\text{ }\mu\text{K}$ for microwave operation.

We measure a depumping rate of $0.064(5)\text{ }\mu\text{s}^{-1}$ from $F = 4$ to $F = 3$ at our operating D_1 optical pumping beam intensity when the D_2 repump is shuttered off. The intensity of the D_2 repump is then increased until the fidelity saturates. Factors that limit the state preparation include imperfect linear polarization purity, spatial variations in the pump laser intensity due to interference fringes arising from the surface of the science glass cell, and heating incurred during the optical pumping. Modeling our magnetic field coils, we estimate that the local direction of the bias magnetic field deviates by $< 10^{-5}$ radians for distances of $\sim 1\text{ mm}$ from the geometric center, and this has a negligible impact on the state preparation of our large scale array. Other state preparation schemes with higher fidelity have been demonstrated previously on smaller arrays and could be implemented in our system in the future [34, 73].

Microwave

In order to drive the $\sim 9.2\text{ GHz}$ qubit transition directly between the $m_F = 0$ magnetic sublevels of the hyperfine ground state, we use a linearly polarized microwave horn. The setup used to drive microwave transitions is pictured in Fig. 2.10. Similarly

⁴It is notable that using the D_2 line instead, even with perfect polarization would allow for off-resonant scattering via the $F = 4 \rightarrow F' = 5$ transitions, thus imperfectly preserving the population in $F = 4, m_F = 0$.

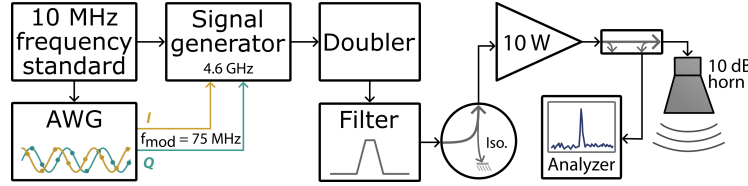


Figure 2.10: **Microwave setup** A schematic that represents the microwave generation setup. Signals are generated by the AWG, referenced to a 10 MHz standard and used to IQ modulate a signal generator. The signal is then doubled and filtered and passed through an isolator. After being amplified, the signal goes to the horn which is pointing at the atoms. A small amount of signal is returned at a bi-directional coupler to be analyzed on an oscilloscope.

to other experiments [74, 75] the RF signal from an arbitrary waveform generator (AWG, Spectrum Instrumentation M4i.6622-x8) IQ-modulates a microwave signal generator (Stanford Research Systems SG386) set at a fixed frequency of 4.6 GHz. The signal is then frequency-doubled, filtered, passed through an isolator before being amplified to 10 W of microwave power (Qubig QDA). A 10 dBi-gain pyramidal horn emits the microwave field on the atom array at a distance of about 15 cm from the cell, directed along the $-\hat{x} - \hat{y}$ direction in Fig. 2.8. We are able to reach a Rabi frequency on the atoms of 24.611(1) kHz using the microwave drive. This is the approach we use for all single-qubit gate operations in this work.

Raman transitions

Another option for driving the qubit transition is to use a two-photon optical process rather than a microwave drive. Since one can achieve higher Rabi frequencies using this method, gates can be implemented more quickly which can potentially be beneficial for increasing gate fidelities. There are several options for implementing such two-photon processes including via two phase-locked counter-propagating lasers with a frequency difference equal to the hyperfine spacing [76], but the approach we implement is based on a single beam driving a two-photon transition in a Doppler-free configuration via amplitude modulation on circularly polarized light [72]. This amplitude modulation is engineered by using phase modulation generated by a Qubig electro-optic modulator (EOM) at 9.2 GHz. A beam that is phase-modulated at the qubit frequency and 345 GHz detuned from the D_1 line is reflected twice off a highly dispersive chirped Bragg grating (Optigrade CBG-894-90), thus increasing the bandwidth. The input and output light to the grating then interferes, creating amplitude modulation. The RF input to the EOM is controlled

with the same AWG-generated, IQ-modulated, and doubled signal via a switch that we have between the microwave and Raman setup. Thus, in this work we have not implemented Raman and microwave driving simultaneously. We use the Raman setup to measure our trap frequencies as presented in Fig. 2.27.

How do selection rules differ for electric dipole (E1) and magnetic dipole (M1) transitions?

Selection rules determining which transitions can be allowed via a polarized oscillating field can be broken into two categories: angular momentum selection rules, and parity selection rules.

The electric dipole (E1) and magnetic dipole (M1) transitions rely on two different mechanisms for changing the state of the atom. The former operates based on the electric field interaction with the *charge* of the electron, and the latter operates based on the magnetic field interaction with the *spin* of the electron and the nucleus.

Both M1 and E1 transitions must satisfy angular momentum conservation with absorption of a spin-1 photon, so the angular momentum selection rules are shared for both cases: $\Delta F = 0, \pm 1, \Delta m_F = 0, \pm 1$

However, the magnetic dipole interaction Hamiltonian is given as $H_{M1} = -\vec{\mu} \cdot \vec{B}$ for the case of an M1 transition in place of the E1 Hamiltonian given by $H_{E1} = -\vec{d} \cdot \vec{E}$. The dipole operator has odd parity given that the operator $\vec{d} = e\vec{r}$, thus adding a requirement $\Delta L = \pm 1$. The magnetic dipole moment on the other hand has even parity and contains no spatial operators, thus requiring $\Delta L = 0$.

The ratio of the magnitude of the electric and magnetic dipole interaction energies is $\frac{U_{E1}}{U_{M1}} \approx \frac{1}{\alpha} \approx 137$. So the RF power needed to achieve a particular Rabi frequency on an M1 transition is orders of magnitude higher than the optical power needed for an E1 transition.

2.3.4 Cesium transitions for quantum circuit operation

Successful quantum circuit operation involves the confluence of many factors. Here, we discuss the Rydberg transitions that are used in the neutral atom quantum platform to implement two-qubit gates, as well as transitions that can be used for hiding and readout in Cs.

Rydberg transitions

Rydberg states are highly excited electronic states with large principal quantum number n . When an atom is excited to such a state, the valence electron can be far from the nucleus, causing the effective atomic radius to go from Angstrom scale to hundreds of nanometers. This allows for atoms that are held several microns apart to strongly interact, as explained in further detail in chapter 5.

In order to excite alkali atoms to a Rydberg state there are a few possible pathways, as illustrated in Fig. 2.5. To reach an $nS_{1/2}$ state, one must excite through an intermediate P state via a two-photon transition, typically with a large detuning to avoid scattering. In Cs typical two-photon transition pathways use $7P_{1/2}$ or $7P_{3/2}$ as an intermediate state.

One can also attempt to directly excite an atom to a Rydberg nP state via a single photon excitation. Because Cs is heavier than other alkali atoms, the wavelength range for this excitation is more favorable than in other cases: in the range of 318 – 320 nm.

The dipole matrix element between any particular state and a highly excited Rydberg state scales as $(n^*)^{-3/2}$ and thus the Rabi frequency scales as $(n^*)^{-3/2}$, where n^* is the effective principal quantum number. $n^* = n + \delta_l$ where δ_l is the quantum defect, accounting for the finite size of the nuclear core. Thus, for both a single photon transition from the ground state, as well as for the upper leg of a two-photon transition, higher power is required to maintain high Rabi frequency and overcome this scaling.

For this work, we utilize the $6S_{1/2} \rightarrow 7P_{3/2} \rightarrow nS_{1/2}$ excitation pathway with corresponding two-photon wavelengths of 455 nm and 1060 nm. We use two Ti:Sa lasers tuned to 910 nm and 1060 nm. We reference these wavelengths to a crossed cavity, and then we double 910 nm to 455 nm using an Agile Optics doubler. The crossed cavity allows some noise sources to be shared as common mode between the two arms, while allowing for different choices of Finesse for the different cavities which is useful because 910 nm is doubled to 455 nm. At 455 nm, we have about 2W of optical power after the doubling. For the upper Rydberg leg of 1060 nm, we input the Ti:Saph light to seed a Fiber Amplifier from Precilasers which can output 130 W. Both Rydberg beams are incident on SLMs for beam shaping and then sent to the atoms in a counter-propagating configuration, with the 455 nm beam being σ^+ polarized, and the 1060 nm beam being σ^- polarized. The approximate Rabi

frequency that is reached on the atoms for this two photon transition is 5 MHz, when using 2 GHz detuning, where the two-photon Rabi frequency is given as

$$\Omega_{2ph} = \frac{\Omega_{455}\Omega_{1060}}{2\Delta} \quad (2.47)$$

where Ω_{455} and Ω_{1060} are the single photon Rabi frequencies of the two legs of the two-photon excitation, and Δ is the detuning from the intermediate state during the excitation. In order to implement gates, the phase and amplitude of the 455 nm beam is modulated while the 1060 nm beam is held constant. See chapter 5 for further details.

Hiding transitions and 685 nm Quadrupole Transition

During quantum circuits, in order to implement mid-circuit measurement, some qubits must be measured and others must be kept coherent. This requires local or zone-based imaging beams, but it also requires that the data atoms that are not being imaged are protected from resonant scattered light that could cause decoherence. One way to do this is to ‘hide’ some atoms from the light by shining strong light that is resonant between $6P_{3/2}$ and a $D_{5/2}$ state, such that the imaging line experiences a strong shift out of resonance with the imaging and scattered light [58, 77]. A few of these such transitions in Cs are labeled in Fig. 2.5 and exist at 685 nm to the $7D_{5/2}$ transition (not pictured), 917 nm to the $6D_{5/2}$ and 3491 nm to the $5D_{5/2}$ transition.

Since such hiding light is used on-resonance in order to achieve the largest possible shift of the imaging state, the approximation used to calculate the AC Stark shift of a two-level system in equation 2.8 can no longer be used, which assumes $\Delta \gg \Omega$. In the opposite limit of $\Delta \ll \Omega$, and setting $\Delta = 0$ the energy eigenvalues of the interaction Hamiltonian result in the *Autler-Townes Splitting*, with energies:

$$E_{\Delta=0} = \pm \frac{\hbar\Omega}{2}. \quad (2.48)$$

Such hiding light could be particularly useful in combination with background-free imaging techniques that would enable atoms to be readout in-situ. Whereas typically cross-talk from attempting local imaging would cause decoherence to surrounding atoms or require large atomic spacings to avoid crosstalk, hiding light combined with a background-free imaging scheme can protect from this. Such a scheme in Cs uses the 685 nm narrowline quadrupole (E2) transition $6S_{1/2} \rightarrow 5D_{5/2}$, whose subsequent decay pathway is fully through the D_2 transition (branching ratio of 1)

in order to perform imaging without shining D_2 light directly on the atoms to obtain a cycling transition [78, 79]. A clear and detailed explanation regarding the strengths of quadrupole transitions with respect to beam geometry are discussed in chapter 3 of [80]. To optimally drive the $F = 4, m_F = 4 \rightarrow F = 6, m_F = 6$ transition requires that the beam propagation, polarization, and the quantization axis be mutually orthogonal.

2.4 Tweezer Array Generation

Scaling up the number of atoms that can be used as qubits requires not only choosing a favorable atom and trapping wavelength, but also requires utilizing optical engineering techniques and devices that can generate enough traps for single atoms to be loaded into. There are several approaches that have been demonstrated for engineering and loading optical tweezer arrays. Some combine optical lattices with tweezer arrays [56, 57, 81] and others directly generate optical tweezers using either active or passive devices that can apply phase or amplitude modulation to the light which is subsequently Fourier transformed into an engineered spatial distribution through a high-NA objective, as described in section 2.4.2. Some such active devices that have been used include Digital Micromirror devices (DMDs) [82, 83], Spatial light modulators (SLMs) [30, 32, 84–86] and Acousto-optical deflectors (AODs) [31, 87], and passive devices include Micromirror arrays (MMAs) [55], Metasurfaces [88, 89], and other diffractive optical elements (DOEs).

In this work, we use SLMs to generate $\sim 12,000$ static optical tweezers, and we use AODs to generate optical traps that can move on the timescale of quantum circuits for the sake of deterministic rearrangement as well as connectivity across the array. We utilize fiber amplifiers providing up to 130W at 1055 nm and 1060 nm to create the traps, combining light using polarizing beam splitting cubes, as pictured in Fig. 2.9. In Chapter 5, we discuss scaling further past 12,000 tweezers using combinations of the presented techniques in this thesis along with other active and passive components for tweezer array generation.

The tweezers generated by our SLMs are relayed with 4f telescopes to an objective that focuses them down to the diffraction limit. In this section we present considerations regarding scaling up numbers of optical tweezers related to SLM diffraction and phase retrieval. We also discuss aberration considerations for optical pathways involving large tweezer arrays.

2.4.1 High NA objective

In order to create a large array of optical tweezers for trapping atoms, we must first have access to a large area within a high-NA objective field of view at wavelengths that can be utilized for trapping. Such an objective must also be compatible with the thickness of the glass cell, with long enough working distance to focus into the vacuum. We additionally need to image atoms across the full field of view as well in order for the space to be useful for any array operations.

We thus iterated with Special Optics on a design for objectives with a field of view of 1.5 mm, that have $NA = 0.65$ at 1030-1080 nm for trapping and at $NA = 0.55$ ⁵ at 852 nm for imaging. The objectives are compensated for our 6.35 mm thick glass cell, and have a working distance of 12.5 mm vacuum + 6.35 mm glass + 1 mm air. The effective focal length of our objective is 7.97 mm, the back aperture size is 10.4 mm, and the housing diameter is 60mm. Details regarding objective alignment can be found in Appendix C. The housing for the objectives is made with Ultem.

We note that mounting the objectives is a particularly important experimental consideration, as very slight misalignment of the objective relative to the glass cell can cause significant aberration and decreased tweezer trap depth. In this aspect, we do not recommend our solution, though it has been satisfactory for our purposes. In our experiment we have a custom aluminum mount into which we screw the objective which is threaded. This is then mounted on a Newport 9081 5-axis stage (picomotor stages are rated for lower torque), with additional springs that have been added to compensate for the torque of the objective weight. For future experiments, we would recommend designing around a stage that is both motorized and can support the necessary torque required due to the size of the objective.

Having objectives that give us access to a large FOV is essential, but we also have to diffract light into tweezers across that area in order to utilize the space for trapping. We discuss techniques for accomplishing that large scale, controllable diffraction in the following section.

2.4.2 SLM Diffraction

For the work presented here, we use LCOS (Liquid Crystal on Silicon) SLMs for phase modulation in the Fourier plane of our atoms in order to generate tweezers. These devices use voltage-controlled crystals on each pixel whose index of refraction

⁵Note that this definition of NA is based on the tangent of the opening angle, being defined by Special Optics as $NA = \frac{D}{2f}$ where D is the aperture and f is the effective focal length of the objective.

can be changed with voltage in order to impart a variable phase-delay on the light that is reflected from that pixel. The phases applied to the SLM pixels are optimized in this work using the Weighted-Gershberg Saxton (WGS) algorithm in order to optimize for a target spatial distribution in the focal plane.

Several metrics are relevant for determining the usefulness of an SLM for generating a large tweezer array for a given trapping wavelength.

Reflectivity (or 0th-order diffraction efficiency) and light-induced damage threshold (LIDT) which may, depending on the wavelength, increase with water cooling capacity, are important especially for high power usage. These can be often related to the materials that are used for making the SLM, or whether coatings are used for individual pixels to increase reflectivity.

Modulation depth, or the range of phases that can be applied at each pixel can also decrease with thermalization, so having over 2π depth is useful, and is dependent on manufacturing. Increasing modulation depth can give access to various techniques for increasing diffraction efficiency or reducing pixel crosstalk [90, 91]

For scaling, pixel numbers are particularly important as we will discuss moving forward, and are closely related to the number of tweezers that we can make and to the 1st order diffraction efficiency as a function of deflection angle that an SLM can achieve.

Pixel Numbers and Resolvable Spots

In order to discuss the impact of pixel number on the number of tweezers that can be made using an SLM, we make a few assumptions. We assume that the spherically Gaussian beam incident on the SLM is mapped such that the waist of the beam is half the length of the shortest dimension of the SLM. We then additionally assume that the optical system thereafter maps this Fourier plane directly to a high-NA objective, such that the waist of the beam maps the radius of the aperture of the objective.

The field of view ($\text{FOV}_{x/y}$) of the SLM is determined by the largest input angle θ_{max} that is deflected by the SLM (as mapped to the back-aperture of the objective). This is given by the diffraction grating equation, derived from Huygen's principle which tells us that

$$\sin \theta = \frac{m\lambda}{\Lambda} \quad (2.49)$$

where Λ is the grating spacing. In order to find the maximum angle θ_{max} for the first diffraction order ($m = 1$), we must have the smallest possible grating— which, with an SLM is given by alternating pixels being bright or dark. Thus for this case, $\Lambda = 2Ma$, where a is the distance between the centers of adjacent pixels and M is the magnification between the SLM and the back aperture of the objective.

To find the FOV then, we use the small angle approximation to set $y_{max} = x_{max} = f \tan \theta_{max} \approx f \sin \theta_{max}$, thus finding:

$$x_{max} = \frac{f\lambda}{2Ma} = \frac{FOV_x}{2} \quad (2.50)$$

If we assume that the created tweezer is a Gaussian-diffracted beam, then the diffraction-limited beam radius, as introduced in equation 2.27 is given by $w_0 = \frac{2\lambda f}{\pi D}$, where D is the objective back aperture diameter, which in this scenario is $D = NMa$, where N is the number of pixels on the short axis of the SLM which has been mapped to the back aperture.

We then find that the linear number of resolvable spots s that the SLM can access in the image plane is given by

$$s_x = \frac{FOV_x}{2w_0} = \frac{\pi}{4}N. \quad (2.51)$$

For our SLM parameters, this results in $\frac{\pi}{4} \times 1200 = 942$ resolvable spots.

However, due to the beam clipping on the back aperture of the objective at its waist radius, the beam size is larger at the image plane, and thus the number of resolvable spots is fewer, as shown in Fig. 2.2.

For the particular case of our 1060 nm trapping light being clipped at its waist radius, we find a theoretical trap waist of $\omega_0 = 759$ nm by integrating over the field multiplied by an aperture function as described in equation 2.32. The theoretical linear FOV for the SLM in our experiment with $a = 8\mu m$ pixel pitch, magnified by $M = 10.2/9.6$ and an objective with effective focal length $f = 7.97$ mm, is $FOV_x = 0.994$ mm, thus resulting in 655 resolvable spots, as opposed to the prediction from the Gaussian limit.

Theoretical Diffraction efficiency

While we then have access to this full FOV using our SLM, we must also take into account the power usage required to cover this area, since the diffraction efficiency

of a particular tweezer in the first order is dependent on its deflection distance from the 0th order.

In order to find the diffraction efficiency as a function of tweezer position, we again work in the Fraunhofer Approximation as described in equation 2.26, where at the focal plane, ($z = f$), the field in the image plane is proportional to the Fourier transform of the input light field:

$$E(x, y, z = f) = \frac{e^{i2kf}}{i\lambda f} \mathcal{F}(E_0(x', y')). \quad (2.52)$$

We strongly reference [64] for the mechanics of the following derivation and the Fourier transforms of the relevant functions. In the case of the SLM, the input light field can be given by a sum over the field given by individual pixels. For a given pixel $\{m, n\}$ with SLM-applied phase $\phi_{m,n}$ the field in the plane of the objective back aperture is then:

$$E_{m,n}(x', y') = A_{mn} e^{i\phi_{m,n}} e^{-\frac{(ma)^2 + (na)^2}{w_{in}^2}} \text{Rect}\left(\frac{x' - ma}{p}\right) \text{Rect}\left(\frac{y' - na}{p}\right) \quad (2.53)$$

where p is the length of the active part of a pixel, determined relative to the distance a between pixel centers using the *fill fraction* ζ , where $p = a\sqrt{\zeta}$. A_{mn} is an aperture function that is unity when $(ma)^2 + (na)^2 \leq (\frac{D}{2})^2$ and zero otherwise. We make the approximation that the input Gaussian beam is uniform across individual pixels, defined by the position of the center of the pixel.

Now, we can sum over the pixels on the short edge of the SLM that has been mapped to the back aperture diameter, to find the total input field

$$E_0(x', y') = \sum_{m=-\frac{N}{2}}^{\frac{N}{2}} \sum_{n=-\frac{N}{2}}^{\frac{N}{2}} A_{mn} e^{i\phi_{m,n}} e^{-\frac{(ma)^2 + (na)^2}{w_{in}^2}} \text{Rect}\left(\frac{x' - ma}{p}\right) \text{Rect}\left(\frac{y' - na}{p}\right). \quad (2.54)$$

Taking the Fourier transform of this input field, we find the field at the focal plane to be

$$E(x, y) = \frac{e^{i2kf}}{i\lambda f} p^2 \text{sinc}\left(\frac{px}{\lambda f}\right) \text{sinc}\left(\frac{py}{\lambda f}\right) \sum_{m=-\frac{N}{2}}^{\frac{N}{2}} \sum_{n=-\frac{N}{2}}^{\frac{N}{2}} B_{mn} e^{i\phi_{m,n}} e^{-\frac{i2\pi a}{\lambda f}(mx+ny)} \quad (2.55)$$

where we have incorporated the amplitude of the beam and the aperture function into:

$$B_{m,n} = A_{m,n} e^{-\frac{(ma)^2 + (na)^2}{w_{in}^2}} \quad (2.56)$$

and where we apply the following Fourier transform found in section 4.1.1 of [64] :

$$\mathcal{F}\left\{\text{Rect}\left(\frac{x'}{p}\right)\text{Rect}\left(\frac{y'}{p}\right)\right\} = p^2 \text{sinc}\left(\frac{xp}{\lambda f}\right) \text{sinc}\left(\frac{yp}{\lambda f}\right). \quad (2.57)$$

Here we take $\text{sinc}(x) = \frac{\sin \pi x}{\pi x}$. The factor of $e^{-\frac{i2\pi a}{\lambda f}(mx+ny)}$ arises as a result of the *shift theorem* (derived in appendix A2 of [64]). The theorem states that if $\mathcal{F}\{g(x', y')\} = \mathcal{G}(f_x, f_y)$ for spatial frequencies f_x and f_y , then a shift in the input plane results in a linear phase shift in the Fourier plane:

$$\mathcal{F}\{g(x' - \xi, y' - \eta)\} = \mathcal{G}(f_x, f_y) e^{-i2\pi(f_x \xi + f_y \eta)}. \quad (2.58)$$

Substituting our optical spatial frequencies $f_x = \frac{x}{\lambda f}$ and $f_y = \frac{y}{\lambda f}$, and our physical pixel shifts $\xi = ma$ and $\eta = na$, yields the phase term in our sum.

The intensity of this beam is then given as

$$I(x, y) = \frac{p^4}{(\lambda f)^2} \text{sinc}^2\left(\frac{px}{\lambda f}\right) \text{sinc}^2\left(\frac{py}{\lambda f}\right) \left| \sum_{m=-\frac{N}{2}}^{\frac{N}{2}} \sum_{n=-\frac{N}{2}}^{\frac{N}{2}} B_{mn} e^{i\phi_{m,n}} e^{-\frac{i2\pi a}{\lambda f}(mx+ny)} \right|^2 \quad (2.59)$$

The phase profile of the SLM is sampled discretely, with pixels spaced by a . Accordingly, the summation in Eq. 2.59 is a discrete Fourier transform, and the focal-plane field is periodically repeated (these repetitions can be understood as Nyquist replicas, or Brillouin zones) with spacing $\lambda f/a$. The Nyquist criterion therefore limits the first Brillouin zone field of view (what is often called the ‘first order’) to approximately $\pm \lambda f/(2a)$ about the 0th diffracted order. The intensity envelope of the diffraction whose spatial distribution we optimize with the SLM is proportional to $\text{sinc}^2(\frac{px}{\lambda f}) \text{sinc}^2(\frac{py}{\lambda f})$, thus decreasing with distance from the 0th order position of $(x, y) = (0, 0)$.

We plot the relative diffraction efficiency across the SLM FOV in the focal plane, as given by the intensity envelope in Fig. 2.11 a. It becomes clear that though the accessible FOV of the SLM is rectangular, in order to maximize the diffraction efficiency of a tweezer pattern, the most efficient distribution of traps will be circular, centered on the 0th order of the SLM in the focal plane.

Not only do we wish to create a large number of tweezers across the FOV, for the purpose of atomic physics and particularly qubit operation, uniform trap depths are

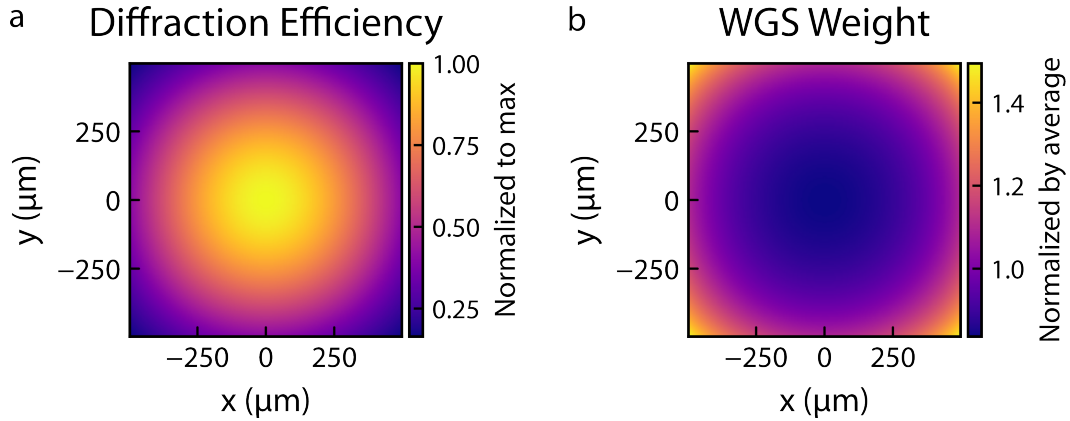


Figure 2.11: **Diffraction efficiency and WGS compensation** For the 2D FOV given by the SLM that is used in this work as calculated in subsection 2.4.2, we plot **a**, the diffraction efficiency given by the intensity envelope across the FOV, as derived in equation 2.59 and normalized to the maximum diffraction efficiency. In **b**, we plot the theoretically required weight W_i in the WGS algorithm we use with $G = 0.6$ in order to compensate for the decreased diffraction efficiency with distance from the center, as described in subsection 2.4.3.

required. Choosing cooling parameters that work well for a full array requires that the trap depths and trapping frequencies be uniform. Additionally, differences in trap depth contribute to differential light shifts between qubit states (See section 2.5.2 and Fig. 2.26). While the impact of these light shifts can often be minimized using dynamical decoupling, there are scenarios in which this impact will remain. Thus, we aim to solve the problem of creating a large array of uniform tweezers across our SLM FOV.

SLM 0th order filtering

Even if an SLM has very high diffraction efficiency into the 1st order, there will be some light that is not diffracted into this order. Though some of that light will be in higher orders, much of the light will be directly reflected into the 0th order, around which we are centering the tweezer array to maximize the diffraction efficiency. The amount of light in this order can be around 40% of the reflected light, so when using fiber amplifiers the total power can be tens of watts. Having this amount of power focused in the tweezer plane can be dangerous on the SLM optical path and create a large background even if defocused relative to the 1st order.

A typical approach used is to generate the tweezer array adjacent to the 0th order and then remove it using a D-mirror on the optical path, but this will not allow us to

create a large array, since we need to optimize the use of the whole SLM FOV. Thus, we utilize a technique for spatial filtering in the focal plane after the SLM tweezer generation. On the 1061 nm path, we glue two D-mirrors to an adjustable stage, and allow the 0th order light to pass through the slit between them (Fig. 2.12a and c). We fill in the gap between the tweezers that are reflected off of each half using tweezers created on a secondary pathway, also generated by an SLM.

On the 1055 nm path, we utilize a custom mirror with a 300 μm diameter hole, designed with Layertec, in order to allow the 0th order to pass through, while reflecting the 1st order. This mirror is pictured in Fig. 2.12b and d. In the tweezer pathway upgrade that we have implemented since [59], we have switched to using this mirror with a hole technique for both pathways since it works well and maximizes SLM tweezer usage.

2.4.3 Phase retrieval: WGS algorithm

We now arrive at the question of choosing the phases $\phi_{m,n}$ to be applied at each pixel of the SLM in order to realize a phase pattern that can generate the spatial distribution of $\sim 12,000$ optical tweezers for trapping atoms. We require the tweezers not only to be created, but also to be of equal depth such that differential shifts between tweezers are minimized. The spatial light modulator only has control over the phase distribution in the back Fourier plane, and the incident Gaussian beam is fixed. The intensity distribution in the focal plane is determined by both the input amplitude and phase, and thus we must utilize a *phase retrieval* algorithm in order to find phases that interfere in the necessary way to compensate for the fixed input amplitude. The fact that the phases in the focal plane are not fixed allows this to be a solvable problem [92].

There are several such algorithms that have been developed for various purposes [93, 94]. Conjugate gradient descent optimizes on a cost function in order to directly approach a solution [95] and Gerchberg-Saxton (GS) type algorithms apply iterative Fourier-transforms in order to optimize the phase [96].

In our experiment for tweezer generation, we utilize a Weighted Gerchberg-Saxton Algorithm (WGS) [97] that allows us to optimize the generated phase pattern to create traps with equal amplitude at specified locations. This algorithm has been used for creating optical tweezers in the atomic physics community for atom trapping [98–100] and has been studied and optimized across disciplines [101].

Recently, we have also pivoted to implementing an algorithm more similar to that

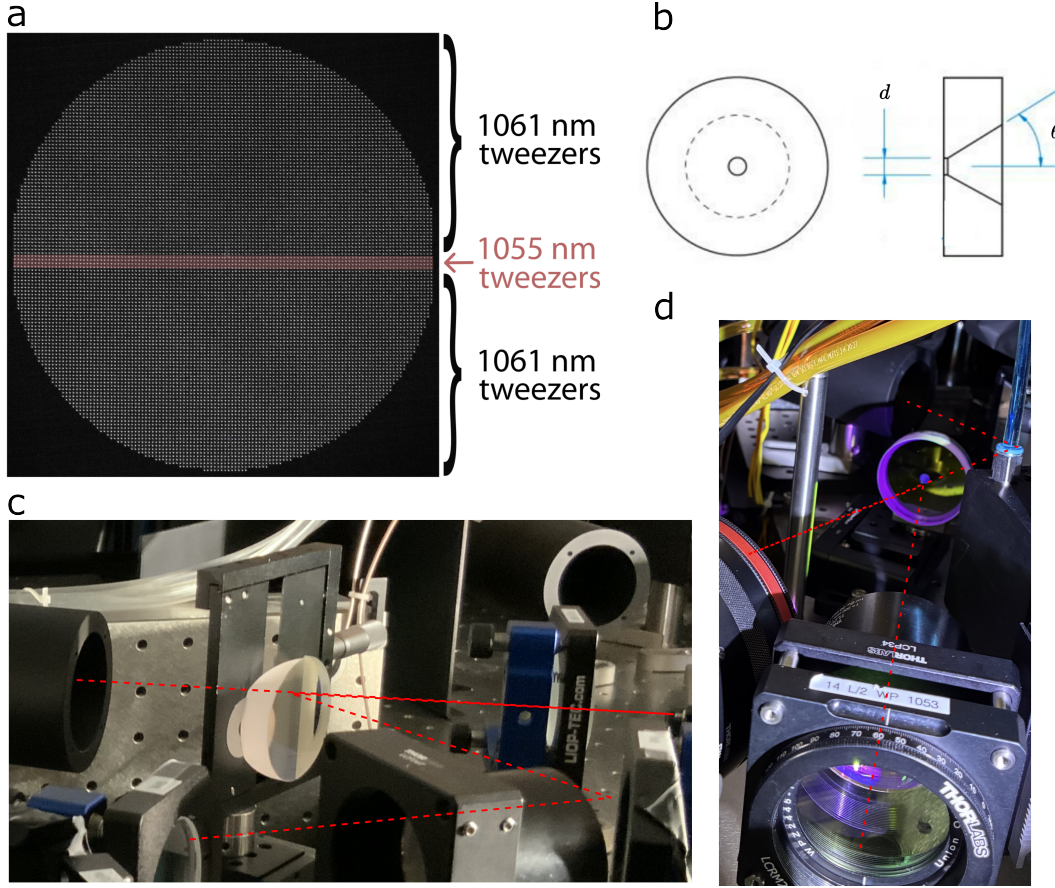


Figure 2.12: Spatial filtering SLM 0th order We utilize two different approaches for spatial filtering the SLM 0th order in the focal plane after our SLM phase pattern is focused through a lens. In particular we use a D-mirror based approach (a and c) and a mirror-with-a-hole based approach (b and d). **a**, Pictured are the tweezers created by two fiber amplifiers and labeled on the averaged atomic image. We create 11,513 (488) tweezers with laser light at 1061 nm (1055 nm), for a total of 12,001 tweezers. The 1055 nm tweezers fill the gap created by the spatial filtering of the 0th order in the 1061 nm tweezer pathway by the pair of D-mirrors. **b**, A schematic drawing of the mirror with a hole. In this approach (applied on the 1055 nm tweezer pathway), the array is aligned such that the 0th order light travels through the $d = 300 \mu\text{m}$ hole into a beam dump, and the 1st order is reflected such that the SLM light can be relayed to the objective. We utilize a mirror with a $\theta = 15^\circ$ opening. **c**, An image of the optical pathway using D-mirrors to filter the SLM's 0th order, where the two D-mirrors can be seen glued to a separation-adjustable mount and on a six-axis stage. **d**, An image of the optical pathway in which we use a mirror with a hole to filter the SLM's 0th order, mounted to a six-axis stage.

presented in [92]. This method requires a number of Fourier transforms that scales linearly with the number of tweezers being created and thus has a significant time cost. However, for the cases where we are only generating static tweezers, this

time cost may be worth the flexibility provided by local feedback (in particular we find local Zernike aberration correction very helpful). For significantly larger arrays (created with diffractive elements that have more pixels), the scaling of this technique may eventually become problematic.

We present below the technique that we used for generating the traps in [59]. Techniques described here such as phase fixing, diffraction correction, and the momentum-based feedback can also be ported over to other iterative feedback methods, such as those implemented more recently in our experiment for local control.

WGS Algorithm

The WGS algorithm proceeds computationally in the following manner, as illustrated in Figure 2.13. We assume as inputs a fixed Gaussian beam A_0 at the SLM plane, and a target amplitude $A_{target,i}$ that we wish to achieve in the focal plane (dependent, in our implementation on the loading-based uniformization cycle described in the next section and indexed by i). For the case of a target distribution of M optical tweezers, $A_{target,i}$ is an array of zeros with nonzero values at target coordinates $\{k\}$. This array may also be zero-padded [99] to improve the FFT result, a technique that we utilize in our implementation. We begin with a phase pattern ϕ_0 whose SLM pixel phases are randomly sampled from $[0, 2\pi]$.

To obtain the field in the focal plane, we apply an FFT to the field in the SLM plane. For the first iteration, this input field is given by $A_0 e^{i\phi_0}$ and thereafter, the phase is updated to ϕ_n from the $((n-1)^{th})$ iteration, while the amplitude is reset to A_0 . The FFT produces an optical field in the image focal plane of the form $A'_n e^{i\phi'_n}$ with an amplitude A'_n and phase component ϕ'_n . We can now look at the amplitudes at target tweezer locations ($A'_{n,k}$). We define the weighted mean amplitude at the tweezer locations as

$$\left\langle \frac{A'_{n,k}}{W_{i,k}} \right\rangle = \frac{1}{M} \sum_{k=1}^M \frac{A'_{n,k}}{W_{i,k}}, \quad (2.60)$$

where $W_{i,k}$ is the feedback weight associated with tweezer site k during the i^{th} experimental uniformization iteration, determined from an external metric such as atomic loading probability.

A recursive weight $g_{n,k}$ is updated to minimize the discrepancy between the realized and mean amplitudes:

$$g_{n,k} = \frac{W_{i,k}}{A'_{n,k}} \left\langle \frac{A'_{n,k}}{W_{i,k}} \right\rangle g_{n-1,k}, \quad (2.61)$$

where $g_{0,k}$ is a matrix of ones.

Finally, an inverse FFT is applied to bring the distribution back to the SLM plane: $\mathcal{F}^{-1}\{A_{target,i}g_n e^{i\phi'_n}\}$. This weighted update rather than just taking $\mathcal{F}^{-1}\{A_{target,i}e^{i\phi'_n}\}$ is what differentiates the WGS from the original GS algorithm [97].

The field amplitude is then replaced by the fixed Gaussian input beam, A_0 , and the cycle continues for N total iterations. As was demonstrated in [98], the weighting can have a stronger impact and allow for faster and more effective convergence if the focal plane phase ϕ'_n is fixed after some number of iterations. Intuitively, once a suitable focal-plane interference geometry has been established, continuing to update the phase in the focal plane can introduce fluctuations that redistribute tweezer power. For this work, we thus fix ϕ'_n as constant after around $n > N/2$ iterations, where $N \sim 50$, allowing the weighted amplitudes to continue to evolve to balance tweezer intensities. In figures such as 2.16, this phase fixing marks a significant drop in the error signal. However, if applied too early, the error signal will not converge.

Loading-based uniformization

After the WGS algorithm has converged numerically, additional experimental feedback can be applied using either measured tweezer intensities or an atomic signal in order to compensate for imperfections not captured in the theoretical model. For example, there may be imperfect 4f telescope matching, or the beam may not be perfectly centered on the SLM with the exactly expected waist. In our experiment, building on previous work [100], we feed back directly based on our tweezer loading rate. This only works if the atoms are loading into tweezers to begin with, so sufficient alignment and tweezer fitting (section 2.6.1) is required. Additionally, the initial weights can be chosen to compensate for the diffraction efficiency of the SLM, ensuring nonzero loading across the array at the start of the feedback procedure. In order to converge to a sufficiently uniform array, we typically only require about 5 iterations of feedback.

To implement this feedback, we choose a site-dependent weight $W_{i,k}$ to affect the feedback in equation 2.61 which incorporates compensation for the loading probability, as described in equations 2.62 and 2.64.

The weight matrix for the i^{th} iteration at the k^{th} tweezer site is chosen as:

$$W_{i,k} = 1 - G(1 - \sqrt{H_{i,k}}). \quad (2.62)$$

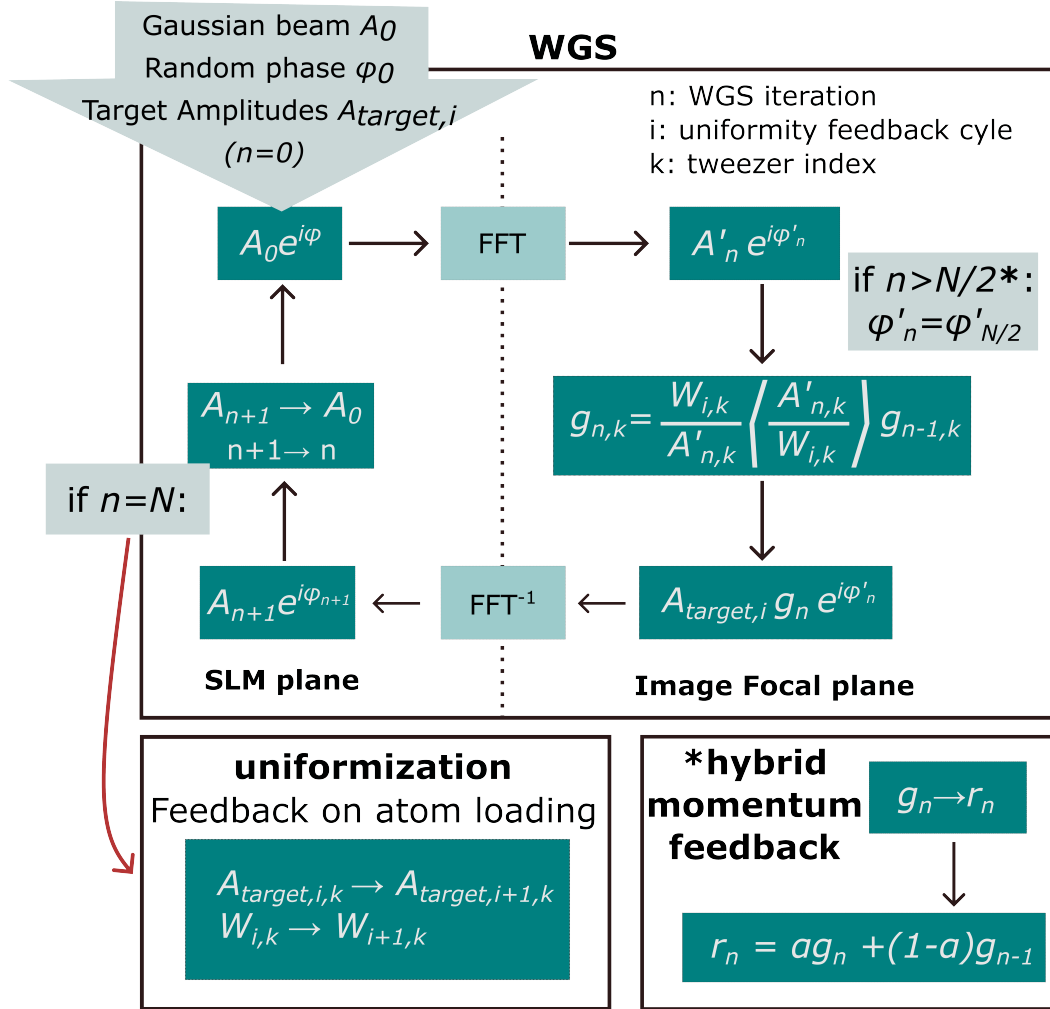


Figure 2.13: **WGS Algorithm.** A schematic describing the process of the Weighted Gerchberg-Saxton (WGS) algorithm, iterated over index n , for the i^{th} round of uniformization at tweezer sites $\{k\}$. Included in the teal arrow are the input parameters. The left-hand side of the box labelled ‘WGS’ describes parameters of amplitude and phase in the *SLM plane*, and the right-hand side describes parameters of the amplitude and phase of the light in the *Image Focal plane*. We use about $N = 50$ iterations and fix the phase ϕ'_n after $N/2$ iterations. The red arrow indicates the iteration and update applied from the feedback of an atomic loading signal. We implement about 5 rounds of feedback on this signal, updating the $A_{target,i,k}$ in the WGS algorithm each time based on the weights determined from tweezer loading. The asterisk indicates a new *hybrid momentum* WGS that allows for both more consistent convergence and lower error as seen in Fig. 2.17.

Initializing the WGS weights $A_{target,i,k} = \sqrt{H_{i,k}}$ yields an effective feedback W^{eff} which interpolates between slower amplitude-based feedback and faster intensity-based feedback, bridged by the weighting factor G .

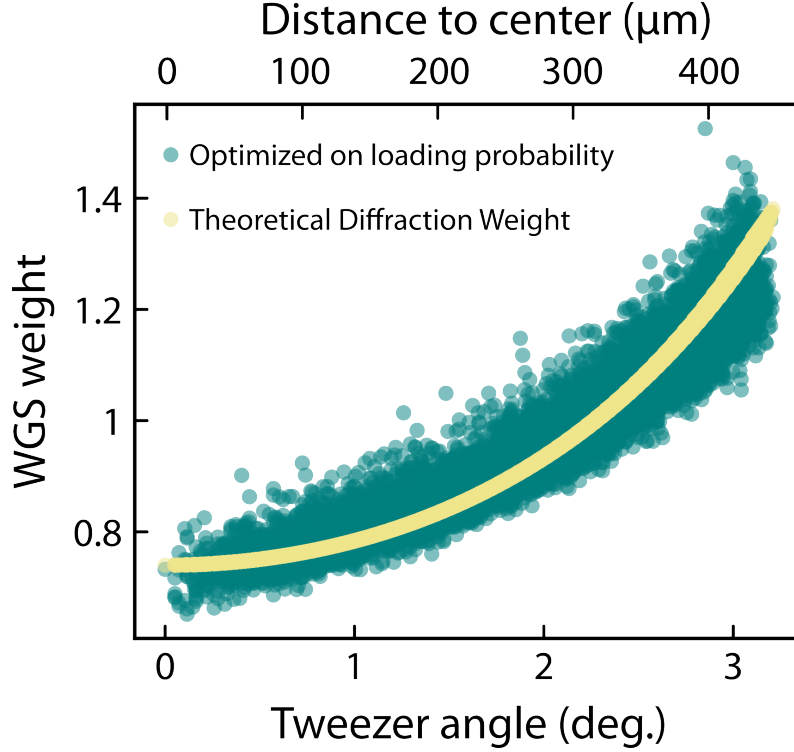


Figure 2.14: **WGS Weights after Loading-based uniformization** The WGS weights given to tweezers during the tweezer homogenization procedure, as a function of angular distance from the 0th-order reflection off the SLM, with the physical distance this corresponds to given our optical setup shown on the upper axis. In teal is plotted the weights obtained after the tweezer depths are uniformized experimentally based on loading probability. Note that this plot corrects the version in the original extended data of [59], where the weights were calculated for a gain of $G = 1$, though the data had been generated with $G = 0.6$. In yellow is shown the weight compensation that would be expected based on the diffraction efficiency, as shown in 2D in Fig 2.11. The weight increases with larger angle in order to compensate for the diminishing diffraction efficiency as a function of tweezer distance to the center. This additionally informs our decision to create a circularly shaped array.

$$W_{i,k}^{\text{eff}} = GH_{i,k} + \sqrt{H_{i,k}}(1 - G). \quad (2.63)$$

where $H_{i,k}$ is determined by adjusting the value from the previous iteration using the loading probability per tweezer $P_{\text{load},k}$, normalized by the average loading probability,

$$H_{i,k} = H_{i-1,k} \left[1 + h \left(1 - \frac{P_{\text{load},k}}{\langle P_{\text{load}} \rangle} \right) \right]. \quad (2.64)$$

$H_{0,k}$ can be chosen to match a theoretical diffraction efficiency across the target

array, in order to compensate for this effect, and commence feedback with a higher loading signal. This choice well captures the required weight distribution found with loading-based uniformization when beginning with a uniform initial $H_{0,k}$, as pictured in Fig. 2.14.

We now take an example, assuming a starting condition of $H_{0,k} = 1$. We can see that if a tweezer does not load at all, $H_{i,k} = (1 + h)$ is above 1. Then the application of the normalized weight in equation 2.61 will cause the algorithm to add more weight to that particular tweezer.

In the above expressions, G and h are used as gain factors in order to avoid oscillatory behavior in convergence [100] and are empirically determined (in our case we set both to 0.6). We have found it helpful in practice to cap the allowed values for $H_{i,k}$ to improve convergence stability.

WGS for scaling

We now look at the efficacy of the WGS algorithm as we scale up the system size of tweezers that are being generated. First, we examine the theoretical convergence of the algorithm itself, without adding additional experimental ‘uniformization’ feedback based on loading that was described in the previous section. Instead, we work with fixed goal weights across the array given by $W_k = 1$, which is a uniform array of ones unless otherwise specified in this section. We particularly study this for a standard SLM with pixel count 1920×1200 , and where relevant, fix the parameters of our experiment including a trapping wavelength of 1061 nm and an effective focal length of our objective to be 7.97 mm. As a metric for convergence, we study the error signal at each iteration of the WGS algorithm. We define the error signal $E_n = \sigma(C_n)$ on the n^{th} iteration as the standard deviation of the elements of C_n . $C_{n,k}$ represents the multiplicative correction required on the n^{th} iteration at tweezer site k :

$$C_{n,k} = \frac{W_{i,k}}{A'_{n,k}} \left\langle \frac{A'_{n,k}}{W_{i,k}} \right\rangle. \quad (2.65)$$

We study this quantity first as a function of array size (specified by FOV coverage) with a fixed tweezer number (Fig. 2.15a and b). Secondly we also examine the scaling of tweezer number with fixed inter-tweezer spacing, covering incrementally more of the FOV by adding more tweezers (Fig. 2.16a and b). For these data, we average over 50 repetitions of the WGS algorithm implementation without feedback,

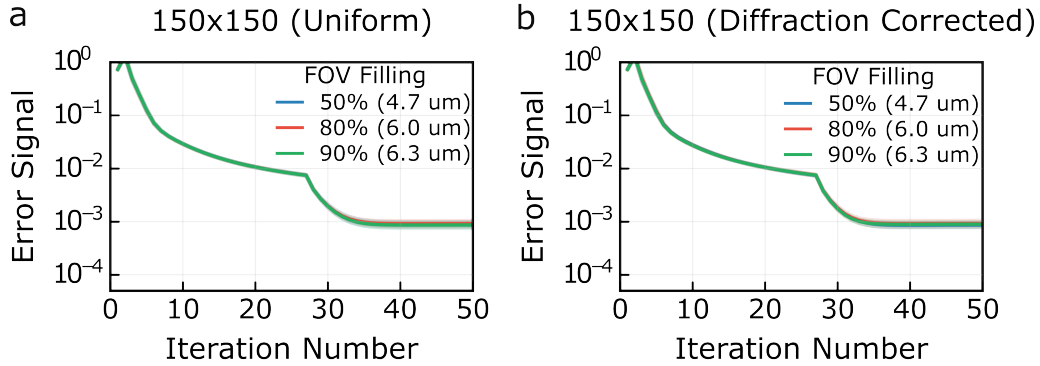


Figure 2.15: **WGS scaling for baseline and hybrid WGS** The WGS convergence for an array with a constant number of 150×150 tweezers, covering different areas of the SLM FOV both for **a**, uniform goal tweezer weight distribution and **b**, a goal weight distribution that compensates for the diffraction efficiency of the SLM such that the amplitude weight is given as $W_k = 1/\text{sinc}(\frac{px}{\lambda f})\text{sinc}(\frac{py}{\lambda f})$. The convergence is quantified by the error signal as defined in equation 2.65, and the area covered is scaled by changing the spacing between the tweezers in the 150×150 tweezer array. 50 WGS iterations are implemented and the phase is fixed on the 27th iteration for faster convergence [98]. The shaded regions represent the standard deviation of the data.

with an amplitude goal of uniform tweezer depth. We also fix the phase, as described in the previous section, at iteration 27 for each repetition, as is visible from the kink in the error signal curves around this value.

We see that for a fixed number of tweezers with increasing FOV, the algorithm converges to the same error signal within error bars for a 150×150 array (Fig. 2.15). We have also checked, not pictured, the same for a 50×50 array. We look at this scaling both for a uniform goal amplitude as well as one that is compensated for diffraction efficiency such that more amplitude is directed towards edge tweezers. Both converge to the same error within error bars. This makes clear, as may be intuitive, that scaling up the size of the array itself with fixed tweezer number does not impact the performance of the algorithm.

However, when we scale up the *total number of tweezers* that are being generated, we see that the error signal increases with increasing tweezer number, which holds true both for $7.2 \mu\text{m}$ spacing and $5.0 \mu\text{m}$ spacing, as seen in Fig. 2.16a and b. The degradation of the WGS performance with scaling can thus be attributed to the increasing tweezer number rather than the increasing area covered in the FOV.

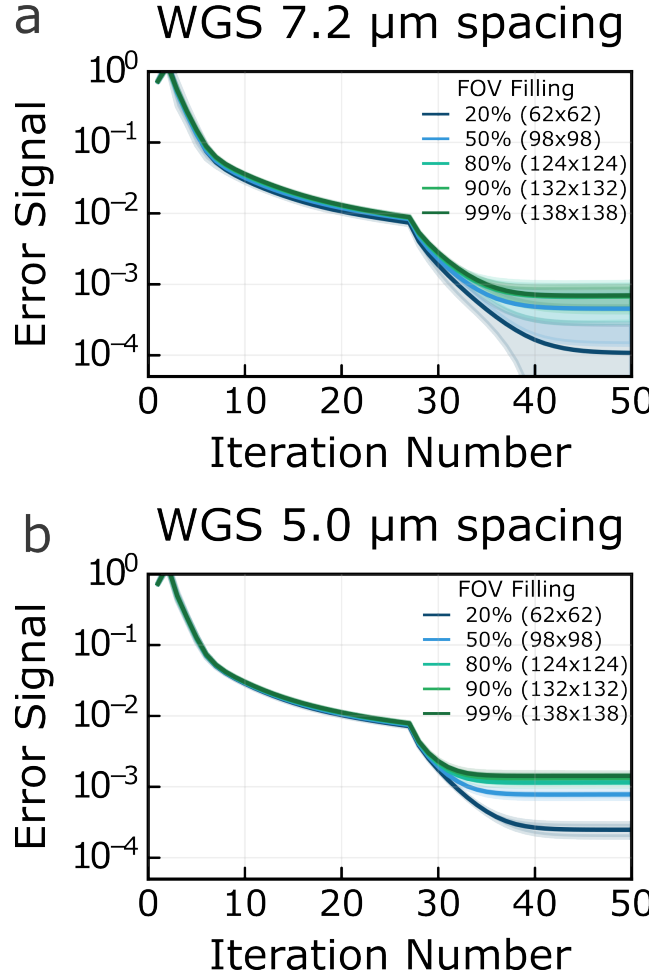


Figure 2.16: **WGS Convergence scaling for two trap spacings.** The error signal is plotted as a function of the WGS iteration number, averaged over 50 runs for each value of ‘FOV filling’. The spacing is maintained in the generated arrays, and so the percentage of the FOV that is filled determines the number of tweezers being generated. We do this for both **a**, 7.2 μm spacing as well as **b**, 5.0 μm spacing and plot FOV fillings: {20%, 50%, 80%, 90%, 99%}. The shaded regions represent the standard deviation of the corresponding data. We begin with random phases and update the phase as described in the previous section for iterations 0 through 26, fixing the phase thereafter. It is clear for both cases of spacing that increasing the number of tweezers limits the error signal that is reached after 50 WGS iterations to a higher value.

Higher tweezer numbers reduce the number of pixels that can be used per tweezer, thus increasing phase interference that may be detrimental. We plot the error signal as a function of number of tweezers in Fig. 2.18, including all data pictured in Fig. 2.16 for the ‘baseline’ case. We also make a comparison with another hybrid algorithm that we implement based on established momentum techniques for phase

estimation [94]. This hybrid-momentum WGS algorithm (HM-WGS) allows us to converge to lower error signal for every array size.

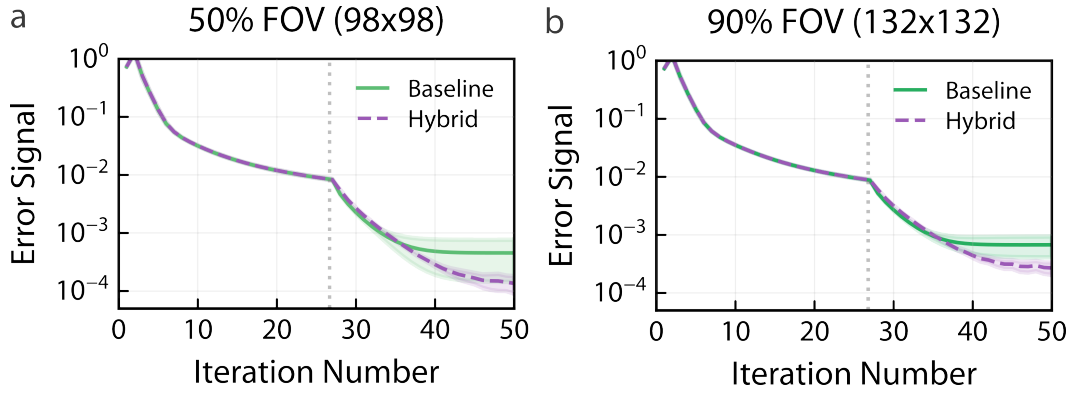


Figure 2.17: **Hybrid feedback for $7.2\mu\text{m}$ spacing WGS** Convergence of baseline (green, solid) and hybrid (purple, dashed) approaches for two array sizes **a**, 98×98 and **b**, 132×132 . The shaded regions represent the standard deviation of the 50 repetitions that we average over. The vertical dashed line represents the point at which the phase is fixed and the feedback switches from the standard approach to the momentum-based approach described in equation 2.66. The hybrid approach outperforms the baseline by a factor of ~ 3 for these array sizes with significantly smaller error bars.

Hybrid-momentum WGS

We observe that in the convergence of the WGS algorithm, the error bars on the final signal can be variable for different array spacings as seen in Fig. 2.16. With inspiration from input-output algorithms [94] which use a ‘momentum’ term dependent on a previous step in order to prevent stagnation and encourage faster convergence, we apply this in conjunction with the WGS algorithm. In particular, we replace the cumulative weight $g_{n,k}$ with $r_{n,k}$, such that

$$r_{n,k} = \alpha g_{n,k} + (1 - \alpha)g_{n-1,k}. \quad (2.66)$$

When using this feedback throughout all N WGS iterations, we observe non-monotonic convergence prior to fixing the phase after the first $\sim N/2$ iterations. Thus, we decide to implement the standard WGS algorithm prior to the phase being fixed, and afterwards we implement the momentum-based WGS. Thus, we dub this method a ‘hybrid-momentum’ WGS approach for the purposes of this study ⁶.

⁶We have implemented this updated approach after the work published in [59].

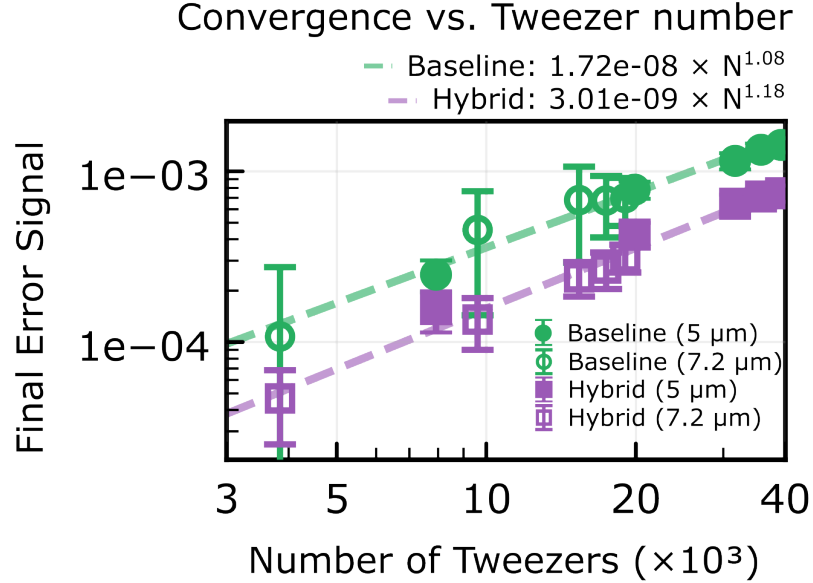


Figure 2.18: **WGS scaling for baseline and hybrid WGS** Convergence for the WGS algorithm baseline (green, circular) and hybrid-momentum (purple, square) approaches as a function of tweezer number. We plot the final error signal, from the 50th iteration of the convergence algorithms, averaged over 50 repetitions. Solid markers represent data taken for arrays with 5 μm spacing, and open markers represent data taken for arrays with 7.2 μm spacing. Both trends are linear on the log-log plot, and so we fit the power law scaling for both cases. We observe slightly steeper scaling for the hybrid case, but a coefficient that is ~ 6 times smaller, thus allowing for the hybrid case to outperform the baseline for all cases plotted. The R^2 values for the fits are 0.97 and 0.98 for the hybrid and baseline respectively.

We compare the hybrid-momentum approach with the baseline for several cases, including for the 98×98 and 132×132 arrays with 7.2 μm spacing that are also shown in Fig. 2.16a. For the data generated here, we use $\alpha = 0.8$. We compare the baseline WGS approach described previously with the hybrid-momentum approach and observe in both cases a decrease in the error signal on the final iteration by a factor of 2.5-3 as well as a standard deviation that is around 4 times smaller in the case of the 132×132 array and 7 times smaller in the case of the 98×98 array.

We plot together the data that we collect with the hybrid and baseline approaches as a function of tweezer number in Fig. 2.18 and observe power law scaling of the error signal with respect to the number of tweezers in both cases. The scaling of the hybrid approach with tweezer number— $N^{1.18}$ — is slightly steeper when compared with the baseline WGS which scales as $N^{1.08}$. However, the coefficient for the hybrid case is almost a factor of 6 smaller than for the baseline, thus clearly outperforming

the baseline for all the cases we test, up to $\sim 40,000$ tweezers with the 1920×1200 pixel SLM.

In order to create this number of tweezers, not only does there need to be enough power that can be diffracted to the tweezer location and compensate for the lower diffraction efficiency at the edges of the SLM FOV, we also need to be able to pack tweezers close together. We next look at factors that limit the spacing that we can obtain when using one SLM.

Off-axis traps from WGS

The phases that are being applied to the SLM pixels in the Fourier plane of the objective are interfering with each other in the focal plane to create the tweezer amplitude distribution that we target, as described in the last section. However, out-of-plane from the focal plane, we have not specified a distribution, and the phases are not uniform across the tweezers. Thus, out-of-plane tweezer interference is not predictable.

We can, however, calculate the out-of-plane tweezer interference close to the focal plane by returning to the formalism laid out in section 2.1.4, and with the Fresnel approximation in equation 2.24, and substituting $z = f + \delta z$. This results in the following equation for the electric field at a distance δz from the focal plane:

$$E(x, y, f + \delta z) = \frac{e^{ik(2f+\delta z)}}{i\lambda f} \iint E_0(x', y') e^{\frac{ik}{2f}(x'^2+y'^2)(1-\frac{(f+\delta z)}{f})} e^{-\frac{ik}{f}(xx'+yy')} dx' dy' \quad (2.67)$$

$$\propto \mathcal{F}\{E_0(x', y') e^{-\frac{ik\delta z}{2f^2}(x'^2+y'^2)}\} \quad (2.68)$$

In order to understand the out-of-plane interference from a phase pattern generated by the WGS algorithm, we can just define $E_0(x', y', 0)$ in equation 2.68 to be an aperture-clipped Gaussian beam with an input phase defined by the WGS-optimized SLM phase pattern. In Fig. 2.19, we show out-of-plane interference for different tweezer spacings calculated for slices of δz in the range $\delta z = \{-15\mu m, 15\mu m\}$.

Comparing Fig. 2.19a ($7.2\mu m$ spacing) and b ($5.0\mu m$ spacing) makes clear that the closer the spacing, the more out-of-plane interference will occur, with out-of-plane traps at $5.0\mu m$ spacing being just as strong as the traps at $z = 0$ that we intend to use for trapping atoms. The images shown are vertical slices along a row in a 122×122

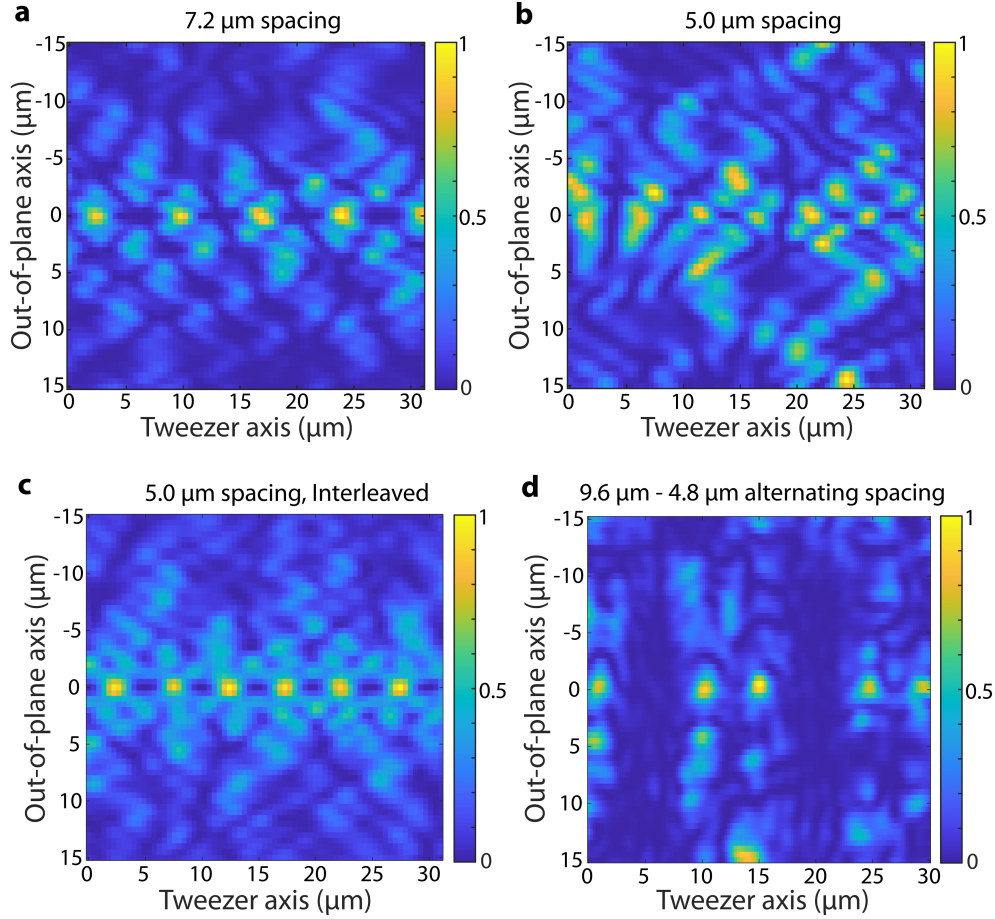


Figure 2.19: **Out-of-plane traps** Calculation of an out-of-plane intensity profile from a section of a 122 by 122 site tweezer array at 1061 nm, whose phase pattern is generated by a WGS algorithm. The tweezer axis (x-axis in plots) is centered along a selected row in the array and the out-of-plane axis (y-axis in plots) is perpendicular to the focal plane of the tweezers, along the direction of light propagation. The focal plane for the tweezers is at 0 μm , and we simulate for different spacing between the tweezers: **a**, 7.2 μm and **b**, 5.0 μm . We also show in **c**, the case for which 5.0 μm spacing is achieved by alternating traps generated with two different lasers such that they do not interfere by using for example orthogonal polarization or sufficiently different wavelengths. In **d**, we show the out-of-plane interference for the case of alternating spacing between tweezers of 9.6 μm and 4.8 μm as is used in section 4.2.4. The color scale for each case is normalized by the highest intensity in the simulated slice.

array generated with the baseline WGS algorithm for different spacings, assuming our experimental parameters of wavelength 1061 nm and objective effective focal length of 7.97 mm. If instead of creating a 5.0 μm spaced array using a single SLM, we combine two 10.0 μm spaced arrays using two SLMs with light from different

wavelengths or polarizations, we can avoid the strong ‘spurious’ trap interference and create a uniform array in the focal plane(Fig. 2.19c). For some applications, an alternating patten of tweezers being closer and further is useful for e.g. avoiding atom loss during movement, and so we also show a calculation of the out-of-plane interference for the case of alternating $9.6\text{ }\mu\text{m}$ and $4.8\text{ }\mu\text{m}$ spacing in Fig. 2.19d.

Note that we observe these spurious traps in practice to be experimentally trapping atoms even at a spacing of $7.2\text{ }\mu\text{m}$. We implement several procedures in order to expel atoms trapped in these out-of-plane traps, as previously described in section 2.3.2, which requires over 400 ms of our experimental sequence, constituting almost half the time.

2.4.4 Acousto-Optic Deflectors

In addition to SLMs, Acousto-Optic Deflectors (AODs) have also been used to create optical tweezer arrays and to rearrange atoms into defect-free arrays [30, 31]. In tandem with static arrays created with SLMs, AODs have recently been used in order to move atoms coherently for implementing circuits between atomic qubits, and even for creating and implementing operations on logical qubits [35, 51].

AODs utilize radio-frequency (rf) waves traveling through a crystal in order to diffract light. The larger the active aperture of the crystal, the more peaks/troughs of the acoustic wave are sampled by the input beam, thereby decreasing the effective spacing of the diffraction grating. Thus, the larger the crystal, the more resolvable spots the AOD can access in the FOV. In our experiments we utilize pairs of 15mm active aperture AODs (Gooch & Housego AODF 4085) in a crossed configuration such that one AOD controls diffraction in the x-axis and the other controls diffraction in the y-axis. We map the output after the pair of AODs to the back aperture of the objective using a telescope with 3:2 demagnification to match the aperture diameter of the objective.

Several aspects about AODs require optimization with scaling to larger array sizes, which we examine in further depth in chapter 4 when we demonstrate long-distance transport.

2.4.5 Large Array Aberration Considerations

The objective that we have for tweezer creation has been designed to correct for aberrations across a full 1.5 mm FOV, and so for the smaller SLM FOV, the objective is certainly not limiting. That being said, the SLM is reflective and the 0^{th} order

also needs to be filtered out, and we must implement at least one telescope in order to map the Fourier plane at the SLM to the back of the objective. This telescope can be a 4f direct mapping, or it can be a (de-)magnifying telescope to match the SLM size to the objective aperture diameter.

The optical elements throughout the relay paths are thus important to account for when it comes to aberrations that can be imparted to the tweezers. In particular, large arrays are subject to aberrations arising from large beams incident on lenses causing spherical aberrations, and beams with large angles of incidence on lenses that give rise to astigmatism.

We picture a schematic version of the current experimental pathways that are being utilized in our experiment (Fig. 2.20), and proceed to describe a few considerations for aberrations that can arise for large arrays along the tweezer paths which are combined. We note that this experimental setup has been updated since the first work that was published [59]. We then describe ways in which we actively correct for aberrations across the array using Zernike coefficients—both globally and locally.

Relay lenses for large cone half-angles

As pictured in the above Fig. 2.20, we utilize many 4f relays in order to combine AOD and SLM pathways. Thus, these lenses must be chosen such that they do not induce aberrations on the optical tweezers. In particular, considerations arise regarding the size of optics as well as the *cone half-angle* (CHA) that they can accept.

We utilize Thorlabs lenses (TL200-CLS2, FTH160-1064, and FTH254-1064) throughout the pathway. For a lens after our SLM, to be aberration free for tweezers across the full SLM FOV, the lens must account for a CHA of $\sim 3.8^\circ$, as determined from the grating equation 2.49. If the beam is magnified to a larger size with a 4f magnifying telescope, the CHA requirement on future lenses will subsequently decrease. We discuss further some of the fundamental limitations on optical elements for scaling up array sizes in Chapter 5.

Combination cubes

We note that there are several combination cubes throughout the optical pathway, including both polarizing beam-splitting (PBS) cubes as well as non-polarizing beam-splitting (NPBS) cubes. All of the cubes in the relay paths are 2 inches

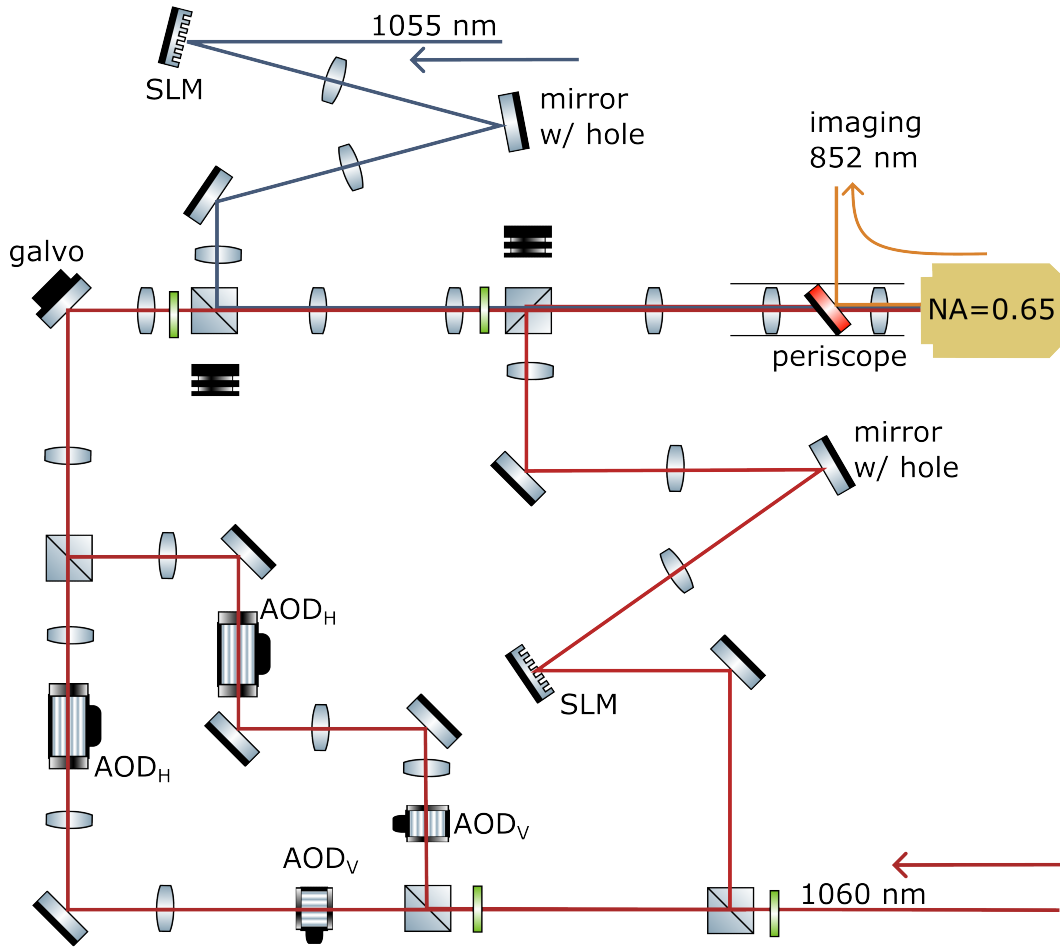


Figure 2.20: **Updated Tweezer paths.** SLM and AOD tweezer combination pathways are pictured as a schematic to depict array combination and pathway complexity. Two SLMs are used (Meadowlark 1920x1200 pixel HDMI-controlled SLM on the 1061 nm path and Meadowlark HSP1920x1152 pixel PCIE-controlled SLM on the 1055 nm path). The AODs are combined and then incident on a galvo mirror with 5 ms switch time, that will allow them to cover the full FOV of the SLMs. We note that in the depicted configuration, polarization is used for recombination twice, and thus power will inevitably be lost along some path. However, a dichroic approach to recombination could be used—transmitting for example S and P polarization of one wavelength and reflecting one polarization of another wavelength.

and are made of N-SF1 with an index of refraction of $n = 1.69$ at the relevant wavelengths. From Snell's law, we are aware that the geometric focusing will be affected by the index of refraction of the material through which the light is traveling. For a lens with working distance D with a cube in the path of diameter d , the focal

plane will be a distance

$$f_{\text{new}} = D + d \left(1 - \frac{\cos \theta}{\sqrt{n^2 - \sin^2 \theta}} \right), \quad (2.69)$$

where θ is the angle between the incident beam and the normal vector to the surface of the cube. For small θ , this reduces to $f_{\text{new}} = D + d(1 - 1/n)$, which represents a global focal plane offset for the whole array, with corrections for individual tweezers that have different incident angles. When using 2-inch cubes, this global offset in the focal plane can be around 2 cm and is necessary to account for when aligning relay optics.

SLM Zernike Corrections

As has been previously shown, loading across a tweezer array can be increased using Zernike polynomials to correct for tweezer aberrations and increase the average trap depth [102]. This can be done globally for a whole array, or to correct particular tweezers if combined with a tweezer generation approach that combines phase patterns for individual tweezers [92, 103], rather than the global FFT procedure that we described in this thesis. Early on in this work, we have utilized such Zernike polynomials globally to compensate for aberrations across the full array.

Zernike polynomials are mutually orthogonal on a unit disk, whose radius we map to the radius of the Gaussian beam that is incident on the SLM. However, given that we create such a large tweezer array, there may be different coefficients that are optimal for different parts of the array. As a result, there is indeed some coupling between polynomials when it comes to optimizing on a metric of average loading probability for instance. Additionally, there is coupling between Zernike optimization and uniformization of intensity for each trap. For instance, if more power is diverted to aberrated tweezers in order to compensate for the diminished trap depth when uniformizing the array, then this may cause Zernike optimization based on trap-depth or loading rates to be skewed. There can also be cross-talk between Zernike polynomials themselves if, for instance, the beam is not centered on the center of the applied Zernike phase on the SLM. In this scenario, a defocus phase pattern can both move the focus in z, but also in x or y (by effectively applying a tip/tilt Zernike phase), depending on the relative spacing of the beam center.

As a result, when optimizing the polynomials globally, it has become clear that it is best to compensate for the theoretical diffraction efficiency such that there is sufficient loading in each trap, and then when sweeping Zernike polynomials, to

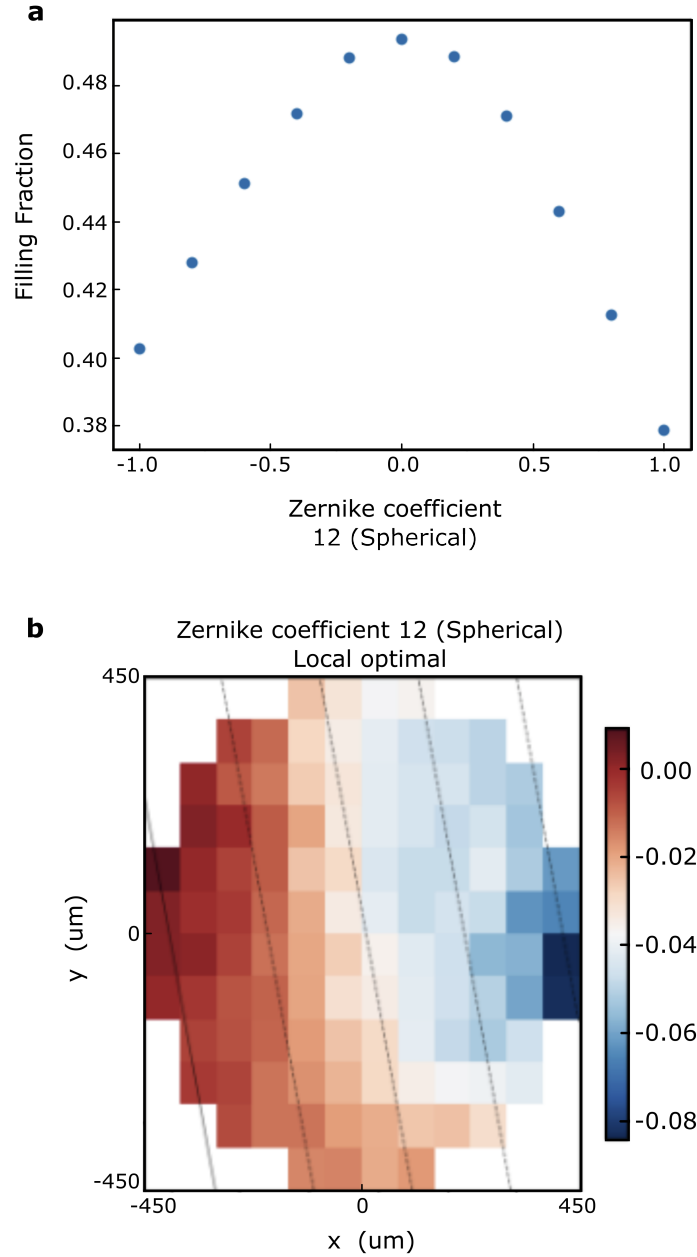


Figure 2.21: **Zernike Polynomials.** We include **a**, An example of one polynomial (the 12th order polynomial, associated with spherical aberrations) being swept for loading probability averaged across all sites, after which it will be set to the optimal value. **b**, We depict optimal Zernike values for different parts of the array, averaged by section to demonstrate array size effects. We recently have pivoted to feedback based on local Zernike being implemented in tandem with iterative feedback following the algorithm presented in [92]. Note that the data presented in **a** and **b** are from two separate optimization scenarios at different times.

update and save the new optimal as the baseline for the sweep of the next polynomial. We depict one example of an optimization sweep of the coefficient for the 12th Zernike polynomial which impacts spherical aberrations in Fig. 2.21a, where we use the metric of filling fraction, averaged across the array. In this scenario, we set the coefficient to the optimal value, before choosing another polynomial at random to optimize. From implementations of global Zernike optimization, we have been able to increase trap depths by around 7 – 10% across the full array. However, as can be seen in Fig. 2.21b, if the filling fraction is evaluated more locally, there is a different local optimum. Thus, we find further improvement when implementing local Zernike optimization [92]. In particular, we find that we can increase the filling fraction of an array from $\sim 25\%$ to $\sim 50\%$ purely by sweeping through local optimization of Zernike polynomials⁷.

2.5 6100 atoms

Using the techniques described in section 2.4, we generate 11,998 tweezers⁸ in which we trap over 6,100 atoms on average in every experimental sequence. In Fig. 2.22, we show both a single-shot image of our atomic array, as well as an averaged image where every tweezer site can be seen. This number of traps and qubits represents almost an order of magnitude advance from the state-of-the-art tweezer array sizes that had been demonstrated until this point.

We space tweezers by about 7.2 μm to avoid strong off-axis traps (Fig. 2.19). The full array is about 900 μm in diameter and made up of tweezers at wavelengths of 1061 nm and 1055 nm generated by two different SLMs, and combined on a cube in the configuration pictured in Fig. 2.12a.

2.5.1 Tweezer loading

We demonstrate uniform loading across the sites in the array. To load single atoms in the tweezers, we cool and then parity-project [27] from a $\sim 1.6 \text{ mm } 1/e^2$ diameter magneto-optical trap using D_2 light. Before imaging the atoms, we use a multi-pronged approach to filter out atoms in spurious out-of-plane traps, residual from the SLM tweezer creation (Fig. 2.6).

⁷See details for this implementation in future thesis from the group by Elie Bataille.

⁸We exclude 3 of the total 12,001 generated tweezers from our scientific results. These tweezers are in the center of the array and are impacted by some remaining 0th order SLM light that is imperfectly filtered. These tweezers occasionally load more than a single atom, and they have somewhat reduced imaging fidelity due to higher background. We note that quantum computation can be done despite defects in an array, and that thus one can choose to avoid a small number of traps [104]

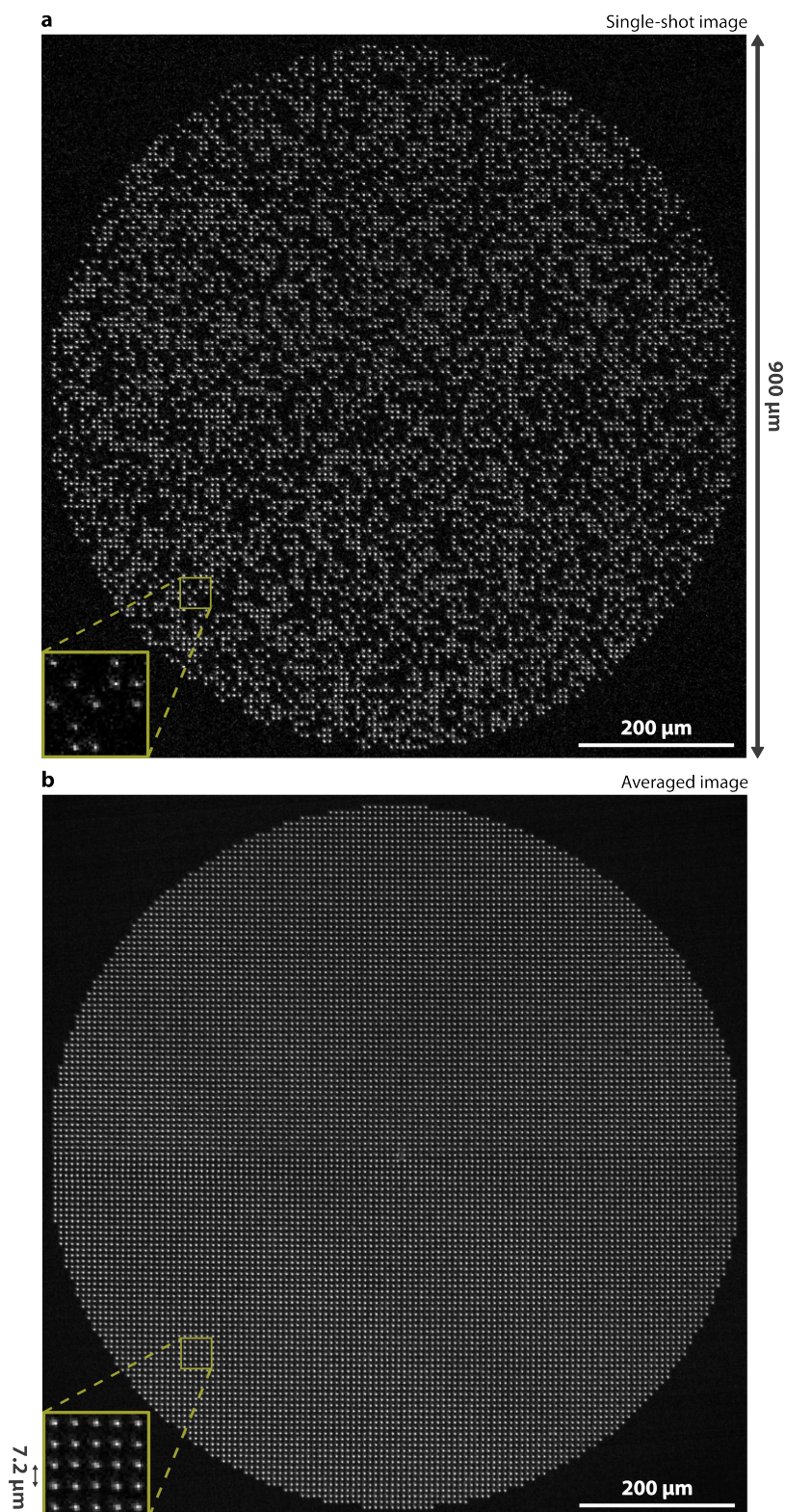


Figure 2.22: **Atom loading in 11,998 tweezer array** **a**, A single shot image of the stochastically-loaded array. **b**, An averaged image over the array. Insets are close-up views of the selected subsection. Atoms are spaced by 7.2 μm and the array is about 900 μm in diameter. Contrast is enhanced for visual clarity

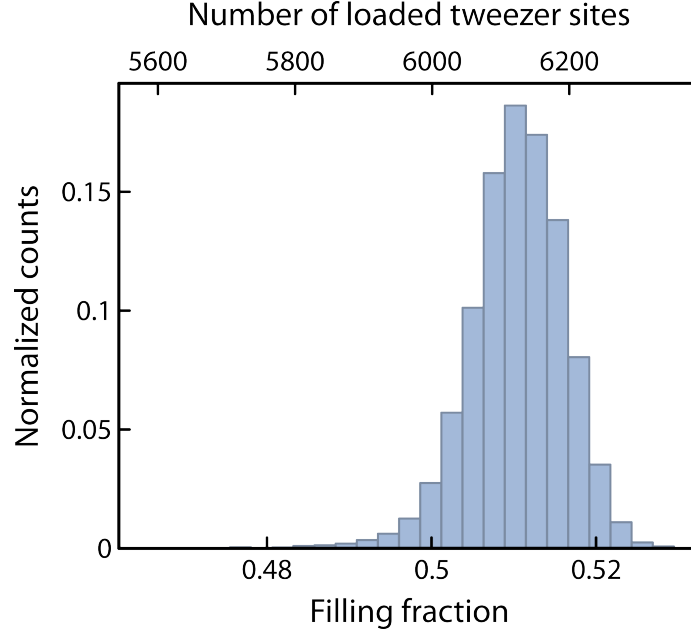


Figure 2.23: **Filling fraction** Histogram of number of atoms loaded into tweezers in each image, averaged over 16,000 images. On average, we load 6,139 single atoms (51.2% of the array), with a relative standard deviation of 1.13%.

We then zero the magnetic field and apply two-dimensional polarization gradient cooling (2D PGC) in the atom array plane for fluorescence imaging of single atoms, which simultaneously cools the atoms. Imaging light is applied for 80 ms, and photons are imaged on a quantitative CMOS camera (qCMOS). We load 6,139 single atoms on average per experimental iteration (51.2% of the array on average), with a relative standard deviation of 1.13% (averaged over 16,000 iterations). The histogram of the number of atoms we observe in each image is pictured in Fig. 2.23.

Not only do we on average load over 6,100 atoms per shot, we also verify that each tweezer loads about half the time. In Fig. 2.24a, we plot a color map of each tweezer's loading probability, which is also imaged in histogram in Fig. 2.24b. We find an average loading probability per site of 51.2% with a relative standard deviation of 3.4%. We note that we use the loading rate per site as the metric with which to feedback on the WGS weight to uniformize the array trap depth (as first demonstrated in [100]).

2.5.2 Tweezer parameters

While we load the tweezers uniformly, the atoms care additionally about the uniformity of other tweezer array parameters. For instance, trap depth non-uniformity

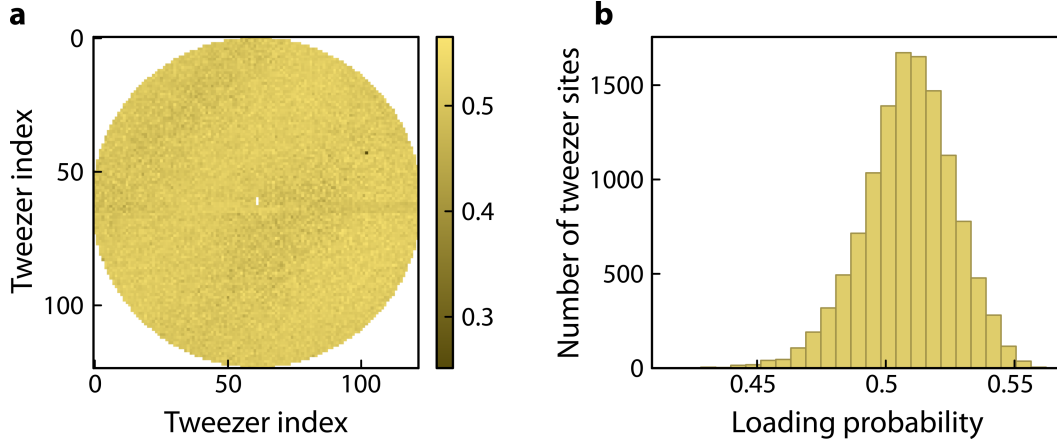


Figure 2.24: **Loading probability per-site** **a**, Map of the loading probability per site across the tweezer array. **b**, Histogram of loading probability per-site across the tweezer array. The lowest loading probability is 25.1% for one tweezer, which is the only tweezer not shown in the histogram, but included in the quoted average and seen in the map. This tweezer, however, does not exhibit a significant difference in imaging survival probability, coherence time, or single-qubit gate fidelity. Therefore, this tweezer site could also be utilized in the experiment.

contributes to dephasing during single-qubit gate operations since the resonant drive is detuned by amounts corresponding to site-dependent light shifts. Trap frequency across the array is important for understanding heating considerations among other aspects impacted by tweezer shape.

Trap depth

We subsequently measure the trap depth across the atomic array to verify the array uniformity (Fig. 2.25). Such a measurement can be done with a few approaches involving differential light shift measurements—comparing the difference in Stark shift experienced by states that have different polarizabilities, in order to determine the intensity of the laser or depth of the trap.

We first describe a measurement approach that utilizes a differential light shift between the $6S_{1/2}$ and $6P_{3/2}$ states [102]. In particular, we prepare the atoms in the $F = 4$ hyperfine ground state by leaving the repump on longer than the PGC beams after loading and cooling. We then scan the frequency of the $F = 4 \rightarrow F' = 4$ which is pulsed for 10 μs , depumping the $F = 4$ population into the $F = 3$ ground state when it is on resonance. We fit the $F = 4$ population loss to a Lorentzian and calculate the difference with the resonance frequency, constituting the differential light shift. This is the measurement that we use to take the data presented in

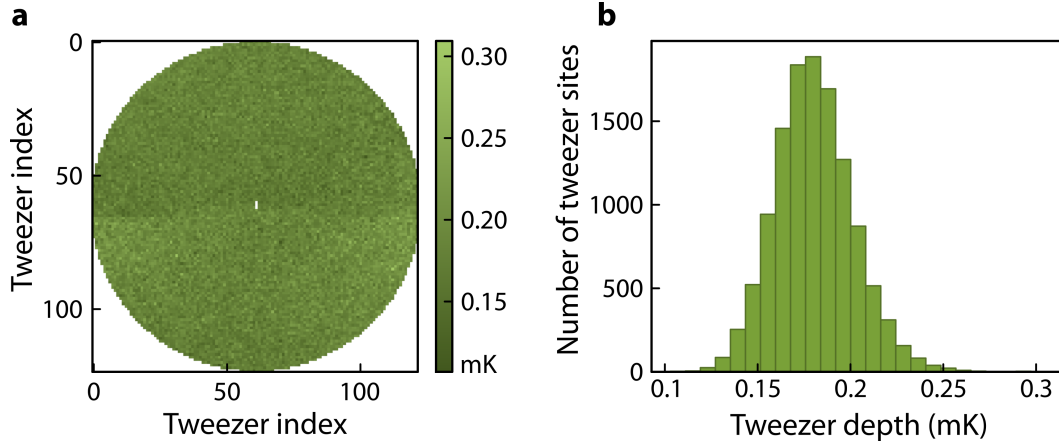


Figure 2.25: **Trap depth per-site** Site-resolved trap depth **a**, map and **b**, histogram. Across the 12,000 tweezer array using the $F = 4 \rightarrow F' = 4$ depumping [102], we perform a differential light shift measurement and observe an average trap depth of $k_B \times 0.18(2)$ mK with a standard deviation of 11.4% across the sites.

Fig. 2.25. The measurement has the downside of the trap depth being washed out by the difference in hyperfine magnetic sublevel polarizabilities, as can be seen in Fig. 2.26.

The per-site tweezer depth map and its histogram is pictured in Fig. 2.25. This is obtained by measuring the differential light shift on $F = 4 \leftrightarrow F' = 4$ D₂ transition [102]. We measure an average trap depth of $k_B \times 0.18(2)$ mK with a standard deviation of 11.4% across the sites.

We would like to characterize the standard deviation of intensity differences across our 11,998 tweezers, which is combined in this measurement with the atomic distribution across hyperfine Zeeman sublevels. In order to isolate the standard deviation of the intensity, we can interpret this measurement in the following manner. Given that the trap depth $U \propto I\Delta\alpha$ that we measured and plot in Fig. 2.25 is a function of independent random variables I and $\Delta\alpha$, where I is the beam intensity and $\Delta\alpha$ is the differential polarizability, we can write the square of the measurement standard deviation as

$$\left(\frac{\sigma_U}{U}\right)^2 = \left(\frac{\sigma_I}{I}\right)^2 + \left(\frac{\sigma_{\Delta\alpha}}{\Delta\alpha}\right)^2 \quad (2.70)$$

$$= \left(\frac{\sigma_I}{I}\right)^2 + \frac{\frac{1}{N} \sum_N (\Delta\alpha_N - \langle\Delta\alpha\rangle)^2}{\langle\Delta\alpha\rangle^2}. \quad (2.71)$$

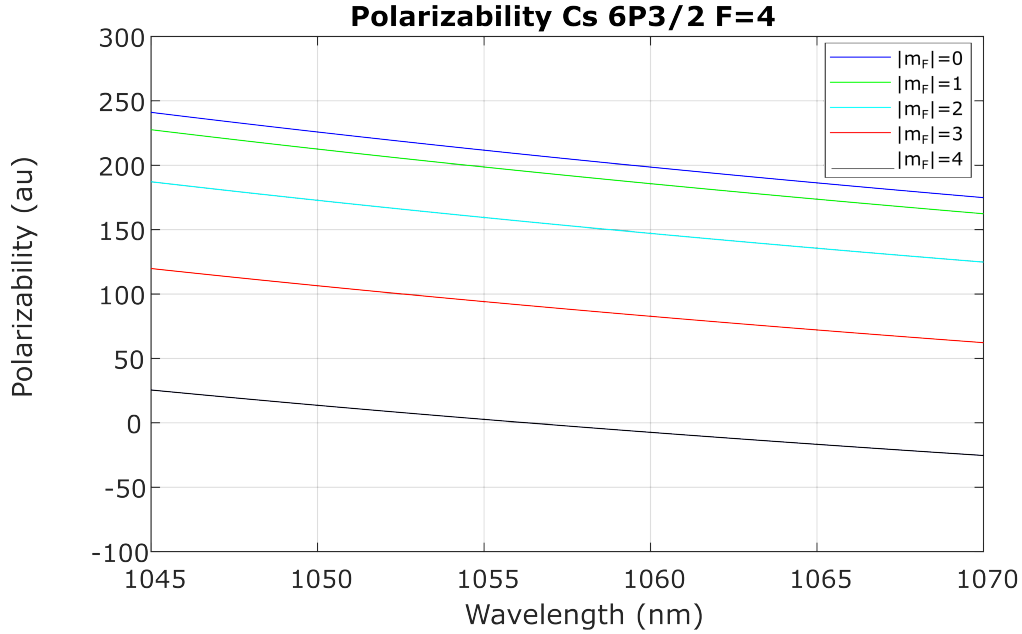


Figure 2.26: **Polarizability of $6P_{3/2}$ in linear trap** Polarizability calculations for the $6P_{3/2}, F = 4$ state of Cs in linearly polarized light. We use the $|6S_{1/2}, F = 4\rangle \rightarrow |6P_{3/2}, F = 4\rangle$ transition for depumping during the trap depth measurement. We note that the ~ 200 au spread of magnetic sublevel state energies leads to a larger standard deviation in our trap depth measurement.

Assuming a ground state polarizability of 1181au at 1061 nm, taken from [61], an average measured trap depth of 0.18mK, and equal Zeeman hyperfine sublevel population distribution with polarizability as calculated in Fig. 2.26, we obtain the result that $\frac{\sigma_I}{I} \sim 0.105$. Thus we expect from this measurement that the non-uniformity across the intensities of the traps is around 10.5%.

In order to make a measurement that does not mix these uncertainties, we have a few options. If we were to prepare all the atoms in a given m_F state prior to measurement, this would allow the differential polarizability for all atoms to be the same, up to a state preparation error, thus giving a more accurate measurement. Another approach is to implement a differential measurement between the qubit states ($F = 3, m_F = 0$, and $F = 4, m_F = 0$) in the hyperfine ground state of $6S_{1/2}$. This approach allows the measurement to be made with higher accuracy using microwave operation or two-photon Raman transitions. In particular to perform this measurement, we prepare the atoms in the $F = 4, m_F = 0$ qubit state, and then perform a Ramsey measurement. This involves first applying a $\frac{\pi}{2}$ pulse at the clock

frequency, such that the atoms are in the state $\Psi = \frac{1}{\sqrt{2}}(|F = 4\rangle + e^{i\phi}|F = 3\rangle)$. We then evolve for *interrogation time* T , after which we apply a second $\frac{\pi}{2}$ pulse. The total phase accumulation $\Delta\phi$ is then given by

$$\Delta\phi = 2\pi \int_0^T \delta\nu(t)dt \quad (2.72)$$

where $\delta\nu(t)$ is the frequency difference between the clock frequency and the actual frequency difference between the states arising from the differential light shift as a function of time. When the shift is constant throughout the evolution, $\Delta\phi = 2\pi\nu T$. The final population in $F = 3, m_F = 0$ will evolve as a function of the interrogation time T , and in fitting this so-called *Ramsey fringe*, we can determine the accumulated phase and thus the trap depth. The differential shift between the two clock states arises from the difference in frequency between the states and the transitions that contribute to polarizability (dominated by the D_1 and D_2 lines) coming from the 9.2 GHz splitting between them.

We note that, as previously described, we optimized the WGS weights in our algorithm based on the metric of loading, rather than trap depth. This requires less data taking to accurately measure. Furthermore, dynamical decoupling can typically be implemented to account for errors arising from trap depth differences across an array. However, in the scenario where higher trap intensity uniformity were required, we anticipate that uniformizing on a trap depth measurement directly would allow for improvement of this metric.

Trap frequency

In order to further characterize the traps, we measure our trap frequencies. Initially we attempt to do so via parametric heating by modulating our trapping AOM, however we believe that we are limited in our measurement by the AOM fall time, so instead we use this measurement as a test of our Raman setup.

To perform Raman sideband spectroscopy, we must operate in a configuration where we can simultaneously implement a two-photon transition and also impart a momentum kick to change the motional state of the atom. The approach that we have chosen, as described in section 2.3.3 on Raman transitions, is a Doppler-free approach where one amplitude-modulated beam can drive a two-photon transition. This approach is beneficial for quantum computing applications where only one beam must be aligned and controlled, but it cannot impart momentum to the atoms, and thus cannot be used for sideband spectroscopy. Therefore, we split the amplitude

modulated light into two beams that are incident on the atoms with wavevectors $\mathbf{k}_1$ and $\mathbf{k}_2$ (both with axial and radial components) in order to measure the trap frequencies, akin to previous works with mode-locked lasers used to address the motion of trapped ions [105, 106]. A schematic of the setup is shown in Fig. 2.27a. In this configuration, the effective Lamb-Dicke parameter is

$$\eta_{\alpha}^{\text{LD}} = |(\mathbf{k}_1 - \mathbf{k}_2) \cdot \boldsymbol{\alpha}| \sqrt{\frac{\hbar}{2m\omega_{\alpha}}}, \quad (2.73)$$

where m represents the mass of cesium-133, and α denotes the radial or axial motion (with unit vector $\boldsymbol{\alpha}$). Out of 1 W of fiber-coupled amplitude-modulated laser light, each beam has 1-5 mW of laser power and a Gaussian $1/e^2$ diameter of ~ 2 mm.

We then scan the frequency of the amplitude modulated beam incident on the atoms after they have been prepared in the $F = 4, m_F = 0$ state. The carrier represents the resonant two-photon clock transition frequency (ν_{clock}) with no change in momentum or motional sideband ($\Delta n = 0$). The radial and axial sidebands respectively represent the case when the frequency of the two photon transition becomes resonant with $\nu_{\text{clock}} \pm \nu_{\text{trap}}$ where ν_{trap} is the radial or axial trap frequency, driving the $\Delta n = \pm 1$ motional sidebands. These resonances allow for increased state change along with the increase in motional energy by one quanta. Therefore, we see increased survival at the carrier and trap frequencies, after pushing out atoms that remain in the $F = 4, m_F = 0$ state, as seen in Fig. 2.27. In fitting these data, we find a radial trapping frequency of 29.30(4) kHz and an axial trapping frequency of 5.64(3) kHz. Trap inhomogeneity broadens the signal— from the lineshape fit, we extract standard deviations across the tweezers of the array to be 4.7 kHz and 1.9 kHz, respectively.

From this trap frequency measurement as well as the trap depth measurement, we can back out the trap waist from the Gaussian approximation (see equations 2.31), inferring a $1/e^2$ tweezer waist $w_0 = 1.17(6) \mu\text{m}$.

2.6 Image characterization

To characterize anything about atom loading including the number of atoms that we have trapped, or to use any atomic signal for trap characterization or computation, one needs to be able to map photons that are collected on the camera to a binarization of atomic presence (1) or absence (0). Additionally, in this work, for all quantum computations, we map atom presence or absence to our qubit states for readout (via

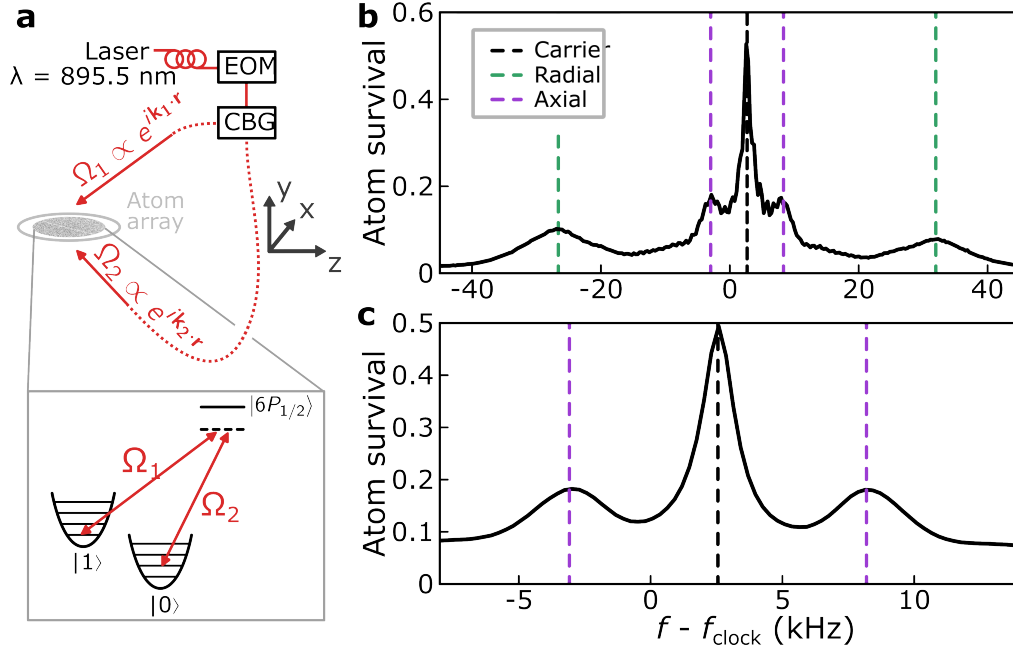


Figure 2.27: **a**, Schematic of the Raman configuration used to address the atomic motion. We use a highly dispersive double-passed chirped Bragg grating (CBG) to obtain a large bandwidth, which we scan across using a 9.2 GHz EOM. **b**, Raman spectroscopy results exhibiting sidebands corresponding to the radial motion (in green) and the axial motion (in purple). We measure a radial trapping frequency of 29.30(4) kHz and an axial trapping frequency of 5.64(3) kHz. The sideband signal is broadened due to inhomogeneities in the array. From the lineshape fit, we extract standard deviations across the array of 4.7 kHz and 1.9 kHz, respectively. The measurement is averaged over the 11,513 tweezers created with the 1061 nm light as shown in Fig. 2.12a **c**, Fine-grained spectroscopy data acquired with a lower Rabi frequency to resolve the axial sideband.

pushout of atoms in $F = 4$). Thus, implementing this characterization with high fidelity across the array is critical. Additionally, imaging atoms without inducing atom loss is necessary for quantum preparation. We thus describe various aspects of determining atomic presence, and the associated fidelities and imaging survival that we obtain across our array.

As described previously, we use 2D PGC beams in the atom plane for cooling and for imaging, at different detunings. We image the atoms onto our q-CMOS camera which has a typical RMS readout noise of 0.27 electrons for the ultra-quiet scan (readout 5 frames/s), and 0.43 electrons RMS for the standard scan (readout 120 frames/s). These readout times assume the use of all the camera rows (4096 x 2304 pixels). We typically use only about half of the rows, and achieve a readout time

of around 5 ms, which requires use of the CoaXPress connection option. The dark current is 0.006 electrons /pixel/s which is not limiting for our experiments.

Besides the noise inherent to the camera, there is inherently background noise from light in the experiment surrounding the camera. To minimize this background however, we include a line filter in front of the camera⁹, and we surround the imaging pathway with dark foil from the dichroic separating imaging light and tweezers, until the light reaches the camera.

2.6.1 Thresholding

After minimizing the background light, in order to determine the presence of an atom, we utilize photon counting statistics, as has been done since single atoms trapped in tweezers have been observed [27]. We expect only two peaks in the distribution due to trapping only one or zero atoms after parity projection which the 3D molasses accomplishes. A threshold, set between these peaks determines whether a trap contains an atom or not depending on the number of photons that have been collected. Thus, the further separated the peaks are, the higher the imaging fidelity.

Identifying atom location

To properly collect photon counting statistics, we must identify the tweezer locations on the image in order to count photons that are from the correct tweezer to determine if it is holding an atom. A naive approach is to use a combination of peak finding on an averaged image and translate the goal tweezer distribution that has been set by the WGS algorithm to find the highest overlap.

However, a large array of atoms may have imperfect spacing on the qCMOS camera due to aberrations either in the array or in the imaging pathway, making a direct mapping from the WGS-defined tweezer locations difficult. Additionally, when attempting to uniformize a newly generated array, some traps may not be loading yet, thus making the matching more difficult. While we have iterated on several approaches, the one that is the most effective for mapping tweezer positions onto an average image takes the following strategy:

We first fit the warp across the array to a 2nd or 3rd degree polynomial. We do so by implementing a gradient descent algorithm that assigns tweezers to peaks that

⁹Note that perpendicular alignment of this line filter is important due to the large image with strong k-vector difference from tweezers across the array.

have been found on the image with some likelihood, and optimizes the matching likelihood across the array. We save these optimal polynomial warp parameters that characterize the mapping between the goal input tweezer array and the actual spacing on the camera, and use these for each experiment. We only refit these parameters if the thermalization on the experiment changes via for example an increased heat load with more incident power in tweezers, a significantly altered cycle time or a change in the idle sequence.

We apply the optimal warp parameters to the whatever tweezer pattern is being used. These are constant, and can be used to fit arrays with different tweezer spacings and positions (for example, if the thermalization remains essentially the same, but we double the spacing between tweezers, we can keep the optimal warp parameters but re-scale the spacing).

After this has been implemented, we additionally add a final step of local adjustment—within a 5×5 box of pixels around the current optimal locations, we find the pixel that has the highest signal, and assign this pixel to the tweezer location.

Image threshold optimizing

After the location of tweezers has been mapped onto the images that we are taking with the qCMOS camera, we must define a surrounding region whose photon counts will be used to determine the presence or absence of the atom in each given image. The averaged point-spread function (PSF) waist radius we measure is 1.7 pixels on the qCMOS camera, corresponding to $7.8 \mu\text{m}$ on the camera plane or $1.5 \mu\text{m}$ on the atom plane. We estimate that, accounting for a finite atomic temperature (up to $50 \mu\text{K}$ in this simulation) and camera sensor discretization, the ideal PSF radius should be 1.25 pixels. We leave an investigation of the discrepancy to future work.

In order to characterize the presence or absence of an atom in a tweezer, we account for photons within 7×7 pixels surrounding the center of the tweezer, and for the given number of photons, determine whether an atom is present based on whether this number is above or below a threshold. The chosen threshold has a strong impact on the fidelity with which we determine atom presence. A threshold that is too high will result in false negative and a threshold that is too low will result in false positives.

In order to optimize the fidelities that we can obtain, we define a weight for each pixel that represents how likely it is that incident photons on that pixel came from

an atom being held at the relevant site [107], using the function $W(u, v)$, pictured in 2.28, rather than counting all incident photons within the 7×7 pixel box with the same weight. The photon histogram in Fig. 2.30a is based on photon-counting that is weighted on each site using this distribution.

The weights we choose are optimized via a quasi-Newton numerical method to maximize the imaging fidelity that we characterize with a model-free approach. This approach, as we describe in the next section, is agnostic of the photon distribution and relies solely on the consistency of imaging outcomes between three images. Thus, we believe that the imaging fidelity we quote is accurate and not artificially larger due to overfitting.

We then compare the signal obtained for each site and each image with a threshold to determine if an atom has been loaded. One can either use a global threshold for all sites, or characterize one for each tweezer site. We optimize the placement of the threshold in order to maximize the imaging fidelity of the array, as depicted in Fig. 2.30c.

To position the threshold and estimate the fidelity, we employ two complementary methods: an analytical model that predicts the shape of the imaging histogram by integrating the loss probability in a Poisson distribution, and a model-free approach that estimates the fidelity by identifying anomalous atom detection results in three consecutive images. The first method *infers* classification errors from the shape of the photon histogram while the second method *detects* errors directly; thus, the first method requires fewer samples to reach satisfactory accuracy, however it relies on a model. The inference method is compatible with any type of experimental runs while the direct detection method requires acquiring three consecutive images. Hence, we use the first method to position the binarization threshold in most experimental runs, as well as for site-by-site analysis; we use the second method to accurately estimate the fidelity with a single array-wide threshold and to benchmark the model-dependent approach.

2.6.2 Imaging fidelity

We distinguish atomic presence in the array with high fidelity. The wide separation between the peaks in Fig. 2.30a demonstrate that the presence of an atom (1) can be well distinguished from no atom (0).

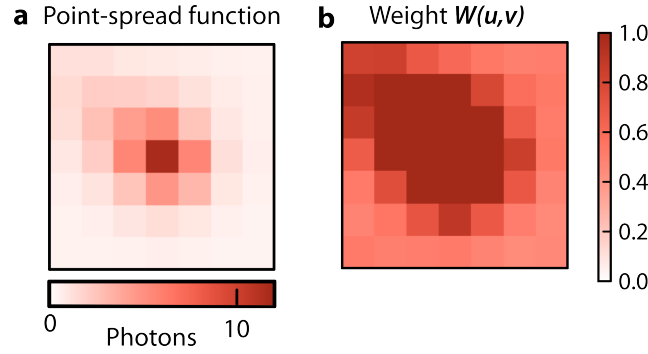


Figure 2.28: **Point-spread function and imaging weights** **a**, The point-spread function (psf) of collected photons, averaged over all sites that are above threshold in each of 16,000 images of the 11,998 tweezer array. The 7 pixel width corresponds to $6.3 \mu\text{m}$ in the tweezer focal plane. The waist radius of the psf is 1.7 pixels with a standard deviation of 0.2 pixels. **b**, The weight function $W(u, v)$ applied to each pixel of the 7×7 square-pixel box around each site. Here, u and v refer to the camera pixel coordinates centered on a given site. The weight function is chosen to optimize imaging fidelity based on the model-free characterization.

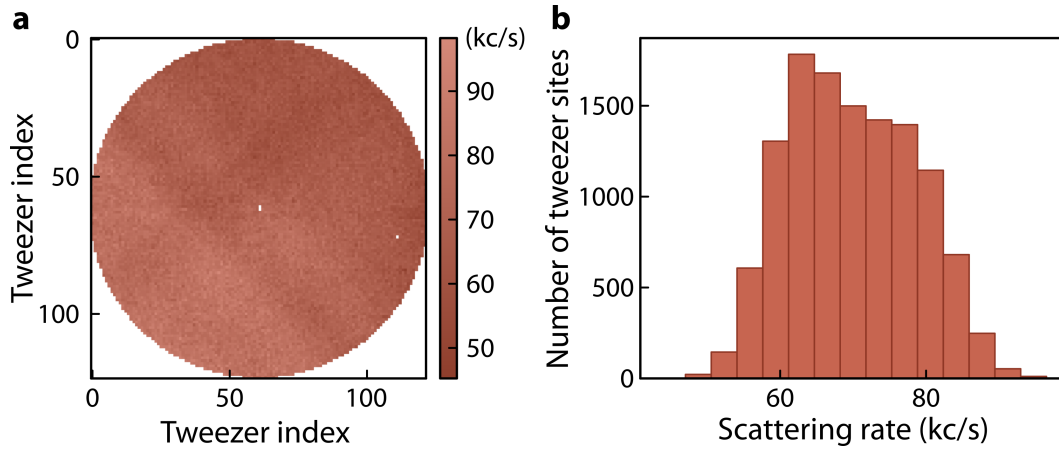


Figure 2.29: **Scattering Rate Uniformity** **a**, Map of scattering rates across the 11,998 tweezer array. Uniformity in scattering rate allows a global threshold to be used across the array to determine atom presence, rather than requiring a site-by-site approach for high fidelity. **b**, A histogram of number of scattering rates across all the sites in the array.

Model-free imaging characterization

We characterize the imaging fidelity, defined as the probability of correctly labeling atomic presence in a site, with a model-free approach, where no assumption is made about the photon distribution from Fig. 2.30a. This approach considers that the imaging survival and fidelity is identical for three successive images [33, 81] and identifies anomalous series of binary outputs. For instance, $0 \rightarrow 1 \rightarrow 1$ would point to a false negative event in the first image, while $1 \rightarrow 1 \rightarrow 0$ could be due to atom loss during the second image or a false negative event in the third one. This approach allows us to precisely decouple inherent atom loss from false negatives or positives. From this we find an imaging fidelity of 99.99374(8)% (note that in reality, we take 4 images, and use images 2-4 since we find the first image has slightly lower fidelity and survival probability). Crucial to this result is the homogeneous photon scattering rate across the array (Fig. 2.29) and the consistency of the point-spread function across the array (waist radius of 1.7 pixels with a standard deviation of 0.2 pixels, Fig. 2.28). Consistent imaging parameters across the array further are evidenced in that we find that treating each site with an individual threshold only marginally improves the imaging fidelity to 99.9939(1)%.

To implement this model-free approach, we estimate the probability $\tilde{P}_{x_1x_2x_3}$ of the presence of an atom in three images being $x_1x_2x_3$, where x_i is 1 if there is an atom and 0 if there is none. These probabilities are given as

$$\begin{aligned}\tilde{P}_{111} &= S^2 F, \\ \tilde{P}_{110} &= (1 - S)SF, \\ \tilde{P}_{100} &= (1 - S)F, \\ \tilde{P}_{000} &= 1 - F.\end{aligned}\tag{2.74}$$

where S is the survival probability during imaging and F is the initial filling fraction. We can then estimate the probability of *detecting* $y_1y_2y_3$ as

$$P_{y_1y_2y_3} = \sum_{x_1x_2x_3} P(y_1|x_1)P(y_2|x_2)P(y_3|x_3)\tilde{P}_{x_1x_2x_3}.\tag{2.75}$$

The conditional probabilities on the detection categorization given the true atomic presence are $P(1|1) = \mathcal{F}_1$, $P(0|1) = 1 - \mathcal{F}_1$, $P(1|0) = 1 - \mathcal{F}_0$, and $P(0|0) = \mathcal{F}_0$, where $\mathcal{F}_1$ and $\mathcal{F}_0$ represent the probabilities of properly characterizing atom presence (true positive) and atom absence (true negative) respectively. We use the method

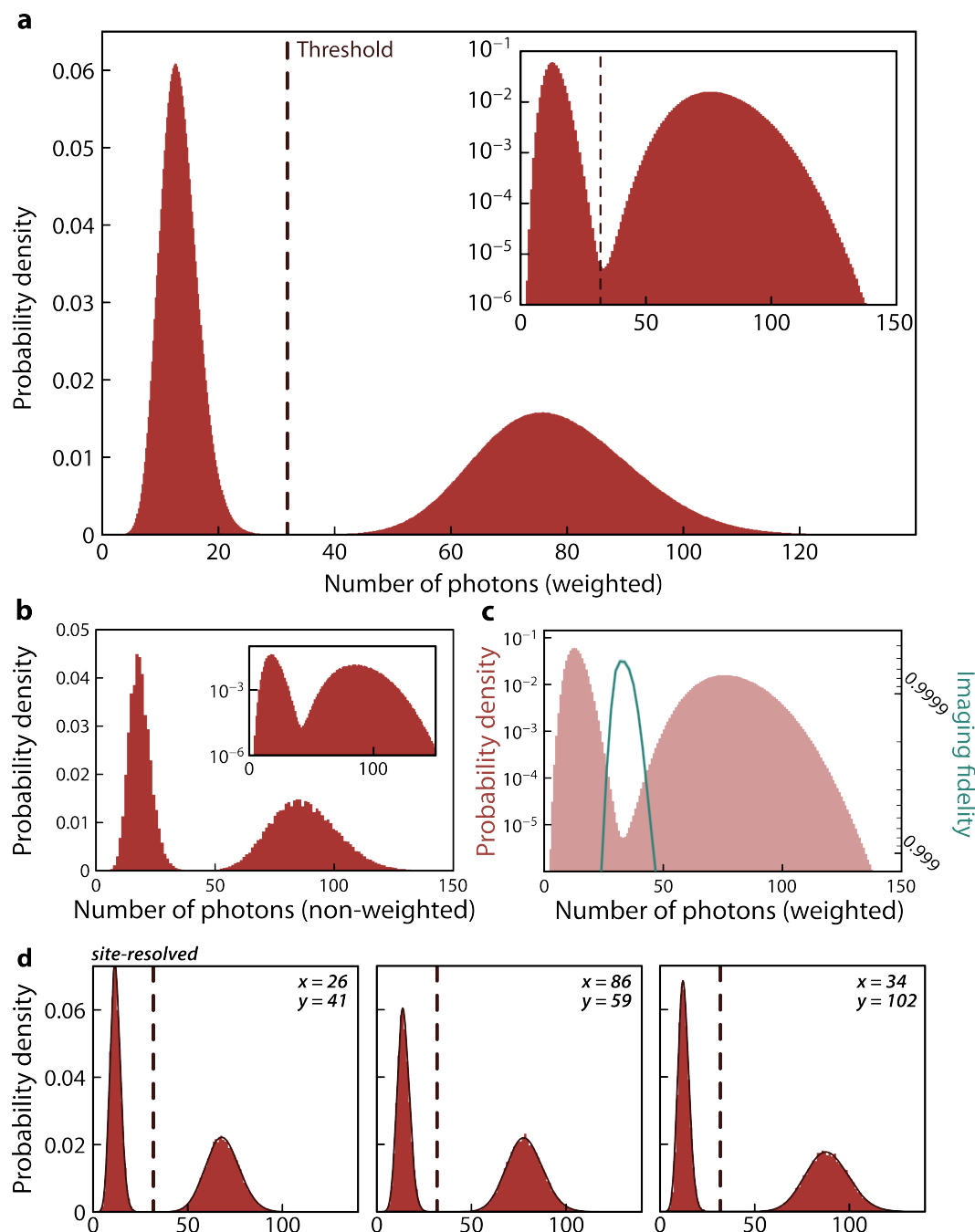


Figure 2.30: **Imaging Histograms** **a**, Weighted photon distribution histogram accumulated over 16,000 images, with threshold designated with which we determine presence (right of threshold), or absence (left of threshold) of a single atom. The inset is the same histogram in log-scale so that the overlapping tails that contribute to errors can be visualized. **b**, Unweighted photon distribution histogram with log-scale inset. **c**, Imaging fidelity overlay with the log-scale histogram, designating the fidelity as a function of where the threshold is positioned relative to the number of weighted photons. **d**, Individual sites are selected at random and compared with a global threshold.

of least squares to minimize the difference between the experimental frequencies of bitstrings $y_1 y_2 y_3$ and the $P_{y_1 y_2 y_3}$ by tuning the four parameters F , S , $\mathcal{F}_0$ and $\mathcal{F}_1$. The imaging fidelity is then defined as

$$\mathcal{F} = F\mathcal{F}_1 + (1 - F)\mathcal{F}_0. \quad (2.76)$$

The array-wide binarization threshold is chosen to maximize the imaging fidelity (Fig. 2.30c). Using this method, we find an imaging fidelity $\mathcal{F} = 0.9999374(8)$, with a false positive probability $1 - \mathcal{F}_0 = 7.01(8) \times 10^{-5}$ and a false negative probability $1 - \mathcal{F}_1 = 5.5(1) \times 10^{-5}$; we find the survival to be $S = 0.999864(2)$, slightly lower than the steady-state imaging survival probability measured by repeated imaging in Fig. 2.32.

Model-based imaging

As has been done in the past [102, 108], we incorporate a Poissonian model for photon collection for each of the background and the single-atom peak. However, we also recognize that the loss of an atom during the time-frame of imaging leads to a deviation from this purely Poissonian model, and results in a larger ‘bridge’ feature in the histogram between the background peak and the single-atom peak.

Thus, we introduce an analytical model that predicts the shape of histogram, which we call the “lossy Poisson model”¹⁰. We fit six parameters: the initial filling fraction (before the first image) F , the mean number of photons collected from the background light λ_0 and the atoms λ_1 , the respective broadening from an ideal Poisson distribution r_0 and r_1 , and the pseudo-loss probability L . The meaning of the parameters are further described below.

We first derive this model in the absence of broadening, from an ideal Poisson distribution. We are interested in the photon distribution given that there is no atom at a given site at the beginning of imaging $P(N = n|0)$ and the photon distribution given that there is an atom at this site at the beginning of imaging $P(N = n|1)$, where N is the number of photons collected. For the background photon distribution, we simply assume a Poisson distribution: $P(N = n|0) = e^{-\lambda_0} \lambda_0^n / n!$.

For the single-atom photon distribution, where loss can impact the resulting distribution, we derive an expression by considering a loss-rate model where each photon

¹⁰We note that this new model was developed by Elie Bataille. Most items in this thesis are work done in close collaboration with others, but given the novelty here, we note this particular individual contribution. Details are reproduced from [59]

collection event (occurring with probability $\lambda_1 dt$) imparts a loss probability L/λ_1 . By integrating over $t \in [0, 1]$ in the system of equations that describes the evolution of the joint distribution of atom presence and photon count, we find the distribution given that one atom was initially present,

$$P(N = n|1) = \frac{(\lambda_0 + \lambda_1 - L)^n e^{-(\lambda_1 + \lambda_0)}}{n!} + \frac{L}{\lambda_1} \frac{e^{\frac{\lambda_0 L}{\lambda_1 - L}} \left(1 - \frac{L}{\lambda_1}\right)^{n-1}}{(n-1)!} \quad (2.77)$$

$$\times \left[\Gamma\left(n, \frac{\lambda_0}{1 - L/\lambda_1}\right) - \Gamma\left(n, \lambda_1 + \frac{\lambda_0}{1 - L/\lambda_1}\right) \right].$$

Here, Γ represents the upper incomplete gamma function. The real loss probability during imaging is then given by $\tilde{L} = 1 - e^{-L}$. This equation illustrates the two mechanisms that limit the imaging fidelity in experiments with single-atom imaging. The first mechanism, represented by the first term in the r.h.s. of the equation, manifests as a Gaussian/Poissonian overlap between the two peaks of the photon distribution, reflecting our ability to record a significant photon count above the imaging noise floor. Finite scattering rate, limited photon collection efficiency, background light leakage from the imaging beams or the ambient light, and readout noise from the camera contribute to this limitation. The other mechanism that limits imaging fidelity is loss of atom during imaging. This manifests as a characteristic ‘bridge’-like feature, and is represented by the second term in the r.h.s. of the above equation. The probability density in the bridge is small but finite across a wide range of photon counts between the two peaks of the imaging histogram [109].

The overall photon probability distribution is then given by $P(N = n) = FP(N = n|1) + (1 - F)P(N = n|0)$. For practical purposes we empirically include a broadening of the Poisson distribution by writing

$$P(N = n) = F \frac{P(N = n/r_1|1)}{r_1} + (1 - F) \frac{P(N = n/r_0|0)}{r_0} \quad (2.78)$$

and by effectively considering non-integer photon numbers (by replacing factorials with the gamma function). For large n this amounts to considering a Gaussian distribution for either of the two peaks, but with the added benefit of including the loss through a physically-motivated derivation using a Poisson process.

Atom imaging as a Poissonian process

Describing photon counts from an empty or a single-atom filled site as a Poissonian process arises in the following manner. We image for a fixed amount of time t , and the number of photons collected during that time from a particular site (occupied or unoccupied) is discrete. To represent this as a Poissonian distribution then, the process must also have a known constant mean rate of emission λ (for each of the occupied and unoccupied cases), and each photon must be collected independently of the time since the previous photon was collected.

Both background photons and single-atoms without loss can be modeled in this manner. Background photons arise from consistent sources—inherent to the camera, or incident due to leakage of other lasers in the experiment. Given that each experimental sequence is repeated identically, this can be taken as a constant mean background rate, which describes the center of the background peak.

Single atoms in a loss-less scenario have a steady photon emission rate that is given by the product of the decay rate of the excited state times the probability that the atom is in the excited state, which reaches a steady state: $R_{sc} = \Gamma \rho_{ee}$

Though the spacing of photons being collected is limited by the time it takes for an atom to decay before another photon can be emitted, this timescale is very fast in comparison with the imaging time. Additionally, only $\sim 2.7\%$ of the photons scattered by the atoms are actually incident on our camera, and the quantum efficiency of the camera at 852 nm is 44%. These ensure that each photon emission can be modeled as independent.

Thus, modeling atom imaging histograms with overlapping Poissonians is physically motivated, however does not account for loss during imaging which motivates the deviation that we capture in the so-called ‘lossy Poissonian model’.

In this model the true negative probability is given by $\mathcal{F}_0 = \int_0^T P(N = n|0)dn$, where T denotes the threshold; and the true positive probability, by $\mathcal{F}_1 = \int_T^\infty P(N = n|1)dn$. Finally the imaging fidelity can be estimated as $\mathcal{F} = F\mathcal{F}_1 + (1 - F)\mathcal{F}_0$ and the op-

timal threshold T can be found by maximizing the fidelity. We find that this model performs well when predicting the shape of the histogram site-by-site (Fig. 2.30d), but fails when the distribution of the background or atom photons in the array is non-Gaussian.

Finally, we can inject the model-free survival probability that we measured (S) into the lossy Poisson model to increase its accuracy (trying to extract the loss directly from the lossy Poisson model would indeed be inaccurate, since losses appear as a small tail feature between the two peaks of the imaging histogram). Using this approach, and fitting each site independently, we find an average imaging fidelity of 99.992(1)%, in reasonable agreement with the model-free imaging fidelity. By setting the atom loss to zero while keeping the five other fit parameters constant for each site, we can estimate a hypothetical imaging fidelity in the absence of atomic loss of 99.999(1)%. This analysis also illustrates that fitting the imaging histogram with a Gaussian or Poissonian model without including losses leads to overestimating imaging fidelities [102].

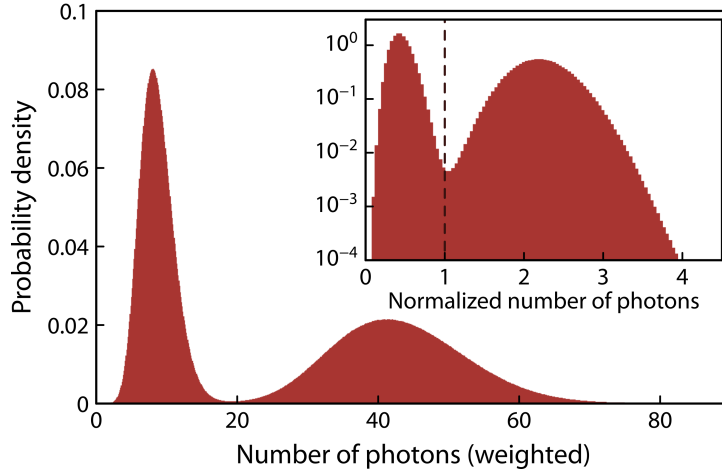


Figure 2.31: **20 ms Imaging** Imaging histogram for 20 ms imaging time. The weight function used is the same as the one shown in Fig. 2.28. Using model-free imaging characterization, we find an imaging fidelity of 99.9571(4)% and a survival probability of 99.176(1)%. Inset: log-scale histogram where the number of photons for each site is rescaled by the threshold for this site.

Fast Imaging

In addition to the high-fidelity (and high-survival) imaging demonstrated and characterized in Fig. 2.30 (and Fig. 2.32), we show in Fig. 2.31 imaging results acquired with an imaging time of 20 ms. This timescale is on the shorter side for high-fidelity

non-destructive imaging with alkali atoms, although even faster imaging has been demonstrated with other species [110] or lower fidelity [111]. It is worth noting that long, high-fidelity images would be used particularly for initial sorting, to characterize the presence of atoms before and after rearrangement to reach high preparation fidelity. Other operations for a quantum processor would likely be accomplished using different techniques, such as fast destructive readout.

Notably, this imaging data was acquired with a PGC detuning of $\Delta = -9.5\Gamma$, rather than -15.5Γ . The equilibrium atomic temperature scales inversely with the detuning in this regime, so we do expect more loss, which we expect to impact both the fidelity and the survival [112]. We detune closer however for a shorter imaging time in order to collect a similar number of photons— the scattering rate increases by a factor of $(\frac{15.5}{9.5})^2 \sim 2.66$, while the time that we image for decreases by a factor of 4. So the number of photons that we collect is about 66% compared with what we presented previously. This lower photon number also impacts the fidelity with which we distinguish between a single atom and no atom.

In this configuration, we measure an imaging fidelity and survival probability of 99.9571(4)% and 99.176(1)%.

2.6.3 Imaging survival

As we have discussed, there is some probability of losing an atom from the trap during imaging due to heating that the atom incurs from scattering photons. Losing an atom while imaging can result in incorrectly assuming that an atom is present when it has been lost. This results in defects in the preparation of an array to be used for quantum computation or simulation. For a given atom imaging survival for a trap depth and beam configuration, the probability of not losing an atom in the array during imaging drops exponentially with the number of atoms in the array. Thus, this number is important to characterize to understand the usefulness of our large array.

We accurately characterize the imaging survival probability, without assuming any parameters, by performing 80-ms repeated imaging (at -15.5Γ detuning) up to 1,000 times, after which $\sim 90\%$ of initially loaded atoms still survive (Fig. 2.32a). We fit this decay to $p(N) \propto p_1^N$, and find a steady-state imaging survival probability of $p_1 = 99.98952(1)\%$, mostly limited by vacuum lifetime. This, to the best of our knowledge, surpasses prior studies reporting record steady-state imaging survival using single alkaline-earth metal [113] and alkali-metal [114] atoms in optical

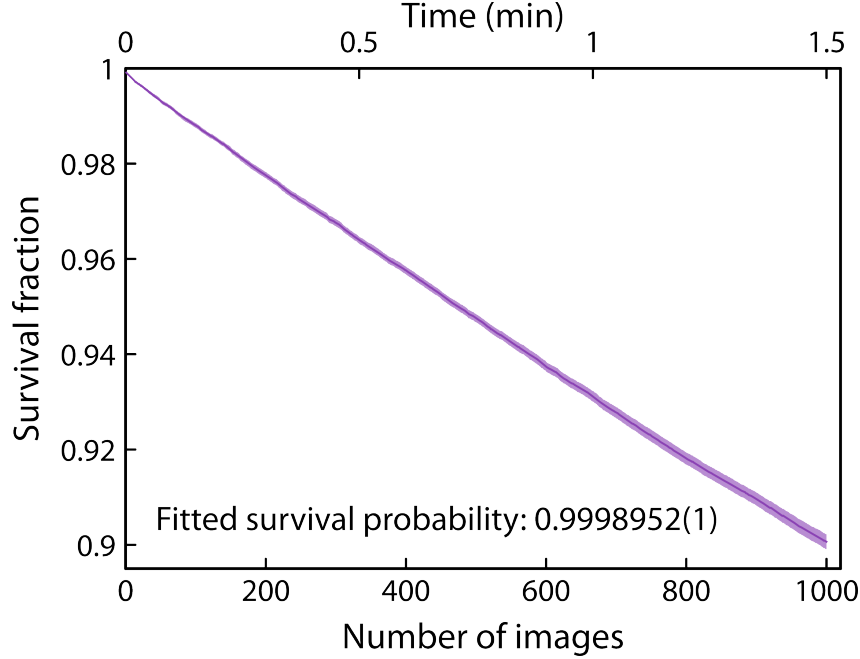


Figure 2.32: **Imaging Survival** Array-averaged survival fraction after many successive images. Between each image, we hold the atoms for 10 ms, without applying any cooling beams. We fit the data with $p(N) \propto p_1^N$, where $p(N)$ is the survival fraction after imaging N times. From the fit, we find a steady-state imaging survival probability of $p_1 = 0.9998952(1)$. The light purple fill shows the estimated 68% confidence interval.

tweezers.

These results, and the uniformity of imaging survival across the array, as can be seen in the map and histogram in Fig. 2.33, enable low-loss, high-fidelity detection of single atoms in large-scale arrays, crucial components for the practical use of the system. We note that the histogram is plotted in log scale in order to show that the site with the lowest survival still has 99.66% survival probability, and the average of the site-resolved imaging survival is 99.985%. These data were obtained with the 3-image model-free fidelity measurement, and are used to determine the survival that we use for updating the lossy-Poissonian model.

As an example of the practical implications of this array-wide survival and how the characterization is affected by the choice of threshold, we plot the probability of initializing an experiment in a defect-free manner, assuming loss-less rearrangement, and only accounting for atom loss during imaging, as a function of threshold choice (Fig. 2.33b). We additionally plot the fraction of atoms that are excluded due to the threshold placement—the defect-free probability is the likelihood of initializing

an experiment without loss due to imaging of the remaining atoms that are not excluded.

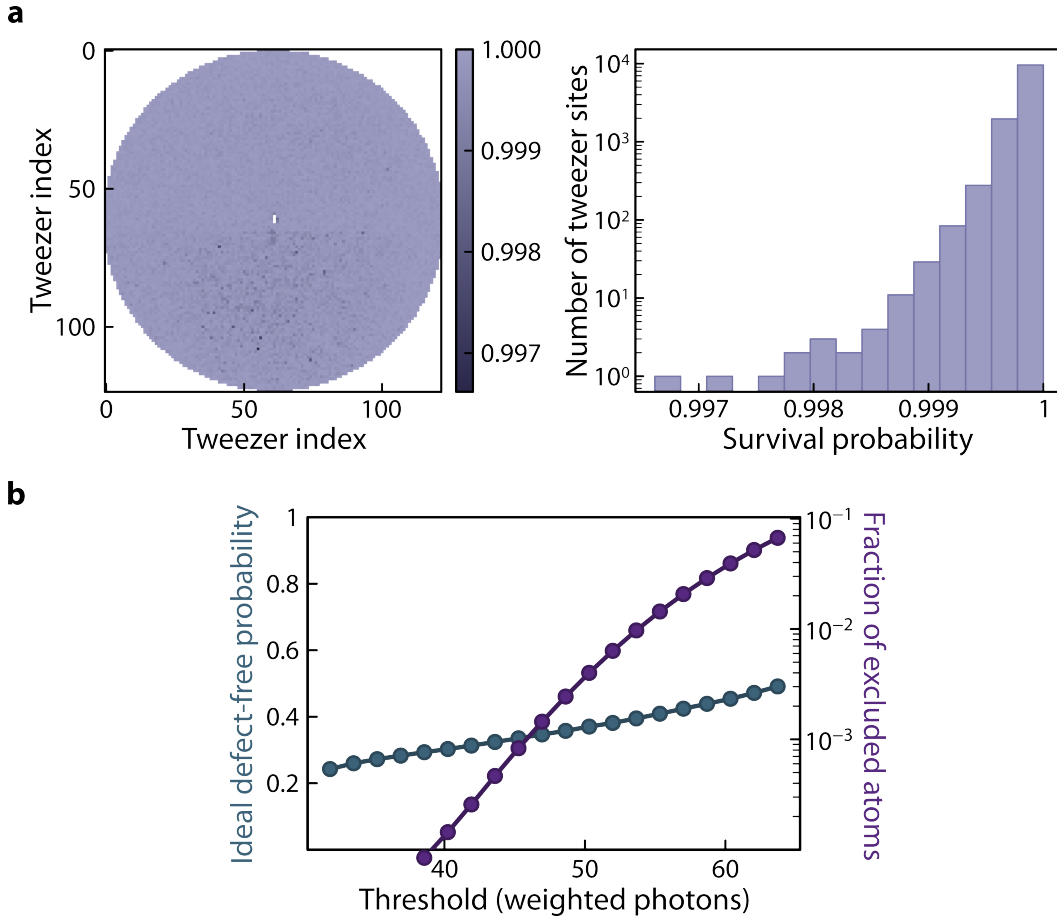


Figure 2.33: **Imaging survival** **a**, Map and histogram of the imaging survival probability per site across the tweezer array, as characterized using the three image analysis of data from 16,000 iterations. Note that the vertical axis of the histogram figure is plotted on a log scale. The mean of site-resolved imaging survival probability is 99.985%, and the minimum value 99.66%. **b**, Predicted upper bound on the probability of detecting a defect-free array after an ideal rearrangement sequence (estimated as $p(1|1)^n$ where n is the number of atoms in the first image), limited by imaging survival and false positives. The threshold in the first image can be displaced to reduce false positives, at the cost of excluding some atoms. Note that we may ignore the issue of false negatives in the first image, since we can always physically eject residual atoms in sites that are determined to be negative.

2.7 Vacuum

The probability of losing no atom in a tweezer array during imaging and due to finite vacuum lifetime both decrease exponentially with the number of atoms in an array, making these crucial metrics to optimize for large-scale array operation. The

vacuum-limited lifetime, in particular, sets an upper bound on the duration during which operations can be executed without loss of an atom in a given experimental run. This can, for example, be applied as an upper limit on the fidelity with which one can achieve a defect-free array via atom rearrangement [30, 31]. Vacuum loss arises from collisions of background atoms with atoms that are being held in tweezers.

We thus particularly work to design an experiment that optimizes vacuum lifetime and minimizes background pressure. We additionally take care with our baking procedure to eliminate leaks and minimize potential out-gassing sources (see Appendix A). These optimizations result in a vacuum-limited lifetime measurement of around 23 minutes, as we describe below.

2.7.1 Vacuum Design

A views of a schematic of our vacuum system is shown in Fig. 2.34, with images of the actual setup also pictured in Fig. 2.37. After the initial chamber assembly and multi-round baking process (see Appendix A), we fire two titanium sublimation pumps (TSPs), mounted such that every surface except the rectangular portion of the glass cell and the interior of the ion pump are covered by line-of-sight sputtering. This creates a vacuum chamber in which essentially every surface is pumping. We do not find it necessary to re-fire the TSPs in order to maintain the vacuum level that we measure. We additionally maintain ultra-high vacuum conditions with an ion pump, connected to the primary chamber via a 45° elbow joint. The secondary, science chamber consists of a rectangular glass cell (JapanCell) optically bonded to a 24-cm long glass flange (also sputtered by the TSP) that connects to the primary chamber. From lifetime measurements of tweezer trapped atoms and collisional cross-sections available in literature [115], we estimate the pressure in the glass cell to be $\sim 7 \times 10^{-12}$ torr, consistent with vacuum simulations using the MolFlow program [116].

2.7.2 Vacuum Calculations

Collisions of background atoms with atoms being held in tweezers lead to vacuum-related loss. We bake the chamber in order to minimize the out-gassing from the elements of the vacuum that can contribute to this background gas (as described in detail in Appendix A). However, there remain a few prominent sources of background atoms. One is from the source of Cs itself—there is a large source of Cs atoms at a much higher pressure (10^{-7} to 10^{-6} torr) that is line-of-sight to the tweezer array,

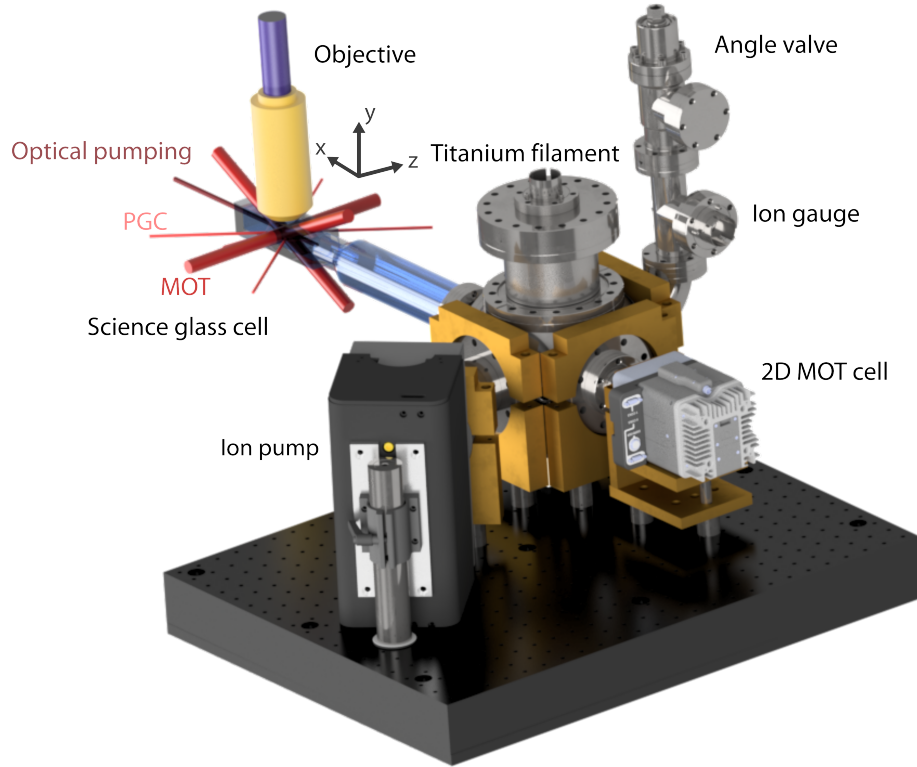


Figure 2.34: **Vacuum chamber** A view of a CAD model of our vacuum chamber with relevant beams labeled as well as relevant vacuum parts and pumps. The titanium sublimation pumps were particularly optimized in order to coat the glass-to-metal transition and provide as much pumping surface area as possible. The 2D MOT source cell, ion pump, and top titanium sublimation pump are visible. At the glass cell, bigger red beams are MOT beams, smaller bright red beams are PGC beams, and the dark red beam parallel to the cell axis is for state preparation.

only separated by a differential pumping aperture of 0.75mm diameter. Additionally, while out-gassing rates decrease with baking, the glass cell is the closest to the atoms, and it is not coated with titanium, which acts as a pumping surface (as other parts of the chamber are). So we additionally consider out-gassing from the glass cell of predominantly H_2O as well as He permeation through the glass cell. We consider He permeation, and not the permeation of other molecules such as H_2 , given that the size and shape of the particle allows it to penetrate glass in ways that are forbidden

for other particles in the background atmosphere.

The lifetime (τ) of a Cs atom in a trap, limited by collisions from background atoms with mass m , at temperature T , with partial pressure P , can be given as

$$\tau = \frac{1}{\sigma P} \left(\frac{\pi m k_B T}{8} \right)^{1/2} \quad (2.79)$$

where σ is the collisional cross-section of the background atom with Cs [117].

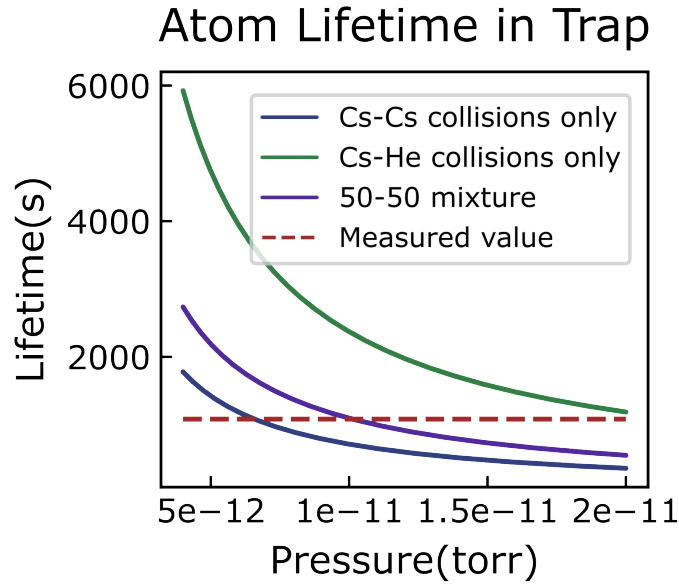


Figure 2.35: **Vacuum-limited Lifetime Prediction** The lifetime of an atom in a tweezer is plotted as a function of the pressure from background gases. We consider the cases where the background pressure is only from Cs-Cs collisions (blue) and where the background pressure only arises from Cs-He collisions (green). We finally consider the case where both have equal partial pressures contributing to the total pressure (purple). We plot alongside these calculations, the measured value (dashed red) of 1374(60) seconds, and note that from MolFlow simulations, we expect that Cs-Cs collisions dominate, and that thus our lifetime measurement corresponds to pressure around 7×10^{-12} torr.

We plot the expected lifetime of a Cs atom in a trap, as a function of background pressure for the case where only Cs-Cs collisions contribute to the background pressure, for the case where only Cs-He collisions contribute to the background pressure, and for the case where they each equally contribute to the background pressure (see Fig. 2.35). To estimate the expected pressure contributions from Cs and He that we expect in our system (as well as for optimizing design parameters and conductance),

we implement simulations using MolFlow [116]. We include simulations (assuming room temperature) for Cs (line-of-sight to 2D MOT), for H₂O outgassing from the glass cell, and for He permeation through the glass cell separately and estimate the partial pressures in each scenario, as pictured in Appendix A (Fig. A.2). We additionally describe the details of assumptions that we make for pump rates and outgassing rates in Appendix A. From these simulations, we conclude that Cs-Cs collisions are dominant on our vacuum, contributing around $\sim 6 \times 10^{-12}$ mbar partial pressure, compared with $\sim 9 \times 10^{-13}$ mbar partial pressure from outgassed H₂O, and $\sim 4 \times 10^{-13}$ mbar partial pressure from He permeating through our 6.35 mm thick glass cell.

2.7.3 Vacuum Lifetime Measurement

We probe the vacuum-limited lifetime in our tweezer array using an empirically optimized cooling sequence consisting of a 10-ms 2D PGC cooling block every 2 s [118]. By fitting the exponential decay of the atom survival, we find a $1/e$ lifetime of 22.9(1) min (Fig. 2.36). This is a significantly long timescale compared with state-of-the-art room-temperature atomic experiments, and within a factor of five of the longest reported lifetime in a cryogenic apparatus [118]. The result indicates that the probability of losing a single atom across the entire array remains under 50% during 100 ms, a relevant timescale for dynamical array reconfiguration and quantum processor operation.

We additionally measure the survival in our atom array after implementing the cases of continuous PGC cooling while holding the atoms, or in another case, no cooling while we hold the atoms. In the case of continuous cooling, we observe a lifetime of 17.7(2) min, and in the case of no cooling, we find a $1/e$ lifetime of around 2.2 min. We anticipate that the loss when cooling is not implemented is due to scattering from the trapping light, as well as tweezer noise that can cause heating, arising from sources such as laser relative intensity noise (RIN), SLM flicker noise, and tweezer pointing stability. In the case of continuous cooling, we anticipate that the equilibrium temperature is reached fairly quickly with a short block of cooling, and that further cooling can stochastically cause atom loss since there is non-zero probability to scatter photons that cause an increase in temperature, though on average the 2D PGC configuration should cool.

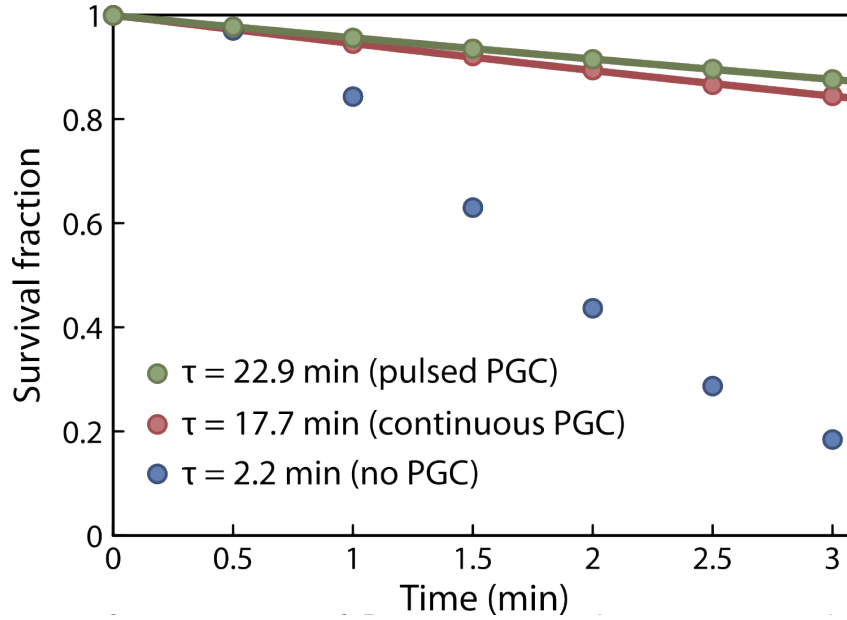


Figure 2.36: **Vacuum-limited Lifetime Measurement** Array-averaged survival fraction as a function of hold time is plotted. Three experiments are shown in the figure: pulsed cooling, continuous cooling, and no cooling. The green markers show data with a 10-ms 2D PGC block applied every 2 s during the wait time (pulsed PGC). The red markers show data with 2D PGC block continuously applied during the wait time (continuous PGC), while the blue markers show the data without cooling during the wait time (no PGC). The error bars are smaller than the markers. We find a $1/e$ lifetime of around 2.2 min without cooling. When the pulsed PGC block is applied, by fitting the data with $p(t) \propto \exp(-t/\tau)$, we find a vacuum lifetime of $\tau = 22.9(1)$ min. When the 2D PGC is applied continuously, we obtain $\tau = 17.7(2)$ min.

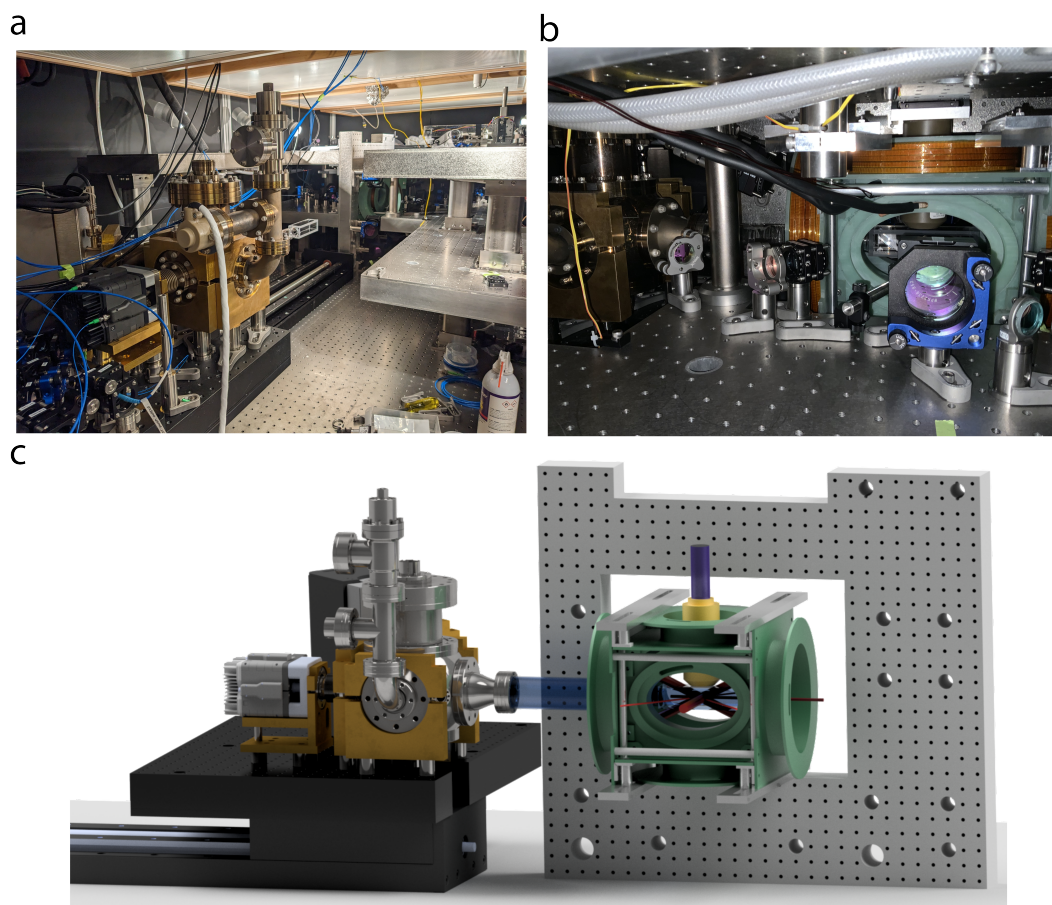


Figure 2.37: Pictured are photos of our vacuum chamber (a,b) as well as a CAD rendering of the design (c). **(a)** A view of the chamber as well as the 3-layer breadboard system on which our experiment is built. Our chamber is on a linear motion stage such that it can be retracted from the science position for the purpose of baking, or beam or objective alignment. **(b)** A view of the science cell. In green are high-field coil holders for water-cooled coils, allowing for up to around 1G/A magnetic field in each direction (see Appendix B). **(c)** A CAD rendering of the vacuum chamber, including the linear stage, coil holders and vertical breadboard.

Chapter 3

FROM 6100 ATOMS TO 6100 QUBITS

We have presented thus far a record number of single atoms trapped in optical tweezers, along with data demonstrating the uniformity of the traps holding the atoms, and effective cooling and long vacuum lifetimes alongside high fidelity imaging and survival— all necessary for preparing a defect free array with high fidelity, and for allowing longer timescale array operations.

In order to demonstrate that these 6100 atoms will in fact be useful for quantum science, we now look at their qualities when used as qubits. We first introduce the hyperfine qubit (3.1) discuss coherence times (3.2), introducing methods of dynamical decoupling for extending coherence times. We then describe randomized benchmarking (RB) for single-qubit gates across the array and present data for single-qubit RB on our 6100 qubits (3.3). We demonstrate record long coherence times for alkali atoms in optical tweezers and high gate fidelities with global microwave operations.

3.1 The hyperfine qubit

Key to recent progress in quantum computing and metrology with neutral atoms is the ability to encode a qubit in long-lived states of an atom, such as hyperfine states [51, 52], nuclear spin states [110, 119–121], or optical clock states [38, 39]. In cesium atoms, the hyperfine ground states ($|F = 3, m_F = 0\rangle \equiv |0\rangle$ and $|F = 4, m_F = 0\rangle \equiv |1\rangle$), which are separated by the *clock frequency*¹ of ~ 9.19 GHz, provide such a subspace for storing quantum information. Furthermore, entanglement via Rydberg interactions can be readily transferred to this qubit to realize high-fidelity two-qubit gates [34]. We demonstrate the storage and manipulation of quantum information in a large-scale atom array by measuring the coherence times and global single-qubit gate fidelity using a microwave horn to drive the hyperfine transition (Fig. 2.10a).

Preserving the coherence of a quantum system as it is scaled up is a known challenge across platforms for quantum computing and simulation [17]. This difficulty persists

¹This transition is called the *clock transition* because it is insensitive to first order to magnetic fields. The clock transition in cesium has been used to define the SI second since 1967 [122]. Precision provided by atomic clocks has been necessary to the development of technology such as global navigation satellite systems.

even for neutral atoms, albeit at a lower level, due to residual interactions with a noisy and inhomogeneous electromagnetic environment, particularly with the tweezers themselves. Thus, as discussed in the last chapter, we have chosen to trap in far-off-resonant optical tweezers to preserve coherence, since at constant trap depth the differential light shift of the hyperfine qubit decreases as $1/\Delta_{\text{tweezer}}$ and the scattering rate as $1/\Delta_{\text{tweezer}}^3$, where Δ_{tweezer} is the tweezer laser detuning relative to the dominant electronic transition [68, 123, 124].

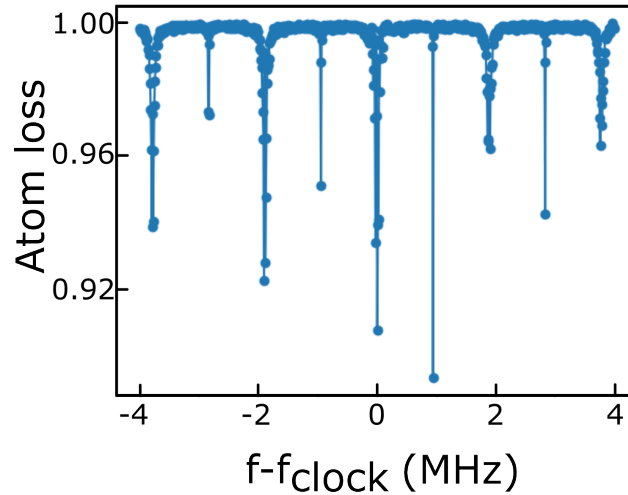


Figure 3.1: **Microwave spectroscopy** We map out hyperfine transitions using microwave spectroscopy after pumping broadly into the $F = 4$ hyperfine manifold. Since we push out atoms in the $F = 4$ manifold at the end of the measurement (such that it corresponds to loss), the dips represent frequencies at which population is transferred to the $F = 3$ manifold. The center transition is $m_F = 0 \rightarrow m_F = 0$, and we determine a 2.7G magnetic field from the hyperfine state separation. We note that the broader transitions are driven by linear polarization ($m_F = 0 \rightarrow m_F = 0$, $m_F = \pm 1 \rightarrow m_F = \pm 1$, $m_F = \pm 2 \rightarrow m_F = \pm 2$), and the alternating more narrow resonances we believe are transitions driven with remaining components of circular polarization in the RF field.

In order to drive the hyperfine qubit at the right frequency, we first implement microwave spectroscopy. We can map out the hyperfine state energies, and thus the transition energy between the clock states f_{qubit} by implementing a broad scan of the microwave frequency, after repumping the atoms to the $F = 4$ state without implementing state preparation. We thus find the $m_F = 0 \rightarrow m_F = 0$ transition (center dip in Fig. 3.1) and additionally characterize the magnitude of B-field that we are applying via the frequency difference between adjacent hyperfine states. After

a more fine-tuned scan, we can then implement state-preparation via selection-rule dependent optical pumping (see section 2.3.3) and a Ramsey interferometry experiment to measure the resonant qubit frequency.

What is Ramsey interferometry?

To measure the difference between the RF drive frequency that we are using and the qubit frequency (with light shift), we can implement a Ramsey interferometry experiment [125].

First a $\frac{\pi}{2}_y$ pulse is applied to the qubits, prepared in state $|1\rangle$, such that the state is now $|\psi\rangle = \frac{1}{\sqrt{2}}(|0\rangle + |1\rangle)$.

Then the atoms evolve for an *interrogation time* τ while the drive is off ($\Omega = 0$), during which time the atoms evolve under the Hamiltonian $H = \frac{\hbar\Delta}{2}\sigma_z$, where Δ is the difference between the qubit frequency and the drive frequency.

Thus, the time-evolution operator $U = e^{-iHt/\hbar} = e^{-i\Delta\tau\sigma_z/2}$, representing a rotation around the z-axis by $\theta = \Delta\tau$. Thus, after time τ , our atoms will be in state $|\psi\rangle = \frac{1}{\sqrt{2}}(e^{-i\Delta\tau/2}|0\rangle + e^{i\Delta\tau/2}|1\rangle)$, which gives a relative phase between the two states of $\Delta\tau$.

We apply a second $\frac{\pi}{2}_y$ pulse, which maps the atoms to the state:

$$|\psi_{final}\rangle = \frac{1}{2}(1 - e^{i\Delta\tau})|0\rangle + \frac{1}{2}(1 + e^{i\Delta\tau})|1\rangle \quad (3.1)$$

We then measure the population in state $|1\rangle$ with probability $P = |\langle 1|\psi_{final}\rangle|^2$, which gives us

$$P = \frac{1}{2}(1 + \cos \Delta\tau). \quad (3.2)$$

The resulting oscillations in average population as we scan τ , are called *Ramsey fringes*.

This qubit frequency measurement can be implemented at different trap depths in order to map the relationship between the trap depth and the differential clock shift (Fig. 3.2). We fit the linear relationship between trap depth (measured by depumping via the $F = 4 \rightarrow F' = 4$ transition) and the differential light shift measured via

a Ramsey interferometry experiment to find η , which we can use alongside our T_2^* measurement to characterize the temperature in the array. For the remainder of the coherence and single-qubit gate fidelity measurements, we adiabatically ramp down tweezers to a depth of $k_B \times 55 \mu\text{K}$ during microwave operation.

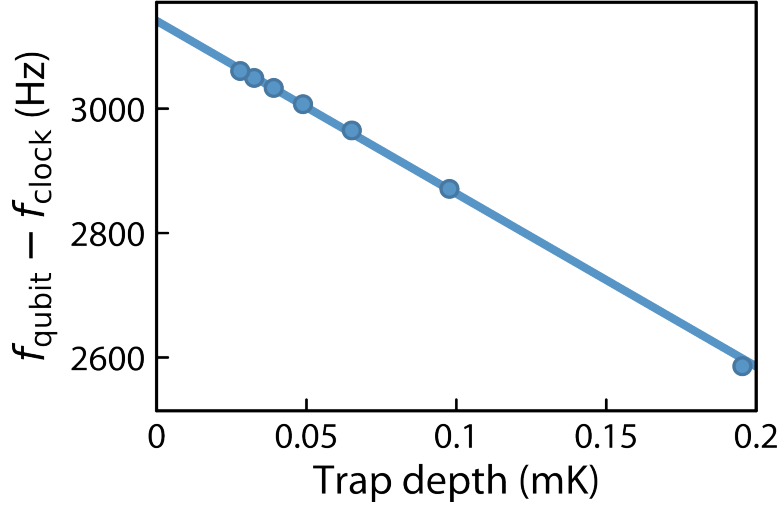


Figure 3.2: **Differential light shift - trap depth calibration** Estimation of η , the ratio of the differential polarizability of the hyperfine qubit to the electronic ground state scalar polarizability. The average qubit frequency is measured by Ramsey interferometry for different trap depths, and the slope is compared with the trap depth inferred from the light shift of the $F = 4 \leftrightarrow F' = 4$ D_2 transition from its free-space resonance.

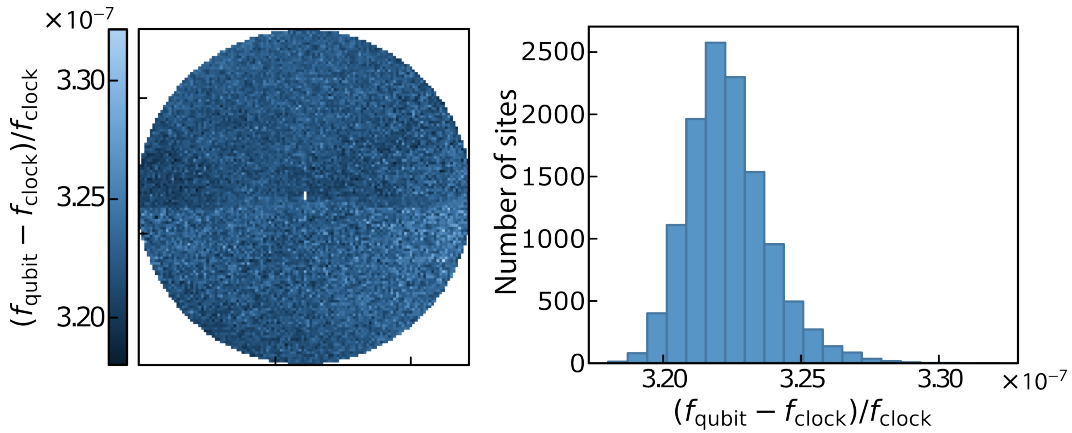


Figure 3.3: **Relative difference of qubit and clock frequency across the array** Relative difference of the qubit frequency with the cesium clock frequency $f_{\text{clock}} \equiv 9,192,631,770 \text{ Hz}$, measured by Ramsey interferometry. The standard deviation is 1.5×10^{-9} , or 14 Hz in absolute value.

We additionally measure the uniformity of the relative difference between the qubit frequency and the clock frequency across the array in a site resolved fashion (Fig. 3.3). This uniformity is important because differences in qubit frequencies result in individual qubits either over- or under-rotating when gates are implemented across the array. While such errors can be mitigated using robust gates or dynamical decoupling (see sections 3.2.3 and 3.3.4), there is a cost in time overhead, and not all operations can be easily made robust. We find a standard deviation across the sites in the array of 14Hz.

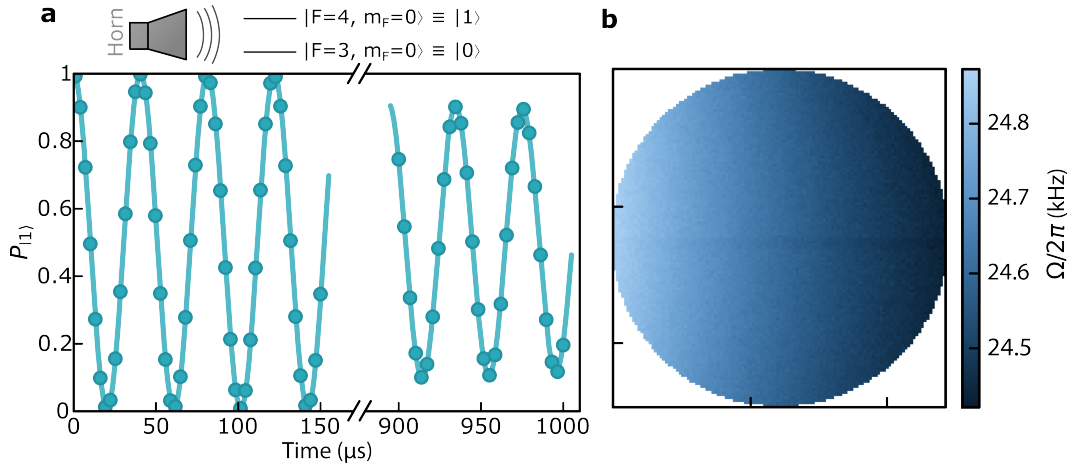


Figure 3.4: **Rabi oscillations** **a**, Array-averaged Rabi oscillations between the hyperfine clock states 0 and 1. The fitted Rabi frequency is 24.611(1) kHz. The observed decay after several hundred microseconds stems from the spatially-varying Rabi frequency across the sites in the array. **b**, Inhomogeneity of the Rabi frequency across the atom array. The reflection off a vertical metallic breadboard near the vacuum cell creates this spatial gradient orthogonal to the propagation axis of the microwaves. The Rabi frequency standard deviation is 0.5%.

3.1.1 Qubit Rabi oscillations

After characterizing the qubit frequency at the operational trap depth (55 μ K), we then drive the qubit on resonance using the microwave horn. The characteristic Rabi oscillations are pictured in Fig. 3.4, alongside site-resolved characterization of the Rabi frequency across the array. We attribute the decay of Rabi contrast after several hundred microseconds to this spatially-resolved inhomogeneity.

The fitted Rabi frequency is 24.611(1)kHz, and depends on the distance from the microwave horn to the atoms r , the input microwave power to the horn P , and the gain of the horn G in the following form,

$$\Omega = \frac{\mu}{\hbar r c} \sqrt{\frac{P G Z_0}{2\pi}} \quad (3.3)$$

where Z_0 is the impedance of free space and μ is the magnetic dipole moment.

We use the early-time Rabi contrast in order to characterize the fidelity with which we have prepared the state $F = 4, m_F = 0$, which we estimate to be about 99.1% fidelity.

3.2 Coherence Times

Coherence times provide a timescale for how long information can be stored in qubit states. We can differentiate between the timescale for which the classical population information can be stored (T_1 time), and how long the quantum phase information can be stored (T_2 time).

Different quantum platforms have vastly different timescales for information storage, and vary as to whether they are limited by T_1 times or T_2 times. For example, transmon qubits have coherence times on the scale of 1 ms, while neutral atoms have coherence times on the scale of 100-1000 ms. However, due to the slower operation of neutral atom processors this does not necessarily constitute higher fidelities, depending on the time-duration and fidelity of the required operations.

Maintaining coherence times while scaling is a challenge across platforms and must be demonstrated for a register of 6100 qubits to be actually useful for computation. For superconducting qubits, each manufactured qubit may have a significantly different coherence time, however atoms, particularly when trapped in far-detuned tweezers where the differential light shifts scale inversely with detuning when trap depth is constant, have the potential for higher uniformity.

We thus measure the coherence times of our qubits, and demonstrate techniques for increasing the bare T_2^* coherence time using dynamical decoupling techniques.

3.2.1 T_1 time

To measure the T_1 time of our qubits, we prepare them in the $F = 4, m_F = 0$ qubit state and hold them there, reading out the information at variable time in order to measure the decay of population. In order to de-convolute state preparation errors, we alternate applying a π pulse prior to readout and measure the difference between the case where a π pulse is applied and the case where it is not, for variable times.

We indeed observe long coherence times, measuring a depolarization time of $T_1 =$

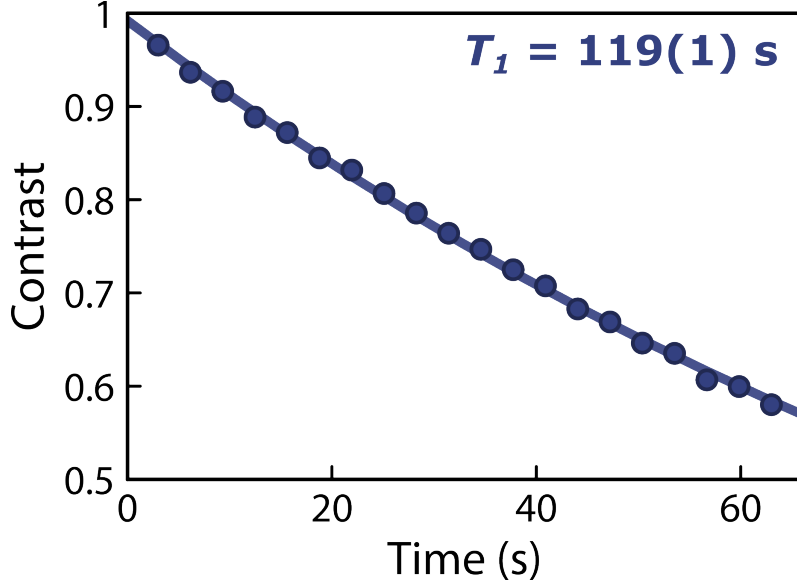


Figure 3.5: **T_1 coherence time** Measurement of the depolarization time T_1 . Atoms are initially prepared in $|1\rangle$. After a given time, the remaining population in $|1\rangle$ is measured, with or without a π pulse before the measurement. The population difference, conditioned by the application of the pulse, constitutes the T_1 contrast. This extended T_1 time is largely enabled due to the choice of far-detuned traps.

119(1) s (Fig. 3.5). This measurement is close to our prediction for scattering-limited T_1 times presented in Fig. 2.4 in section 2.2.1, and constitutes a record for alkali hyperfine qubits to our knowledge, which we attest to our use of far-detuned traps.

3.2.2 T_2^* time

In order to measure the free induction decay (T_2^*) or the phase decay without any spin-echo or dynamical decoupling, we utilize Ramsey interferometry. In particular, we choose to detune our drive frequency from the average qubit frequency by $\Delta = 2\pi \times 1$ kHz. This enables us to see Ramsey fringes which decay with the timescale given by the T_2^* time.

The envelope of the oscillation has a Gaussian decay² with a characteristic time $T_2^* = 14.0(1)$ ms. However, when considering each site individually we find an average $\langle T_2^{*(\text{site})} \rangle = 25.5$ ms with a standard deviation of 3.2 ms (in the per-site case

²The Gaussian decay arises from the signal being given as the integral over phases acquired by qubits multiplied by the probability of that phase being acquired. $S(t) = \int P(\Delta) e^{i\Delta t} d\Delta$. This represents a Fourier transform of the probability distribution of the relative difference between the qubit and drive frequency across the array. For the case that the distribution of differential light shifts for qubits across the array is approximately Gaussian, this decay shape will also be Gaussian.

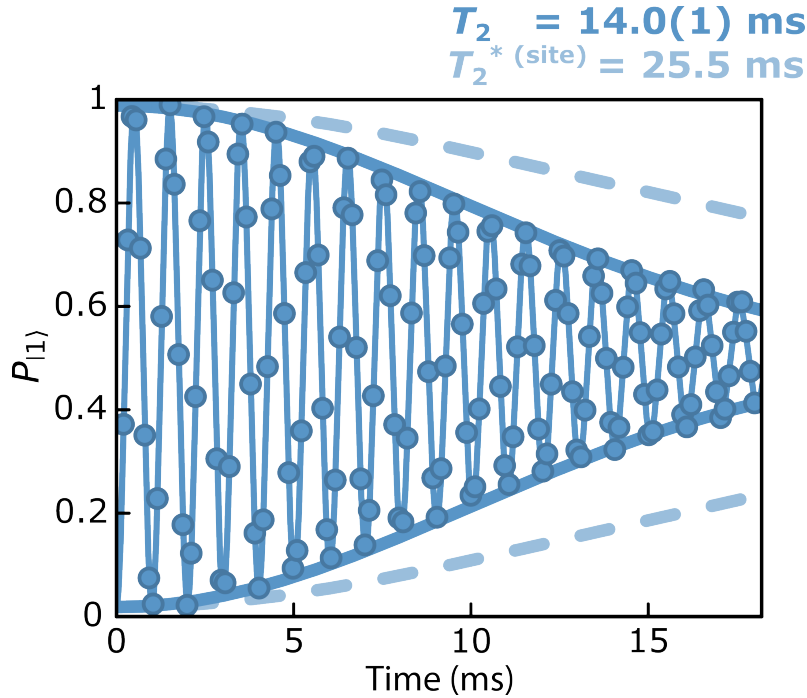


Figure 3.6: T_2^* **coherence time** Array-averaged Ramsey oscillations. During free evolution, the microwave drive field is detuned by 1 kHz, resulting in Ramsey oscillations. The characteristic decay time of these oscillations is $T_2^* = 14.0(1)$ ms from fitting the average signal of all atoms. The light blue dashed line shows the decay time $T_2^{*(\text{site})} = 25.5$ ms from fitting individual sites first and averaging the decay time afterwards.

we fit the oscillation decay with the dephasing decay function from Ref. [123]). This shows that dephasing across the array primarily occurs because of trap depth inhomogeneities (Fig. 2.25): assuming a Gaussian distribution of trap depth with a standard deviation δU , the qubit frequencies in the array also follow a Gaussian distribution, which results in an ensemble-wide dephasing time $T_2^{*(\text{inh})} = \sqrt{2}\hbar/(\eta \delta U)$ where η is the ratio of the scalar differential polarizability of the hyperfine ground states to their polarizability at the fine structure level [123]. On the other hand, finite atomic temperature limits the per-site dephasing time $T_2^{*(\text{site})}$. We observe an uneven distribution of T_2^* across the atom array (Fig. 3.7), with a significantly lower T_2^* measured for atoms trapped in tweezers at 1055 nm than for those trapped in the bottom half of tweezers at 1061 nm. This discrepancy could be due to worse optical aberrations in these areas that decrease the efficiency of polarization-gradient cooling, or due to different intensity noise profiles from the different fiber amplifiers or SLMs used on the two pathways. These data reveal that further investigation of noise sources specific to lasers or tweezer pathways could elucidate limiting factors

on coherence times in neutral atom arrays.

In order to relate T_2^* and trap depth inhomogeneity or atomic temperature, the parameter η can be calculated as the ratio of the differential light shift of the hyperfine states to the electronic ground state light shift (the trap depth), which we can calculate to be $\eta = 1.50 \times 10^{-4}$ ³. (At the few-percent accuracy level, it becomes important to account for higher-order processes [126], but such accuracy is not required here). We corroborate this value by experimentally measuring the differential light shift via Ramsey interferometry at different depths (Fig. 3.2), finding $\eta = 1.3(1) \times 10^{-4}$, in reasonable agreement with the theoretical value. This allows us to estimate the atomic temperature during microwave operation as [123]

$$T = \frac{2\hbar\sqrt{e^{2/3} - 1}}{\eta k_B \langle T_2^{*(\text{site})} \rangle} \approx 4.3 \text{ } \mu\text{K}, \quad (3.4)$$

assuming the temperature is sufficiently homogeneous to invert the fraction and the mean. This temperature may differ from the effective atomic temperature during other points of the experimental sequence that do not include the ramp-down and state preparation steps that may decrease and increase the temperature respectively.

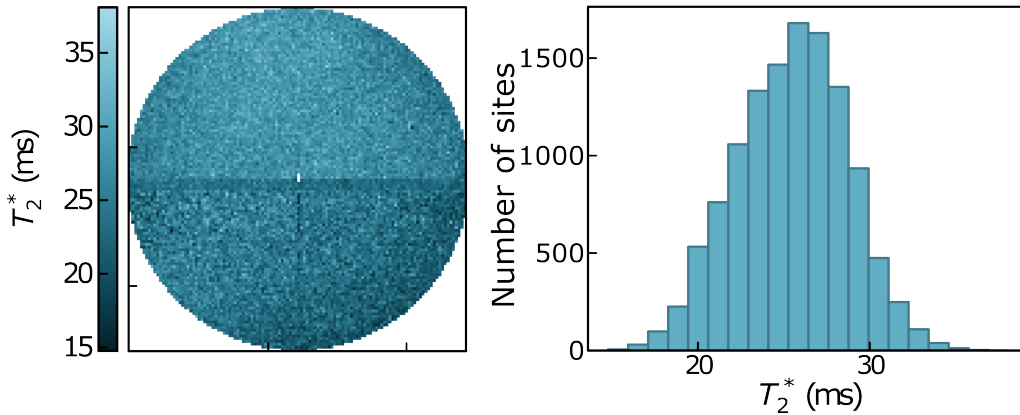


Figure 3.7: **Uniformity of T_2^* coherence time across the array** Map and histogram of T_2^* across the atom array. The average uncertainty per site is 1.5 ms. The average T_2^* for sites generated by the fiber amplifier at 1055 nm is 23.2(1) ms, while it is 25.58(3) ms for sites generated by the fiber amplifier at 1061 nm.

3.2.3 Dynamical Decoupling

In order to extend the operation time of a realistic quantum processor well beyond the dephasing time of the array, we can apply dynamical decoupling to the atomic qubits.

³We calculate the differential light shift between the clock states accounting for the shifts from the D_1 and D_2 lines whose transition frequencies relative to the qubit states differ by ~ 9.2 GHz.

Dynamical decoupling builds on the original spin-echo [127] sequence developed in the context of NMR, which applies a π pulse halfway through an evolution time, such that dephasing caused by static, inhomogeneous fields are canceled. Applying this sequence more frequently to address slowly varying noise was implemented in the Carr-Purcell-Meiboom-Gill (CPMG) pulse train [128].

Spin-echo and CPMG type sequences protect dephasing from a qubit that is pointing along a particular axis, however to protect qubits from dephasing, no matter their current state, the XY4 ($X - Y - X - Y$) sequence was introduced, and soon after, the XY8 ($X - Y - X - Y - Y - X - Y - X$) and XY16 ($XY8 - \bar{X}\bar{Y}8$) sequences, which provide a mirrored structure of pulses that can provide protection against errors in control operations [129]. We utilize in our experiments such pulse sequences with mirror symmetry, such as XY4, XY8, and XY16 sequences [130].

It has been proven that theoretically with a fast enough sequence of unitary pulses, all dephasing could be reversed [131], limited by the error applied due to the pulses themselves.

We test XY4, XY8, and XY16 sequences on our array and find empirically that an XY16 dynamical decoupling sequence with a period of 12.5 ms yields the longest dephasing time for the reduced trap depth of $k_B \times 55 \mu\text{K}$ of the pulse sequences that we test.

3.2.4 T_2 time

By applying cycles of XY16 sequences with a period of 12.5 ms between π pulses, we measure a dephasing time of $T_2 = 12.6(1)$ s, a new benchmark for the coherence time of an array of hyperfine qubits in optical tweezers [51, 52] (Fig. 3.8). In addition, we investigate in Fig. 3.9 the coherence time at different trap depths, yielding notably $T_2 = 3.19(5)$ s at the full trap depth of $k_B \times 0.18$ mK. Although lower, this result also surpasses previous known results with hyperfine qubits in a tweezer array.

To measure the coherence time with dynamical decoupling, we prepare the qubits in $|1\rangle$, and then apply a $\frac{\pi}{2}$ pulse, wait half the period time, and subsequently apply XY16, where each X and Y π pulse is separated from the previous one by the period. After the final pulse, we wait half of the period time and then apply another $\frac{\pi}{2}$ with phase 0 or π . We vary the number of symmetric XY16 cycles and we obtain the coherence contrast by subtracting the population difference of the two readout cases (applying 0 or π phase). We measure a coherence contrast which we fit to an

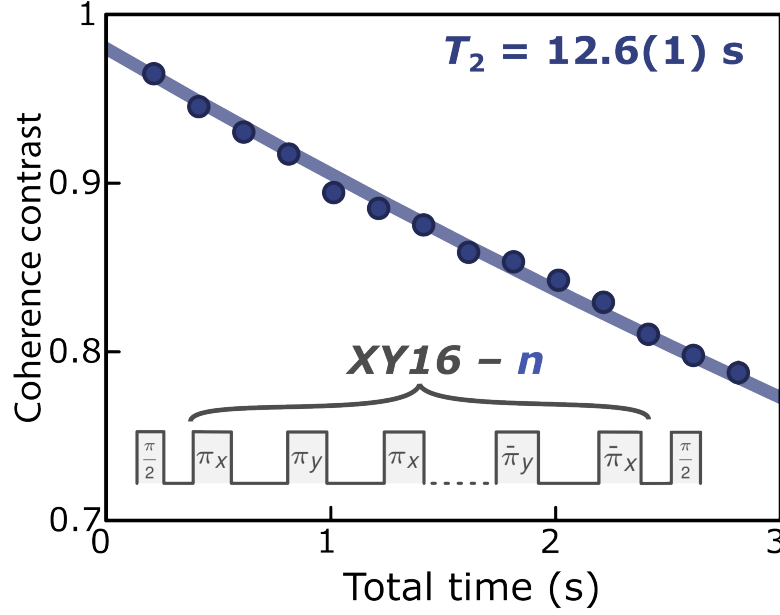


Figure 3.8: **T_2 coherence time.** Measurement of the dephasing time T_2 after dynamical decoupling. After an initial $\pi/2$ pulse, a variable number of XY16 dynamical decoupling cycles with a fixed time $\tau = 12.5 \text{ ms}$ between π pulses are used to offset the reversible dephasing. The phase of the final $\pi/2$ pulse is chosen to be either 0 or π , and subtracting the population difference in these two cases provides the coherence contrast. The contrast decay is fitted to obtain $T_2 = 12.6(1) \text{ s}$.

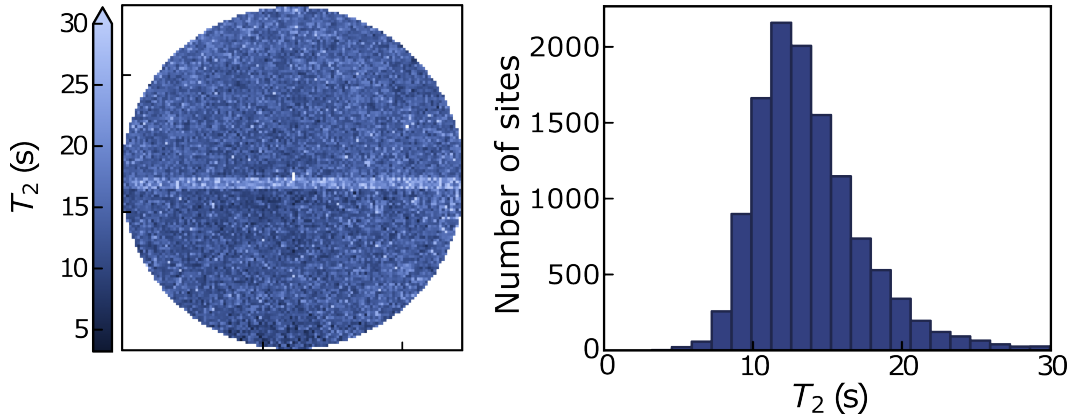


Figure 3.9: **Uniformity of T_2 coherence time across the array.** Map and histogram of T_2 across the atom array. The average uncertainty per site is 2.8 s. The average T_2 for sites generated by the fiber amplifier at 1055 nm is 19.2(4) s, while it is 12.32(6) s for sites generated by the fiber amplifier at 1061 nm. We use averages weighted by the uncertainty on each site, since we observe that the unweighted average results in a bias from the value obtained by global fitting.

exponential decay to find $T_2 = 12.6(1)$ s. This is, to our knowledge, a record for coherence time in alkali atom arrays.

We investigate in Fig. 3.9 the coherence time T_2 as a function of the trap depth for two different periods between π pulses (only for atoms trapped with the fiber amplifier at 1061 nm), 12.5 ms and 6.2 ms. We attribute the different optimal periods at different depths to a trade-off between the unfiltered noise at a specific dynamical decoupling period [131] and the effective depolarization induced by each π pulse. At the full trap depth, we measure a coherence time of 3.19(5) s, which still constitutes a record for hyperfine qubits in a tweezer array.

Considering the Raman scattering rate at a trap depth of $k_B \times 0.18$ mK, we expect that a significantly longer coherence time should be achievable. Based on this observation and the discrepancy in coherence time between atoms trapped at 1061 nm and 1055 nm seen in site-resolved data (Fig. 3.9), we posit that the observed coherence time is limited by intensity noise due to the trapping lasers or the SLMs.

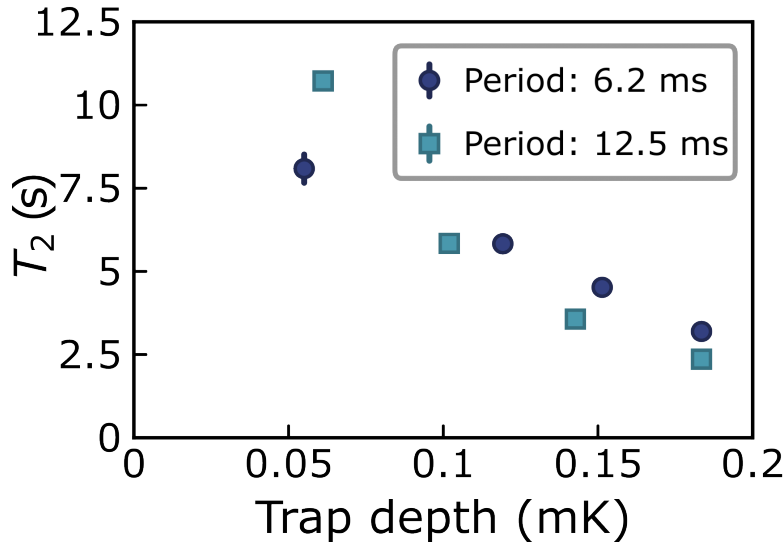


Figure 3.10: **T_2 dependence on trap depth.** Measurement of the coherence time T_2 at different trap depths, for two different periods between π pulses. Error bars indicating 68% confidence intervals are shown when larger than the dot itself. Note that this experiment was performed with the 11,513 tweezers generated by the fiber amplifier at 1061 nm only (Fig. 2.12).

Universal Gates Overview

We provide a reference for the definition of various gates used to construct universal gate sets, as well as some relevant properties.

Single-qubit gates: Single-qubit gates can be represented as rotations on the *Bloch sphere*, whose surface can be defined as the set of states $\{|\Psi\rangle = \cos(\frac{\theta}{2})|0\rangle + e^{i\phi}\sin(\frac{\theta}{2})|1\rangle\}$ accessible by varying (θ, ϕ) . The Hadamard gate H is a single-qubit π rotation around the $\frac{1}{\sqrt{2}}(x + z)$ axis on the Bloch sphere, up to a global phase. In matrix form, H is given as $H = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 & 1 \\ 1 & -1 \end{pmatrix}$, which is equivalent to $H = \frac{1}{\sqrt{2}}(X + Z)$. We note also that $HZH = X$, $HXH = Z$, and $H^2 = I$. We write X, Y, Z to represent Pauli matrices. $R_x(\theta)$ and $R_z(\theta)$ represent arbitrary single-qubit rotations by angle θ around the x -axis and z -axis respectively, and can be written as $R_k(\theta) = e^{-i\frac{\theta}{2}k}$ for $k \in \{X, Y, Z\}$. In matrix form,

$$R_x(\theta) = \begin{pmatrix} \cos(\frac{\theta}{2}) & -i\sin(\frac{\theta}{2}) \\ -i\sin(\frac{\theta}{2}) & \cos(\frac{\theta}{2}) \end{pmatrix} \quad (3.5)$$

$$R_y(\theta) = \begin{pmatrix} \cos(\frac{\theta}{2}) & -\sin(\frac{\theta}{2}) \\ \sin(\frac{\theta}{2}) & \cos(\frac{\theta}{2}) \end{pmatrix} \quad (3.6)$$

$$R_z(\theta) = e^{-i\frac{\theta}{2}} \begin{pmatrix} 1 & 0 \\ 0 & e^{i\theta} \end{pmatrix} \quad (3.7)$$

and in relation to the Hadamard gate, $HR_z(\theta)H = R_x(\theta)$. We note also that $R_z(\theta) = R_y(-\frac{\pi}{2})R_x(\theta)R_y(\frac{\pi}{2})$. The T gate is then a $\frac{\pi}{4}$ rotation around the z -axis, so $T = e^{i\frac{\pi}{8}}R_z(\frac{\pi}{4})$.

Two-qubit gates: We can write the CNOT entangling gate in the following way between qubits a and b :

$$CNOT = |0\rangle_a\langle 0|_a \otimes I_b + |1\rangle_a\langle 1|_a \otimes X \quad (3.8)$$

and a CP (controlled-phase) gate is given by

$$CP = |0\rangle_a\langle 0|_a \otimes I_b + |1\rangle_a\langle 1|_a \otimes e^{i\frac{\theta}{2}}R_z(\theta). \quad (3.9)$$

Here, the CP gate is a CZ gate when $\theta = \pi$. A CZ gate can be transformed to a CNOT gate in the following way:

$$CNOT = (I_a \otimes H_b)CZ(I_a \otimes H_b).$$

3.3 Single Qubit Gates

The coherence times that we measure represent something of a fundamental characterization of the uniformity that qubits in our array experience, and the length of time that information can be held and operated on. We now turn to implementing single-qubit gates and characterizing the fidelity of these operations utilizing randomized benchmarking. We characterize a high single-qubit gate fidelity of over 99.98%, however, we also note that we believe that this does not represent a fundamental limit of our achievable fidelities, since we are implementing these gates using a microwave horn, and could instead pivot to implementing single-qubit gates via a two-photon Raman process which can have two orders of magnitude higher Rabi frequencies. We have demonstrated the use of such transitions for measuring our trap frequencies as was presented in section 2.5.2 and Fig. 2.27

3.3.1 Universal Gate Set

In order to utilize a platform for quantum computing, it must be able to demonstrate a *universal gate set*. According to Nielsen and Chuang [132], ‘a set of gates is said to be universal for quantum computation if any unitary operation may be approximated to arbitrary accuracy by a quantum circuit involving only those gates’.

A few primary examples of such universal gate sets include $\{CNOT, R_x(\theta), R_z(\theta)\}$ and $\{H, T, CNOT\}$. Generally, any two-qubit entangling gate along with arbitrary single-qubit rotations forms a universal gate set [132]. For neutral atom quantum computing, the native gate entangling gate is the controlled-Z (CZ) or the controlled-phase (CP) gate. The CZ gate can be transformed into the $CNOT$ gate via conjugation with Hadamard gates, up to a global phase.

3.3.2 X,Y,Z gates

To implement rotations around $\{X, Y, Z\}$, we utilize a microwave horn to drive the qubits. Implementing X and Y gates involves applying an RF modulation of the form $\Omega \cos(\omega t + \phi)$, where Ω is the Rabi frequency or strength of the RF drive, ω is the drive frequency, and ϕ is the phase. We implement gates by driving the qubit at the qubit frequency. When $\phi = 0$, the qubit rotates around the X-axis, and when $\phi = \frac{\pi}{2}$, we implement a qubit rotation around the Y-axis. In particular, a pulse rotating a qubit by an angle θ with phase ϕ applies the unit operator

$$U(\theta, \phi) = e^{-i\frac{\theta}{2}(X \cos \phi + Y \sin \phi)}. \quad (3.10)$$

In order to implement a rotation around the Z-axis, one could apply combinations of rotations around the X and Y axes, since

$$R_z(\theta) = R_y(-\frac{\pi}{2})R_x(\theta)R_y(\frac{\pi}{2}). \quad (3.11)$$

This amounts to the qubit rotating relative to the X and Y axes. However, one can also instead rotate the reference frame of the axes relative to the qubit instead in a so-called virtual-Z (VZ) gate [133]. Doing so has the significant benefits of not causing control errors, and also not requiring time for implementation, since it is applied simply by shifting the phase ϕ of subsequent operations. We reproduce a helpful example found in [133] to demonstrate that this is an effective approach. In particular, if one implements two $R_x(\theta)$ gates where the second gate is delayed by the phase ϕ , the operator that is applied will be

$$e^{-i\frac{\theta}{2}(X \cos \phi + Y \sin \phi)} R_x(\theta) = e^{i\frac{\phi}{2}Z} e^{-i\frac{\theta}{2}X} e^{-i\frac{\phi}{2}Z} R_x(\theta) \quad (3.12)$$

$$= R_z(-\phi)R_x(\theta)R_z(\phi)R_x(\theta). \quad (3.13)$$

If the phase offset is tracked such that subsequent gates are implemented in the new frame, the final z-rotation will be subsumed into the following gate. And since the measurement is done in the z-basis, this will not affect the final measurement outcome.

3.3.3 Randomized Benchmarking

Characterizing gate fidelities can be done in several ways, but the approach of randomized benchmarking (RB) [134] has become a standard approach. This is because RB has two significant benefits. Firstly, RB allows the fidelity to be determined in a fashion that is independent from state preparation and measurement errors. Secondly, RB is a *scalable* approach because it uses gates from the Clifford group, which is provably scalable, since Clifford gates can be simulated efficiently classically. The Clifford group for two qubit gates C_2 is generated by $\{H, S, CNOT\}$, where S is the so-called phase gate, $S = R_z(\pi/2)$. While it does not include the full universal gate set, stabilizer measurements for quantum error correcting codes are for instance made up of Pauli operators which are included in the Clifford group. The single-qubit Clifford group C_1 is generated by $\{H, S\}$ and contains only 24 operations, in comparison with the 11,520 operations in the C_2 group.

To implement single-qubit RB, we first prepare the qubits in the $|1\rangle$ state, and then implement some number of gates n , after which we apply the inverse of all the gates, and measure the return probability to the $|1\rangle$ state. For each given length n , we

select $U_{n-1}, \dots, U_0$ at random from the 24 unitaries composing the Clifford group. We then apply $U_{-1}U_{n-1} \cdots U_0$ where U_{-1} is the inverse of $U_{n-1} \cdots U_0$, which we calculate on a classical computer. We decompose Clifford gates into elementary rotations around Bloch sphere axes using the zyz Euler angles. Rotations around z are implemented virtually, as described in the previous section, by offsetting the phase of all following x and y rotations [133].

3.3.4 Robust Gates

As was previously presented in Fig. 3.4, we have an inhomogeneous Rabi frequency across the array which causes *pulse length error* in our gates— some qubits experiencing over-rotation and others experiencing under-rotation for a given gate that has a pulse length determined by the average Rabi frequency of the array.

In order to address such errors, that are particular to the noise sources in our system, we turn to error-resilient composite pulses or *robust gates*. There are many families of such gates that allow for platforms with particularly limiting error sources to design composite pulse sequences that together both implement the desired unitary gates, while remaining robust to a particular error source. Common error sources that can be addressed with composite pulses include pulse length errors (such as caused by inhomogeneous Rabi frequency) and off-resonant errors (such as caused by variable light shift across different tweezers). To address our pulse-length errors, a few relevant pulse sequences are the BB1 sequence [135], the SK1 sequence [136], and the SCROFULOUS sequence [137].

Among such common families of error-resilient pulses [135, 137, 138], we find that SCROFULOUS performs the best in our case. The SCROFULOUS in particular is short compared to other pulse-length-robust sequences, requiring only three pulses. Given that we are using microwave transitions rather than optical transitions, our Rabi frequency is already significantly longer, and thus utilizing a shorter robust gate proves optimal.

The SCROFULOUS pulse sequence implements a target rotation $U(\theta, \phi)$ using a symmetric composite pulse $U(\theta_3, \phi_3)U(\theta_2, \phi_2)U(\theta_1, \phi_1)$ where

$$\theta_1 = \theta_3 = \arcsinc(2 \cos(\theta/2)/\pi) \quad (3.14)$$

$$\phi_1 = \phi_3 = \phi + \arccos\left(-\frac{\pi \cos \theta_1}{2\theta_1 \sin(\theta/2)}\right) \quad (3.15)$$

$$\theta_2 = \pi \quad (3.16)$$

$$\phi_2 = \phi_1 - \arccos\left(-\frac{\pi}{2\theta_1}\right) \quad (3.17)$$

In our implementation, the average pulse area for a random Clifford unitary is 2.02π .

3.3.5 Single Qubit Gate fidelity

We thus perform global single-qubit randomized benchmarking, using SCROFULOUS composite pulse [138] to compensate for our inhomogeneous Rabi frequency. We apply n gates, up to $n = 1000$ sampled from the Clifford group C_1 , followed by an inverse operation, and measure the final population in $|1\rangle$ (Fig. 3.11). We average over 60 different strings of n random gates for each number of gates. Fitting the decay as the number of gates increases yields an average Clifford gate fidelity $F_c = 0.999834(2)$, limited by phase noise in our system likely due to magnetic field noise. This could be readily addressed by upgrading the current sources driving the magnetic field coils or by operating at MHz-scale Rabi rates with optical Raman transitions [72]. We additionally plot the site-resolved single-qubit gate fidelity in Fig. 3.13, demonstrating uniform, high-fidelity gates across atoms in all 12,000 tweezer sites.

We extract the average gate fidelity by modeling the sequence as a depolarizing channel. As the number of applied Clifford gates n increases, the sequence depolarizes the initial state toward the fully mixed state, which for a single qubit corresponds to a final population of $\frac{1}{2}$. We fit the decay of the final population with the number of applied Clifford gates as $\frac{1}{2} + \frac{1}{2}(1 - d_0)(1 - d)^n$ where the offset d_0 stems from SPAM errors, d is the average depolarization probability at each gate and n is the number of gates. The average Clifford gate fidelity is then given by [139]:

$$F_c = 1 - d/2 \quad (3.18)$$

3.3.6 Limitations on gate fidelity

Even though the measured single-qubit gate fidelity is competitive with other state-of-the-art atom array experiments [35, 52, 53, 140], single-qubit gate fidelities >0.9999 have been reported [141, 142] in smaller arrays. Moreover, we can estimate

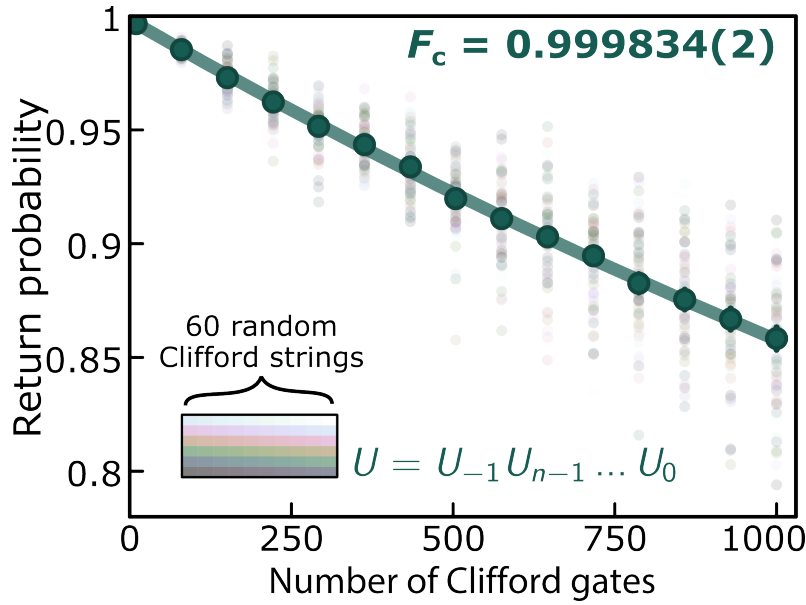


Figure 3.11: **Randomized benchmarking of single-qubit gates.** Randomized benchmarking of the global single-qubit gate fidelity. For each number of Clifford gates, 60 different random gate strings of this length are applied, after which the overall inverse of the string is applied. For a given gate string length, each translucent marker of a given color represents the return probability for a string of gates while the solid green markers indicate the averaged return probability over the 60 different strings. The inset lists all the colors used to indicate the 60 random gate strings for a given length. The decay of the final population to $1/2$ is fitted to $(1 - d)^N$ and $F_c = 1 - d/2$ represents the average Clifford gate fidelity.

the maximal theoretical fidelity achievable for a given dephasing time T_2^* . Assuming the dephasing is only limited by the thermal distribution of atoms yields a theoretical maximum fidelity of [123, 143]:

$$\mathcal{F} = \frac{3}{4} + \frac{1}{4(1 + 0.95(t/T_2^*)^2)^{3/2}} \quad (3.19)$$

where t is the average time needed to apply a Clifford gate, $t = \langle \theta \rangle / \Omega$; $\langle \theta \rangle$ being the average pulse area per Clifford gate. Hence, gate fidelities higher than 0.99999 should be achievable solely based on this value.

Beyond infidelities due to decoherence, other parameters that may limit single-qubit gate fidelities are: (a) amplitude errors due to instabilities in the microwave power; (b) phase errors due to the microwave setup; (c) phase errors due to optical tweezer intensity noise; (d) phase errors due to magnetic field noise. We are interested in which of these factors is limiting the gate fidelity. We rule out (a) because we

observe that the Rabi frequency is very stable shot-to-shot (variations of less than 0.1 %), and we estimate that such variations should be completely suppressed by the SCROFULOUS pulse. We also rule out (c) since reducing the trap depth further does not significantly improve the randomized benchmarking results (Fig. 3.12), and the fidelity is identical for atoms trapped in tweezers at 1055 nm and 1061 nm as seen in Fig. 3.13 (unlike T_2^* and T_2). Although we cannot formally rule out (b), we estimate that it is unlikely since active components in the microwave setup have a very low phase noise, and we observe a sub-10 Hz linewidth of the microwave signal with a spectrum analyzer.

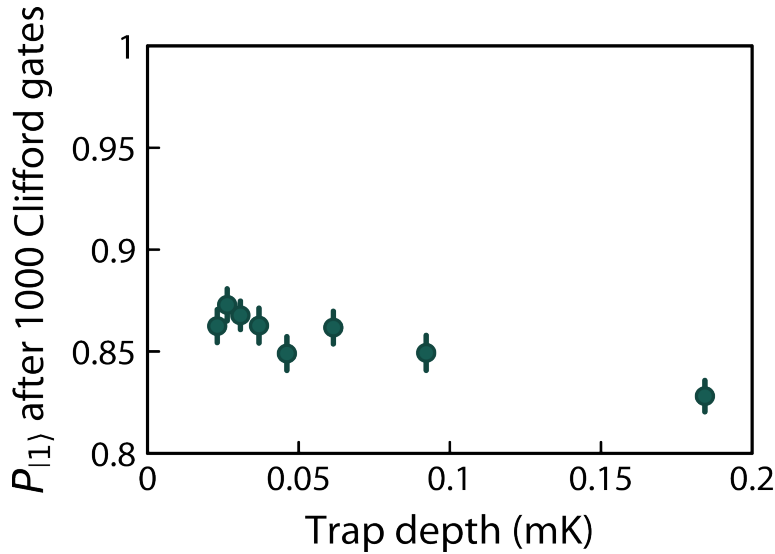


Figure 3.12: **Single-qubit RB dependence on trap depth** The population in $|1\rangle$ after 1,000 random Clifford gates is measured for different trap depths, exhibiting only limited improvement when the trap depth is reduced. This indicates that the single-qubit gate fidelity is likely limited by residual magnetic field noise, which could be readily mitigated by technical improvements. Error bars indicate estimated 68% confidence intervals.

We do, however, notice a dominant phase noise at 60 Hz in the qubit array due to the mains AC voltage. We measure the intensity of this noise with a spin-echo sequence, where the time between each pulse is $\tau = 1/(2 \times 60 \text{ Hz})$ (Fig. 3.14). Although this low-frequency noise cannot by itself explain the single-qubit gate fidelity loss, it points generally to residual magnetic field noise that could be mitigated by shielding the vacuum cell, upgrading the current sources driving the magnetic field coils, and/or by operating at MHz-scale via Raman transitions. This can be achieved, for instance, by utilizing the amplitude-modulation setup used for Raman sideband spectroscopy.

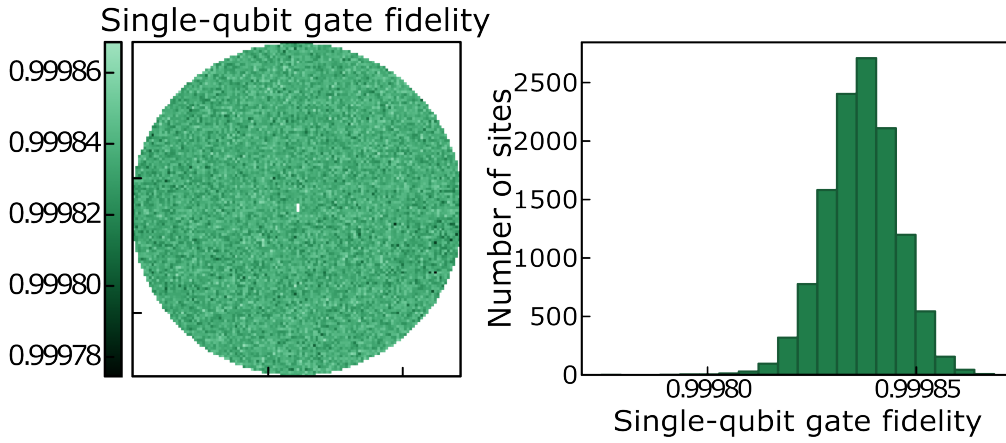


Figure 3.13: **Uniformity of single-qubit gate fidelity across the array** Map and histogram of single-qubit gate fidelity obtained by global randomized benchmarking. The average gate fidelity is 99.9834(2)%.

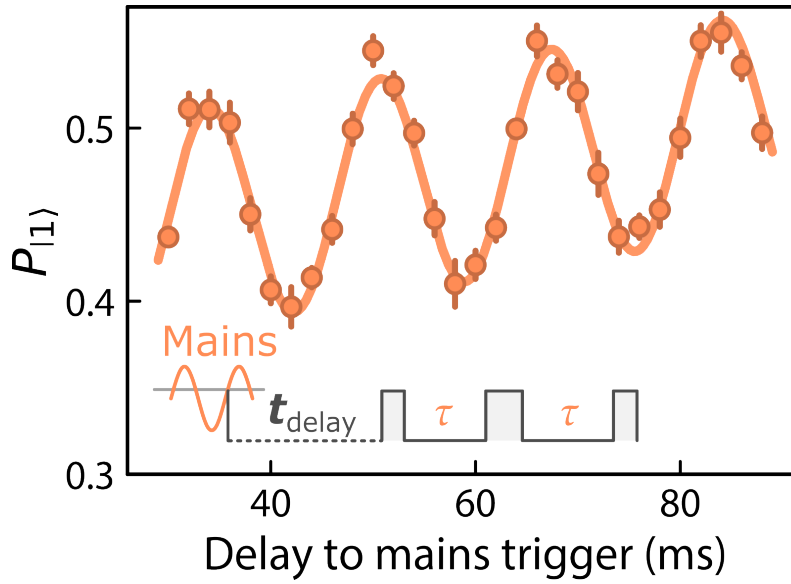


Figure 3.14: **60 Hz Noise** A spin-echo sequence is employed to probe the 60 Hz phase noise in our system. The free-evolution time of each arm, τ , is set to a half-period of 60 Hz, which enhances the noise. By varying the time t_{delay} between the line trigger and the spin-echo sequence, we map the periodic noise at 60 Hz to the population in $|1\rangle$.

Chapter 4

NEUTRAL ATOM ARCHITECTURE: CONNECTIVITY

One central advantage of using neutral atoms for quantum computing lies in the arbitrary connectivity that can be engineered between qubits via atom transport. This enables, for instance, quantum error correcting codes with significantly higher encoding rates than can be achieved with simply nearest neighbor connectivity [12, 13].

Quantum error correction (QEC), in which information is protected from decoherence by being encoded in logical qubits across many physical qubits, is necessary to implement in order to run algorithms or simulations that are useful for real-world problems with high enough fidelity. Encoding the information at all already requires many qubits, and to reach high enough logical qubit fidelities, even more qubits. Thus, utilizing codes with higher encoding rates of the number of atoms or qubits that are necessary for encoding atoms will allow for quantum computers to be more useful sooner, as has been put forward in recent resource estimates [14, 144].

Additionally, implementing universal quantum computation requires single-qubit and two-qubit gates which have been implemented either through single-site addressing [52] or a zone-based architecture [35]. Zone-based architectures leverage the ability to dynamically move atoms in a coherence-preserving manner [51, 145], enabling long-range, non-local connectivity, which allows for less stringent quantum error correction bounds [146]. This architecture also provides a path for mid-circuit readout [35].

Thus, for the purpose of universal quantum computation, as well as looking towards future QEC applications requiring high connectivity, we examine the atom transport at scale in the context also of enabling zone-based computation. While moving atoms using acousto-optic deflectors (AODs) is now a well-established practice to resort them into a deterministic configuration [30, 31] or to transport them coherently [35, 51, 54, 147], the distances across which it will be necessary to transport atoms to enable computation in our 6,100 qubit array are significantly further than previously demonstrated distances for single-atom transport with tweezers [35, 51].

After examining long distance atom transport (4.1) and long distance coherent atom transport (4.2), we then present a proposal including details for utilizing these

atom transport techniques in a zone architecture in our system (4.3). We go on to briefly discuss aspects of qubit hiding (4.4) and two-qubit gates (4.5), presenting some preliminary data, and discussing application to realistic use of the presented zone-architecture.

4.1 Long distance transport

We begin by isolating effects purely due to long-distance transport, loading directly from the MOT into 10 transport tweezers created using 1055 nm light through two AODs mounted in a crossed configuration. We utilize the most wide-band commercially available AODs at near-infrared wavelengths (Gooch & Housego AODF 4085), which cover up to 500 – 600 μm along one axis for the optical parameters used here. These AODs have more resolvable spots due to having a larger crystal active-aperture of over 15 mm. We map the output after the pair of AODs to the back aperture of the objective using a telescope with 3:2 demagnification to match the same beam size at the back aperture of the objective as the beam from the SLM trapping tweezers.

We measure the survival of atoms as a function of the transport time using a constant jerk trajectory¹, and moving by sweeping the frequency along just one axis. Moving longer distances requires a significantly longer time in order to mitigate atom loss, as seen in Fig. 4.1. Additionally, we observe that deeper traps allow for faster movement before there is significant atom loss.

4.1.1 AOD tweezer movement

In probing atom survival as a function of long-distance movement speeds, we find that the speed of transport is strongly constrained by cylindrical lensing—an effect that deforms the trapping potential, and occurs when the AOD frequency is rapidly swept [148]—which becomes increasingly deleterious as the AOD field of view is increased.

This becomes clear because when we implement a diagonal trajectory (by implementing the same sweep on a vertical and horizontal AOD simultaneously), we can move significantly faster during a single move without seeing atom loss than we can for a straight move. Notably, using a pair of crossed AODs for diagonal transport converts cylindrical lensing into spherical lensing, enabling significantly faster movement (Fig. 4.2).

¹We note that in transport data where we implement many moves, we use an adiabatic sine trajectory, which we find incurs less heating.

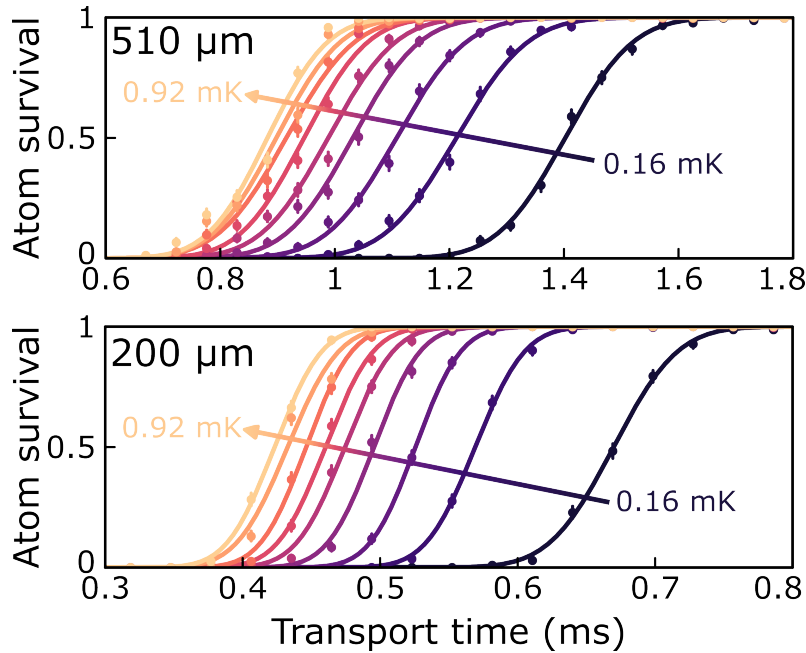


Figure 4.1: **Survival during long-distance atom transport** Experimentally measured atom survival for different trap depths and two distances of 510 μm and 200 μm . The signal for each depth is fitted to an error function. Atoms are transported using the constant jerk trajectory with a straight, single-axis move as opposed to the diagonal moves used in the coherent transport study in the subsequent sections. Note that these moves could be executed considerably faster by mitigating cylindrical lensing, as elaborated on in section 4.1.1.

Cylindrical Lensing

When moving the atoms longer distances, cylindrical lensing becomes an increasingly pronounced problem. Since moving larger distances requires utilizing more resolvable spots, the AOD active aperture and incident beam are both larger. Therefore, as a radio-frequency is swept, different parts of the beam will sample a larger range of frequencies for the same sweep than a smaller beam would, thus causing astigmatism or *cylindrical lensing* at the focal plane. This effect then determines a critical speed for moving tweezers that is impossible to exceed.

As the tweezer is moved using a changing radio-frequency wave in the AOD crystal, the focus of the tweezer is offset from the atom plane by the distance characterizing the astigmatic focal shift, defined as

$$z_s = \frac{f}{V} \dot{x}_0(t) \quad (4.1)$$

where f is the effective focal length of the objective (including any magnification

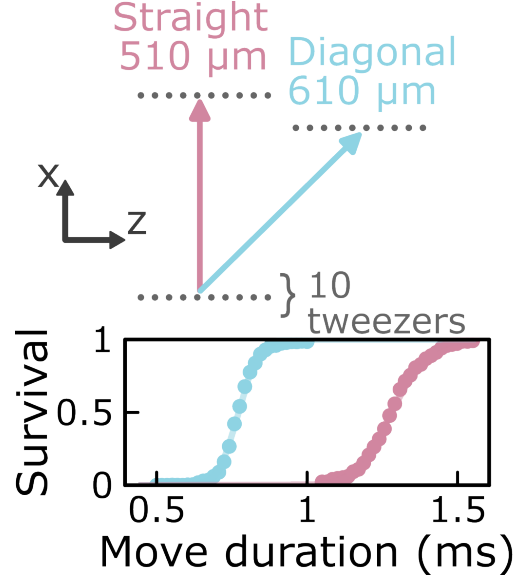


Figure 4.2: **Survival with and without cylindrical lensing.** Schematic and atom survival for a diagonal (blue) or straight (pink) move for ten tweezers (with depth $k_B \times 0.28$ mK) spaced by ~ 10.6 μm . Despite being shorter, a straight move needs to be executed more slowly than a diagonal one due to cylindrical lensing.

after the AOD), V is the velocity of the acoustic wave in the AOD, and $\dot{x}_0(t)$ is the tweezer velocity. This results in a modified tweezer potential $U \propto -\tilde{U}_x \tilde{U}_y$ that can be described, when moving in the x direction as

$$\tilde{U}_x = \frac{1}{\left(1 + \left(\frac{z-z_s}{z_R}\right)^2\right)^{1/2}} \exp\left(-\frac{2x^2}{w_0^2 \left(1 + \left(\frac{z-z_s}{z_R}\right)^2\right)}\right), \quad (4.2)$$

$$\tilde{U}_y = \frac{1}{\left(1 + \left(\frac{z}{z_R}\right)^2\right)^{1/2}} \exp\left(-\frac{2y^2}{w_0^2 \left(1 + \left(\frac{z}{z_R}\right)^2\right)}\right). \quad (4.3)$$

The ratio of $|z_s|$ to z_R (the tweezer Rayleigh range) determines two regimes: when $|z_s| < 2z_R$, the tweezer is still well defined (i.e., there is a single local potential minimum, albeit broadened by cylindrical lensing) and the trap depth is reduced by $1 + \frac{1}{4} \left(\frac{z_s}{z_R}\right)^2$. When the tweezer is moved too quickly in one direction and $|z_s|$ exceeds $2z_R$, the trap splits such that there are two potential minima, causing the most deleterious effects. This motivates us to introduce a *characteristic shift velocity*, v_s , such that

$$\frac{z_s}{2z_R} = \frac{\dot{x}_0}{v_s}. \quad (4.4)$$

Notably, v_s constitutes a velocity that is difficult to exceed during atom transport without atomic loss, and is given by: $v_s = 2\pi V w_0^2 / f \lambda$. By assuming Airy diffraction through the objective, we obtain a simple formula:

$$v_s \approx 5.3 \frac{w_0}{T_a}, \quad (4.5)$$

where T_a is the AOD access time, i.e. the propagation time of the acoustic wave through the input beam [149]. (Note that the exact factor depends on the optical configuration; in our case, we obtain results consistent with a factor 7 – 8 instead of 5.3). This formula highlights an important trade-off: T_a factors in the time-bandwidth product ($T_a \Delta f$), a figure of merit for AODs that is roughly equal to the number of resolvable spots. Thus, as one tries to scale up the addressable field of view of AODs, v_s decreases.

Using a diagonal trajectory (by which we mean that tweezers move with the same rate along x and y ; $\dot{x}_0 = \dot{y}_0$, or equivalently the chirp rates in the AODs are equal: $\dot{f}_x = \dot{f}_y$) allows this splitting to be eliminated, minimizing the impact of lensing, though there remains a penalty due to the off-axis focus shift. This can be seen by recognizing that the overall potential in this case can be written as

$$U \propto \frac{1}{1 + (\frac{z-z_s}{z_R})^2} \exp\left(-\frac{2(x^2 + y^2)}{w_0^2 \left(1 + (\frac{z-z_s}{z_R})^2\right)}\right), \quad (4.6)$$

where there is only one potential minima (at position z_s), that has been spherically lensed due to the symmetry in x and y (the minimum tweezer waist is larger).

Thus, to mitigate the effects of cylindrical lensing in the data that we present in the following sections, we have used diagonal moves. While there is still a penalty associated with the focal shift of the tweezer, it is much reduced compared with the straight move situation. We have leveraged this strategy to realize faster high-fidelity moves, as presented in Figs. 4.5, 4.6, and 4.8. In section 4.3 we describe how a zone-based neutral atom quantum processor can be engineered with these diagonal moves.

How is heating incurred from atom trajectories?

For a stationary atom in a trap, under the harmonic oscillator approximation, the Hamiltonian is given as $H_0 = \frac{p^2}{2m} + \frac{1}{2}m\omega^2 x^2$.

When the trap is accelerated, the atom then feels the force $F(t) = -ma(t)$, so that the perturbative potential is given as $V(t) = F(t)x$. Then from first order time-dependent perturbation theory, if the atom starts in the motional state $|n\rangle$, the probability amplitude (p_{nk}) of finding the atom in state $|k\rangle$ after time T , is given by

$$p_{nk} = \frac{1}{i\hbar} \int_0^T e^{i\omega_{nk}t} \langle k|V(t)|n\rangle dt \quad (4.7)$$

where here, $\omega_{nk} = (E_k - E_n)/\hbar$ is the transition frequency. By inserting our expression for $V(t)$ and $x = \sqrt{\frac{\hbar}{2m\omega}}(a^\dagger + a)$, we find

$$p_{nk} = -i\sqrt{\frac{m}{2\omega\hbar}}(\delta_{k,n-1}\sqrt{n} + \delta_{k,n+1}\sqrt{n+1}) \int_0^T a(t)e^{i\omega_{kn}t} dt \quad (4.8)$$

So then, $\Delta N = |p_{n,n+1}|^2 - |p_{n,n-1}|^2$ which we can calculate to be

$$\Delta N = \frac{m}{2\hbar\omega} \left| \int_0^T a(t)e^{i\omega t} dt \right|^2. \quad (4.9)$$

By taking $s = 2t/T - 1$, we can write the adiabatic sine trajectory given in equation 4.12, and assuming the atom starts at rest and ends at rest, as

$$x(t) = D \left[\frac{t}{T} - \frac{1}{2\pi} \sin\left(\frac{2\pi t}{T}\right) \right], \quad (4.10)$$

which then gives $a(t) = \frac{2\pi D}{T^2} \sin(2\pi t/T)$, where D is the distance of the trajectory, T is the time of the trajectory, and ω is the trap frequency. Taking the Fourier transform, then gives that the increase in the average radial motional quanta ΔN for an adiabatic sine trajectory scales as $\Delta N \propto \frac{D^2}{\omega^5 T^6}$.

In the case of a constant jerk trajectory,

$$x(t) = D \left[3\left(\frac{t}{T}\right)^2 - 2\left(\frac{t}{T}\right)^3 \right], \quad (4.11)$$

we find $\Delta N \propto \frac{D^2}{\omega^3 T^4}$.

4.2 Long distance coherent transport

We not only require transporting atoms without loss, but also that the quantum information is maintained. We thus perform a Ramsey measurement before and after the ten AOD tweezers are transported by $610\text{ }\mu\text{m}$ in 1.6 ms.

We demonstrate in Fig. 4.3 negligible loss of coherence for atoms transported by $610\text{ }\mu\text{m}$ in 1.6 ms. We suppress dephasing with one XY4 dynamical decoupling sequence per move.

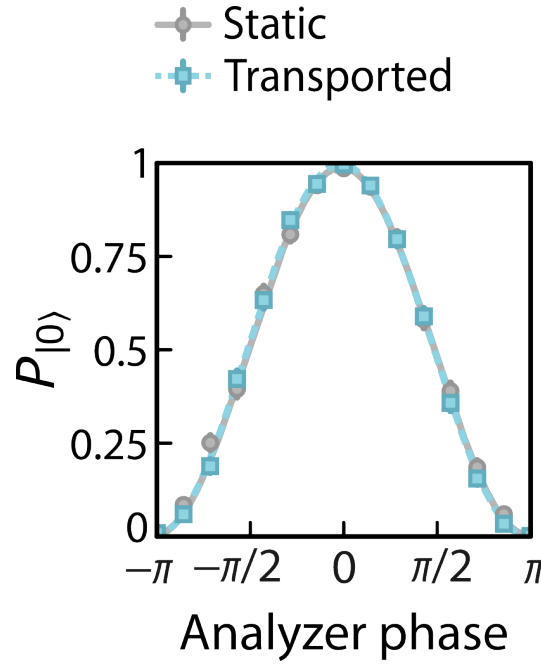


Figure 4.3: **Coherent transport.** Coherence of an atom after being transported diagonally $610\text{ }\mu\text{m}$ (blue) in 1.6 ms or held stationary (gray).

This appears promising, however, we note that realistic applications of coherent transport for quantum computation require many moves, not just single moves as demonstrated thus far. In particular, as an example, for movement of check-qubits in high-encoding rate codes, tens of moves could be needed for syndrome measurements, with increasing numbers of moves required depending on the distance of the code. Thus, we pivot to characterization of the survival and fidelity of many moves using interleaved randomized benchmarking.

4.2.1 Interleaved Randomized Benchmarking

In order to evaluate the fidelity of many long-distance moves, in the absence of cylindrical lensing, we therefore characterize both the survival of atoms after N

diagonal moves, as well as the fidelity of the *quantum channel* defined by coherent transport through interleaved randomized benchmarking [134] (IRB, Figs. 4.5, 4.6). To the best of our knowledge, such a quantitative characterization of transport fidelity in neutral atoms has not been previously demonstrated.

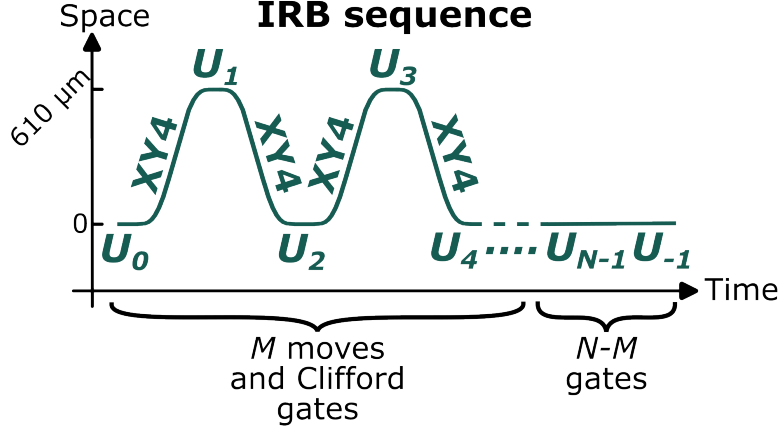


Figure 4.4: **Interleaved randomized benchmarking** IRB sequence used to benchmark the move fidelity. Random Clifford gates are interleaved between each of the $M(< N)$ moves, with the total number of gates N constant.

To perform interleaved randomized benchmarking, we fix a total number of single-qubit gates N drawn from the Clifford group C_1 . We then interleave $M(< N)$ total moves between the first M gates (atoms are held for $\sim 54 \mu\text{s}$ between moves), after which we apply the remaining $N - M$ gates to keep the total number of gates N constant and then apply the inverse of these gates. For the return probability data shown in Fig. 4.5 and Fig. 4.6, we average over 72 sequences of random gates for each number of moves and apply $N = 80$ total random single-qubit Clifford gates. For the static and transported return probabilities in the second panel, we apply the same single-qubit control sequence, including XY4 dynamical decoupling². To maximize the dephasing cancellation, we apply dynamical decoupling in a transformed Clifford frame (as described in a following subsection). As in the case of randomized benchmarking, we utilize SCROFULOUS pulses for implementing the Clifford gates.

For the atomic movement for the IRB, we use an adiabatic sine trajectory described by:

$$x = \frac{1}{\pi} \sin(\pi t) + t \quad (t, x \in [-1, 1]). \quad (4.12)$$

²We empirically find that the asymmetric XY4 sequence [150] performs best by evaluating the coherence contrast after 80 moves (implemented using bare pulses).

We find that we can execute a single move faster with the constant jerk trajectory [51] (which we use for Fig. 4.3 and Fig. 4.1), but that the adiabatic sine trajectory incurs less heating, as we see after implementing several moves.

Note that in the coherent transport data, the tweezer depth change along the trajectory is compensated with RF power which we calibrate beforehand with static tweezers at each position. We believe the transport fidelity can be further increased with more careful compensation of the trap depth including the AOD lensing effect in the future. We have recently implemented a more comprehensive approach for characterization and compensation, following [151].

Fitting return probabilities

At the end of the randomized benchmarking sequence, we measure both the atomic survival and the return probability (note that we apply a final π pulse to map the return state to the non-pushed-out state $|0\rangle$). We fit the atomic survival to a clipped Boltzmann distribution

$$S_n = 1 - \exp(-1/(a + bn)) \quad (4.13)$$

where a and b are respectively the normalized initial temperature and normalized temperature accumulated per move. For the selected durations for interleaved randomized benchmarking, we find that a is negligible. We then fit the return probability to

$$P = (1 - e^{-\frac{1}{bn}}) \cdot \left(\frac{1}{2} + \frac{1}{2}(1 - d'_0)(1 - d')^n \right), \quad (4.14)$$

where d' is the depolarizing probability for coherence, not accounting for atom loss. Owing to the randomized benchmarking procedure, coherence loss also includes the impact of XY4 dynamical decoupling since it would not be necessary without transport.

Since the resulting probabilities depend on the previous number of moves, as loss dominates at large numbers of moves, we define an *instantaneous fidelity*; i.e., the fidelity of the transport channel after a certain number of previous moves. We extract the instantaneous fidelity after n moves as:

$$F_n = \left(1 - \frac{d'}{2} \right) \frac{1 - e^{-\frac{1}{b(n+1)}}}{1 - e^{-\frac{1}{bn}}}. \quad (4.15)$$

Note that this is the most conservative approach and amounts to considering that the channel infidelity due to losses is equal to the loss probability itself. In the

context of fault-tolerant quantum computation, losses could be directly detected, leading to a higher tolerance to such errors than to Pauli errors. One could therefore assimilate loss to a depolarizing channel, which would significantly increase late-time instantaneous fidelities in Fig. 4.5, Fig. 4.6, Fig. 4.8, and Fig. 4.9. It is worth noting that losses are subdominant for early-time IRB results presented in Fig. 4.5 and Fig. 4.6, therefore the quoted early-time fidelity in these figures is independent of the specific model we use for losses. To compute the 68% confidence interval we bootstrap b and d' using the fit results and covariance matrix. We corroborate the obtained fidelity with a simple exponential fit for the first few data points of the return probability, where losses are negligible, and find similar early-time fidelities and error bars. We also notice that the shorter move of $270\text{ }\mu\text{m}$ has a slightly lower early-time instantaneous fidelity of $99.935(2)\%$ compared with the $610\text{ }\mu\text{m}$ -move ($99.953(2)\%$). We believe the discrepancy is likely due a trap depth calibration imperfection, and leave further investigation to future work.

Transformed Clifford Frame

During each move of the benchmarking sequence, we apply XY4 in a *transformed Clifford frame*. Previous works have examined the interplay of dynamical decoupling and quantum operations by; e.g., studying a system Hamiltonian in the “toggling frame” induced by dynamical decoupling pulses [152]. Here, we use related ideas but examine the decoupling operations in the frame rotated by the previously applied Clifford gates.

For instance, ignoring the Clifford gates between moves $k-1$ and k , it is possible to concatenate two XY4 sequences $X-Y-X-Y$ (with a symmetry operation) to obtain an XY8 sequence $X-Y-X-Y-Y-X-Y-X$ that yields higher-order dephasing (and pulse-length error) suppression. However, the random Clifford gate U_k between the two sequences will cancel this effect by twirling the second XY4 sequence with respect to the first one. Thus, we can “counter-twirl” the second XY4 sequence by applying it in a specific Clifford frame: the Pauli operator P becomes $P' = U_k^\dagger P U_k$. Up to a global phase, $U_k^\dagger X U_k$ and $U_k^\dagger Y U_k$ are two distinct elements of $\{X, Y, Z\}$, because U_k is a Clifford gate. If one of these two unitaries is Z , we further conjugate with a Hadamard gate H (or the equivalent basis change unitary between Y and Z) to map these two unitaries into X and Y , or Y and X . This can easily be generalized to n -qubit Clifford gates. A paradigmatic example is the transport between the storage and interaction zone to apply a CZ gate: since $CZ(X \otimes X)CZ = -Y \otimes Y$,

we can appropriately transform the decoupling sequence applied during the return move. This could also be extended to yield higher-order sequences, such as XY16. Importantly, typical architectures for fault-tolerant quantum computation (FTQC) use almost exclusively Clifford gates [153] (e.g., past the initial generation of noisy magic-state inputs, all gates are Clifford). Therefore this technique is fully applicable to FTQC.

4.2.2 Directly loaded AOD with long-distance moves

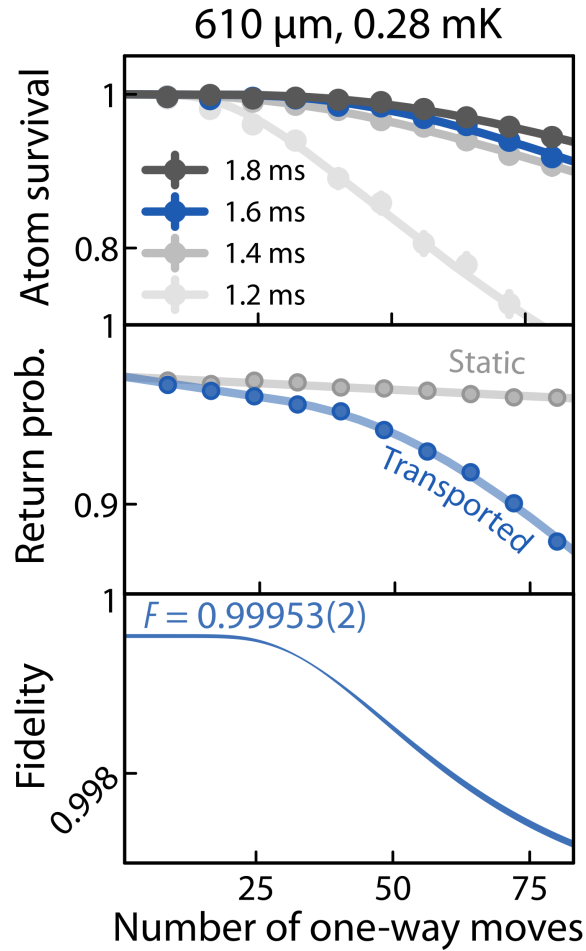


Figure 4.5: **610 μm AOD move survival and coherence.** Benchmarking results for repeated 610 μm diagonal moves. Top: atom survival for varied times, fitted to a clipped Boltzmann distribution. 1.6 ms moves are used for the middle and bottom panels. Middle panel: IRB return probability for static and transported atoms. Curves are fits that include coherence and atom losses. Bottom panel: average instantaneous transport fidelity after a given number of moves, fitted from the IRB return probability. The curve width represents the 68% confidence interval. The instantaneous fidelity of 99.953(2)% is constant for the first ~ 30 moves.

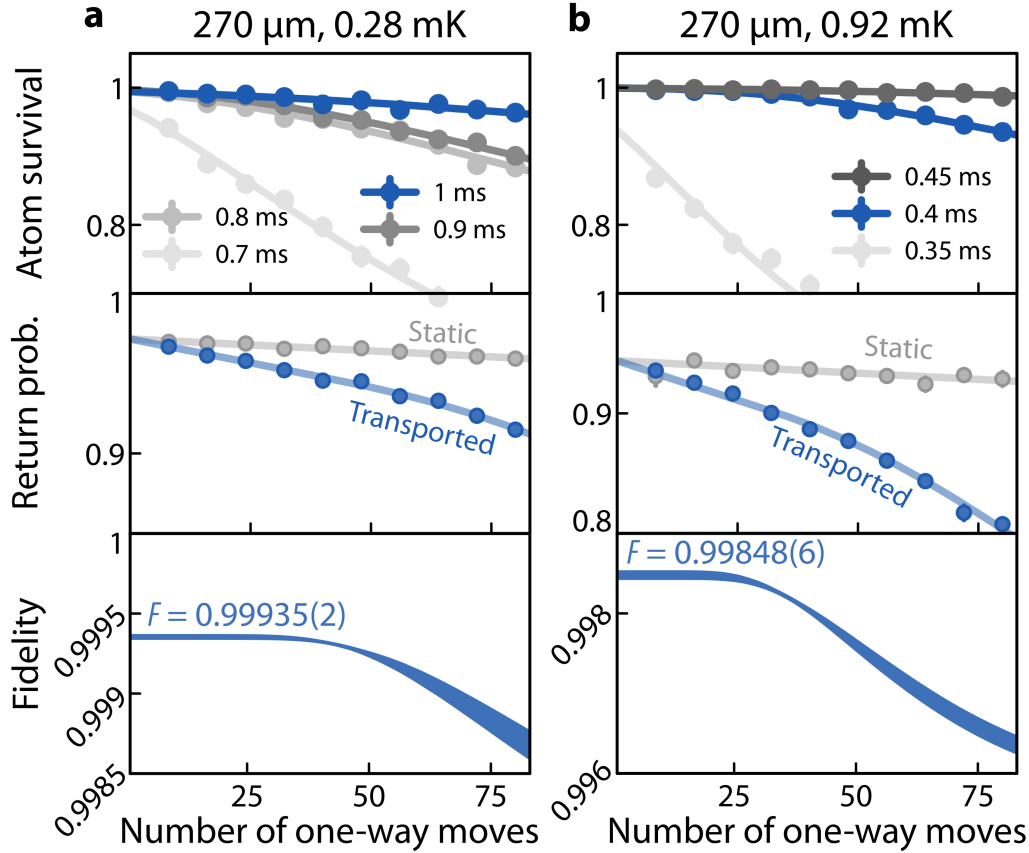


Figure 4.6: **270 μm AOD move survival and coherence.** Results of the randomized benchmarking of transport, for 270 μm at two different trap depths, as specified above each subfigure. Similarly to Fig. 4.5 we present the atomic survival for various move durations (top panel), the IRB return probability for the specific duration—1 ms and 0.4 ms respectively—highlighted in blue (middle panel), and the extracted instantaneous transport fidelity. The curve width in the bottom panel represents the 68% confidence interval.

We perform this benchmarking technique for a movement distance of 610 μm (Fig. 4.5), with diagonal moves, after loading directly from the MOT into ten AOD tweezers. We first measure the survival probability of an atom in a tweezer at the end of the sequence for different move durations (top panel). For a 1.6-ms move using $k_B \times 0.28$ mK-deep traps, we then characterize the return probability to the initial quantum state after IRB as a function of the number of moves (middle panel).

The resulting IRB return probability data is non-exponential in the number of moves, because at large numbers of moves, trap losses become dominant and the fidelity for the transport channel depends on the number of previously executed moves. We thus measure the instantaneous fidelity of the transport channel after a certain number of

previous moves, shown in the bottom panel of Fig. 4.5. The instantaneous fidelity approaches a constant value of 99.953(2)% for small numbers of one-way moves ($\lesssim 30$), where losses are the sub-dominant error. This regime is most relevant for QEC, since data qubits and ancilla qubits can, in principle, be swapped every few layers of gates [154].

We additionally extract the fidelities and survivals using randomized benchmarking at a distance of 270 μm , which is approximately the average distance that atoms would have to move between the storage zone and either the interaction or readout zones in our proposed layout for zone-based computing (section 4.3.1). It is also approximately the average move distance for rearrangement to a compact configuration (section 4.3.2). We measure the fidelity of the quantum channel both for a trap of 0.28mK as well as for a trap depth of 0.92mK. We find an instantaneous-move fidelity with 0.28mK trap depth of 99.935(2)% and an instantaneous move fidelity with 0.92mK of 99.848(6)%. The deeper trap depth allows for faster moves which could be used to enable a faster timescale for rearrangement without loss, but which incurs more infidelity in the case of a coherent move, due to scattering rates from the deeper trap.

4.2.3 SLM-AOD transfer

Until now we have isolated the effect of movement on coherence in absence of the static SLM tweezer array that we have used for the rest of this work. We now implement the necessary operations of transferring atoms between SLM static tweezers and the dynamic AOD tweezers. In this section we characterize purely the transfer operation for 195 AOD tweezers, and in the next section, we investigate the transfer operation in conjunction with a 375 μm move.

We first demonstrate that these operations proceed without the emergence of unexpected technical challenges by performing high-fidelity parallel SLM-AOD transfer across the full field of view of the AOD (Fig. 4.7).

We use 195 AOD tweezers spread across 504 $\mu\text{m} \times 468 \mu\text{m}$ (Fig. 4.7a) to perform and characterize a repeated transfer procedure, post-selected on initially filled SLM sites. As with coherent transport benchmarking, we evaluate the transfer fidelity as a function of the number of one-way transfers through IRB (Fig. 4.8). To execute faster (or higher-fidelity transfers at a given duration), we propose and implement a trajectory where AOD ramp-and-move operations are simultaneously optimized with machine learning techniques to maximize survival (Fig. 4.7c). Compared with

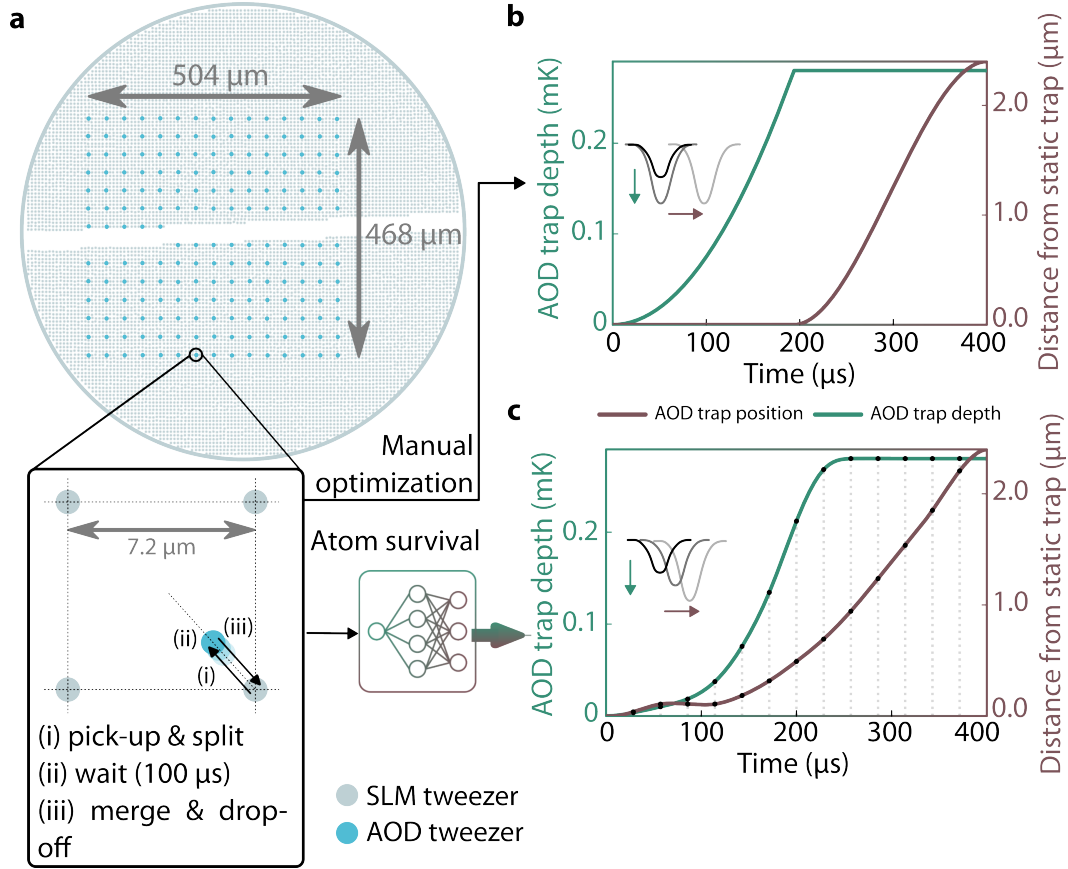


Figure 4.7: Machine-learning optimized SLM-AOD transfer. **a**, Layout of the transfer experiment showing 195 dynamic AOD traps (bright blue) overlapped with 1061 nm static SLM traps (pale blue). Atoms are repeatedly picked up and moved away by 2.4 μm , then held for 100 μs . During this time, the SLM traps are turned off to ensure that atoms left behind in SLM traps are removed (this way, atom survival correctly corresponds to a successful pick-up and drop-off). Then SLM traps are subsequently turned on, and atoms held in AOD tweezers are moved back and dropped off into the SLM traps. For IRB data shown in Fig. 4.8, gates are interleaved between each round-trip transfer. A pick-up and split-move operation (or equivalently a merge-move and drop-off operation) are considered a ‘one-way transfer’. **b**, Best hand-optimized trajectory for trap transfer, using a quadratic depth profile and a constant jerk movement. Here we implement the pick-up and the tweezer separation move in sequence, without overlap. **c**, To speed up atom transfer between static and dynamic traps while preserving high survival, we optimize, via machine learning, a trajectory where dynamic AOD traps are simultaneously ramped and moved. The dashed lines and black dots represent the values that are optimized by the algorithm.

our manually optimized trajectory (Fig. 4.7b), this technique yields significantly higher atomic survival, and enables a one-way transfer fidelity of 99.81(3)% for

$\lesssim 12$ transfers (compare blue and gray data in Fig. 4.8, top panel).

To transfer atoms between static and dynamic traps, we generate an evenly-spaced grid of 15×14 AOD tweezers (with a spacing five times that of SLM sites, Fig. 4.7a), of which 195 sites overlap with SLM tweezers generated with the 1061 nm tweezer laser (out of 11,397 sites in the SLM array). The focal planes are matched by imprinting a Zernike defocus polynomial using the SLM. The position of each SLM site is adjusted in the WGS algorithm to match the corresponding AOD site, first by matching the point-spread function on the qCMOS camera, and then by optimizing the transfer survival. For the data shown in Fig. 4.8, the SLM trap depth is ramped down to $\sim k_B \times 0.14$ mK, which we find is optimal for transfer into $k_B \times 0.28$ mK-deep AOD tweezers. We note that adiabatic ramping between full depth and this depth does not incur noticeable losses.

For hand-optimized trajectories shown in Fig. 4.7b, the AOD trap depth is quadratically increased over the course of 48% of the total ramp-and-move duration, after which the AOD trap is moved with a constant jerk trajectory by $2.4 \mu\text{m}$ during the remaining 52%. These ratios, as well as the ramp and trajectory used, are set to empirically maximize atom survival.

As an alternative, we propose and implement a machine-learned procedure for faster (or equivalently, higher-survival) atom transfer, for which the AOD trap depth and position can be simultaneously changed (Fig. 4.7c). For both trap depth and position, 14 points—from which ramps are obtained by cubic interpolation—are adjusted by a machine learner [155] for a fixed one-way duration of $400 \mu\text{s}$ and 60 consecutive one-way transfers. This trajectory is inverted to merge and drop off atoms back into static traps.

For data shown in Fig. 4.8, AOD tweezers are repeatedly ramped up and moved away from the corresponding SLM sites by $2.4 \mu\text{m}$, and then held static for $100 \mu\text{s}$. The direction of motion is as pictured in Fig. 4.7a, and does not match the direction of transport used in Fig. 4.9 because cylindrical lensing is not detrimental at the speeds being reached for the short move in this pickup-only operation. During the $100 \mu\text{s}$ wait time, SLM tweezers are turned off after which they are turned back on, such that atoms held in SLM tweezers and not AOD tweezers are dropped. This enables us to ensure that atoms that may have remained in traps rather than being successfully picked up, are not counted towards survival. At the end of the sequence atoms are imaged again in SLM tweezers. We use the same dynamical decoupling sequence as for the AOD-only transport experiment, including notably the transformed Clifford

frame technique. Unlike for long-distance transport, we find that the survival as a function of the number of transfers has an exponential component—likely due to experimental imperfections such as trap misalignment. Hence, we fit it to:

$$S_n = p_0 p^n (1 - e^{-\frac{1}{bn}}). \quad (4.16)$$

In order to accurately distinguish between depolarizing effects and atom loss in the IRB return signal, we fit the return probability as $R_n = S_n D_n$ where $D_n = p'_0 (p')^n (1 - e^{-1/b'n})$. The fidelity per move is then extracted as

$$F_n = \left(\frac{1}{2} + \frac{S_{n+1}}{2S_n} \right) \frac{D_{n+1}}{D_n}. \quad (4.17)$$

The uncertainty is obtained by first bootstrapping fitting parameters for S_n , and then, for each sample, by bootstrapping those for D_n . Unlike for AOD-only transport, the choice of convention used to account for losses impacts the early-time instantaneous fidelity. In a scenario where losses can be directly detected, one could assign an infidelity from loss equal to half the loss probability—as in a depolarizing model. In this scenario, with three state readout including loss, the instantaneous fidelity quoted in Fig. 4.8 would increase from 99.81(3)% to 99.88(3)%.

4.2.4 AOD-SLM transfer and move

After characterizing the two steps separately, we now combine SLM-AOD atom transfer with long moves in the setting of a large SLM array where we additionally seek to maintain the survival and coherence of atoms in the static array while moving the atoms.

In order to combine atom transfer and long moves, we change the static tweezer configuration to one featuring alternating spacing as shown in Fig. 4.9. This configuration is motivated by the compatibility with diagonal motion (as schematized in Fig. 4.14), and by the observation of additional losses in the absence of wider pathways for transport. We note that we minimized this loss by adjusting the timing between the horizontal and vertical AOD frequency sweep. If the beams are slightly offset on the AOD crystal relative to the respective RF transducer, then the desired diagonal sweep is not realized unless this is accounted for with an additional delay, and the beam may sweep an arc that interferes with the survival of atoms in static traps near the diagonal path. Despite optimizing this delay to minimize loss, we still find the need to increase the distance to reach acceptable loss rates.

We did not attempt to optimize SLM parameters in the original configuration to mitigate these losses and leave further investigation to future work. We include a

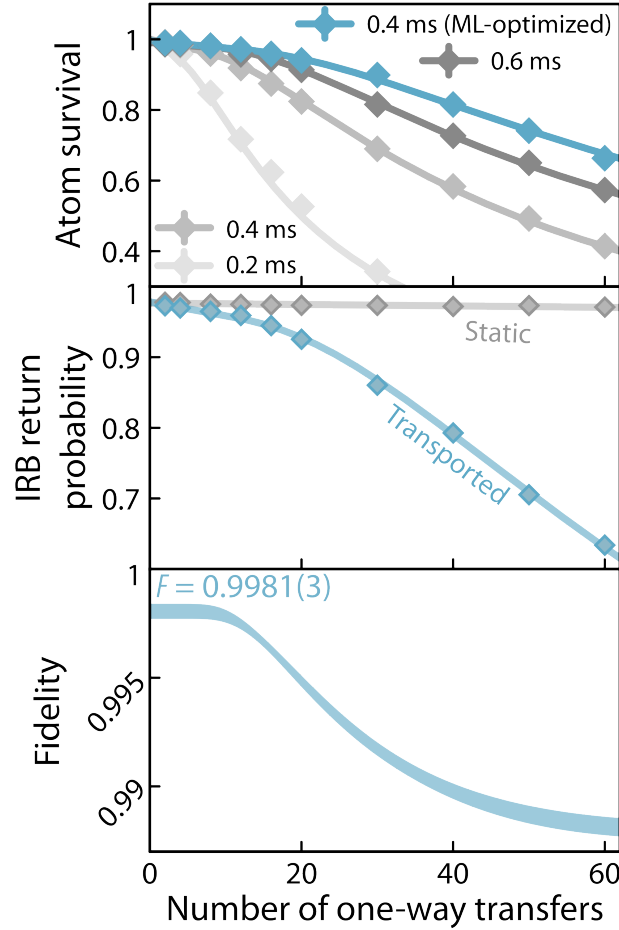


Figure 4.8: **Interleaved randomized benchmarking for SLM-AOD transfer.** Top panel: atom survival as a function of the number of repeated one-way transfers for various one-way ‘pick-up & split’ total durations. In the case of hand-optimized trajectories, the AOD trap depth is quadratically increased 48% of the total ramp-and-move duration, and the AOD trap is moved with a constant jerk trajectory by $2.4\ \mu\text{m}$ during the remaining 52%. A $400\ \mu\text{s}$ trajectory is optimized through machine learning. Middle panel: return probability after IRB for the machine-learning optimized trajectory. Bottom panel: extracted instantaneous fidelity of a coherent one-way transfer as a function of the number of previous one-way transfers.

simulation for the out-of-plane interference for this spatial distribution of tweezers (alternating between $4.8\ \mu\text{m}$ and $9.6\ \mu\text{m}$ spacing) in Fig. 2.19. Apart from a slightly lower imaging fidelity—which, in the context of fault-tolerant quantum computation, matters much less than, e.g., for quantum simulation—we do not expect this array to exhibit different metrics from the configuration characterized in the rest of this paper. We utilize 8×6 AOD tweezers (of which 47 are overlapped with one of the 11,416-site $1061\ \text{nm}$ tweezer array) spanning $285\ \mu\text{m} \times 204\ \mu\text{m}$ to demonstrate

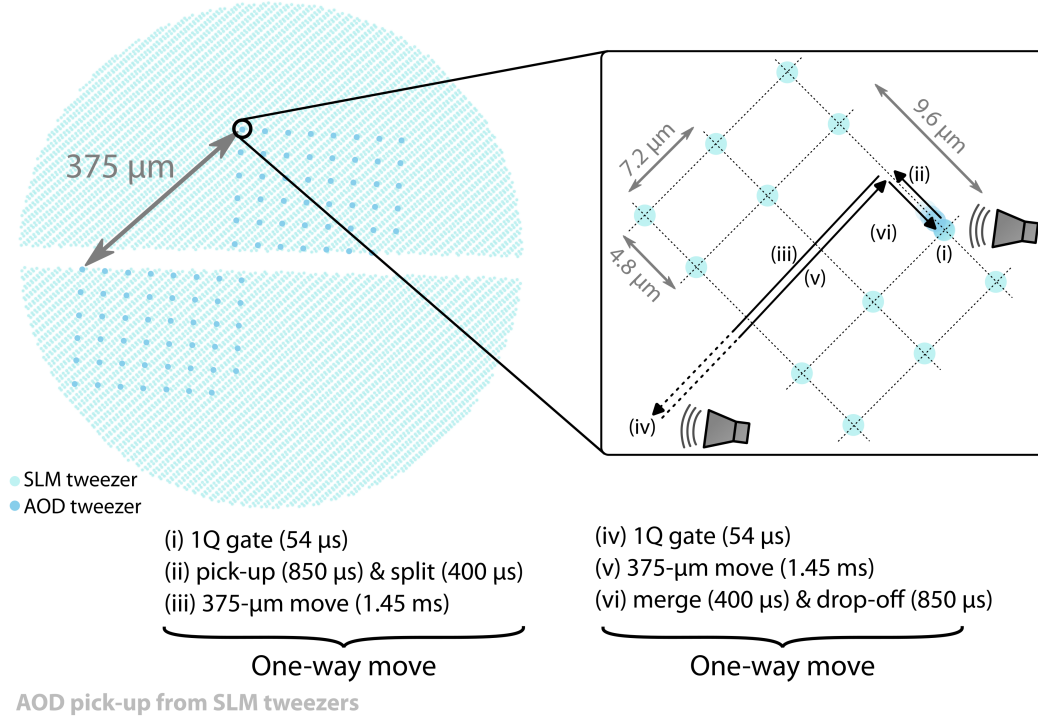


Figure 4.9: AOD-SLM transfer and long-distance transport. Schematic representing the configuration and operations used for coherent transfer and transport of atoms across 375 μ m using 47 AOD tweezers. The SLM layout (totaling 11,416 sites) alternates narrow column spacings of 4.8 μ m and wide column spacings of 9.6 μ m, between which AOD tweezers are moved diagonally. Timescales for each operation composing one-way moves are detailed in the figure.

pick-up as well as 375 μ m movement, spanning the full AOD field of view.

The combined transfer and move sequence is realized as follows (Fig. 4.9): we first apply a single-qubit gate, pick up atoms from the highlighted sites on the top side of the array, and then perform a constant jerk movement for the initial separation move. We then implement the 375 μ m-move using an adiabatic sine trajectory, apply a single-qubit gate while atoms are held in AOD tweezers on the bottom side of the array (shown in highlighted locations), before applying the reverse move and transfers. Timescales for each operation are shown in Fig. 4.9. During the pick-up operation, AOD tweezers are ramped up from 0 to $k_B \times 0.28$ mK while SLM tweezers are ramped down from $k_B \times 0.18$ mK to $k_B \times 0.06$ mK (the trap depth used for the measurement of coherence times in section 3.2). Possible deleterious effects from the repeated ramps on static atoms are captured by the ‘static’ data in Fig. 4.10: the equivalent idle fidelity is $> 99.96\%$. We measure an early-time instantaneous fidelity for a pickup and move of $99.87(1)\%$ for the timescales used here. This is

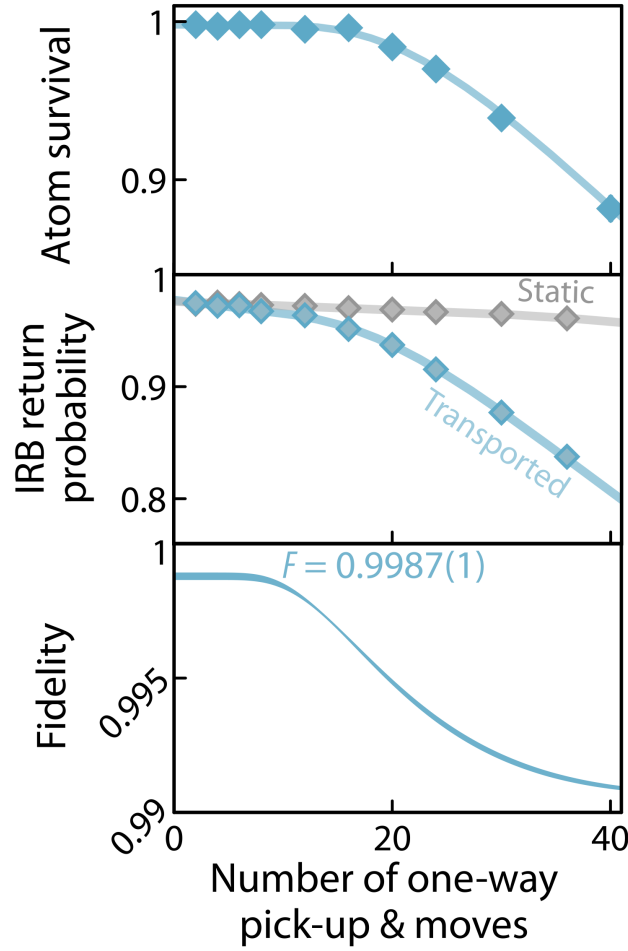


Figure 4.10: **IRB for AOD-SLM transfer and long-distance transport.** Results of the randomized benchmarking of coherent transfer and transport. Top panel: atom survival as a function of the number of repeated one-way pick-up and split, followed by a 375 μm move. Middle panel: return probability after IRB. Bottom panel: extracted instantaneous fidelity of a coherent transfer, split and 375 μm move, as a function of the number of previous such operations.

constant for about the first ten moves.

We notice no significant exponential component in the survival signal. Therefore, when we evaluate the instantaneous fidelity using the technique described in section 4.2, the early-time estimate of instantaneous fidelity is not impacted by the choice of convention for handling loss. We anticipate that the timescales used here can be sped up by leveraging machine-learning to optimize various trajectories and ramps, as demonstrated in Figs. 4.7, 4.8 . Additionally, we envision integrating the short move to split (and merge) AOD and SLM tweezers with the longer move in a

single, curved trajectory.

4.3 Proposed 6,100 qubit zone architecture

Having now demonstrated coherent operations pertaining to dynamic movement of atoms in our tweezer array, we pivot to proposing an architecture that can use the ingredients of this thesis to implement universal quantum computation. Notably, the missing but necessary ingredient is a high-fidelity two-qubit gate. We discuss aspects related to implementing such gates in section 4.5.

We present here a comprehensive outlook on the ingredients necessary for implementing a universal quantum computer utilizing the 6,100 qubit platform we have presented in this thesis, and focusing on prerequisites that represent a challenge due to scale. For instance, we leave aside the question of mid-circuit readout in alkali atoms, for which there is an important ongoing research effort[52, 111, 156], since there is no clear field-of-view or power limitation due to the large size of the array.

While a digital quantum computer is only one of the applications of this atom array, it is the one that requires the highest degree of control, and thus the strategies laid out in this section can be also utilized to implement, e.g., an analog quantum simulator with some degree of local control.

To this end, after first presenting an overview of a quantum processing zone-architecture compatible with our qubit array, we discuss in depth data-based estimates for rearrangement timing, and coherent control at scale. In the following section we also discuss plans for implementing two-qubit gates. Notably, based on long-distance transport data compatible with rearrangement moves, experimentally implemented algorithms, and timed processes, we conclude that the array can be rearranged quadrant-wise in less than 140 ms in parallel or in less than 525 ms sequentially.

This investigation into the usability of the array for quantum computing yields insights into challenges to be expected with scaling up the array size, ultimately revealing a feasible near-term pathway to large scale universal quantum computation with $\sim 6,000$ qubits.

4.3.1 Zone Architecture overview

To execute quantum operations, we envisage the usage of a zone structure [51, 157] for entanglement and readout in combination with global and local hyperfine control. There are multiple constraints that must be considered to ensure that all operations

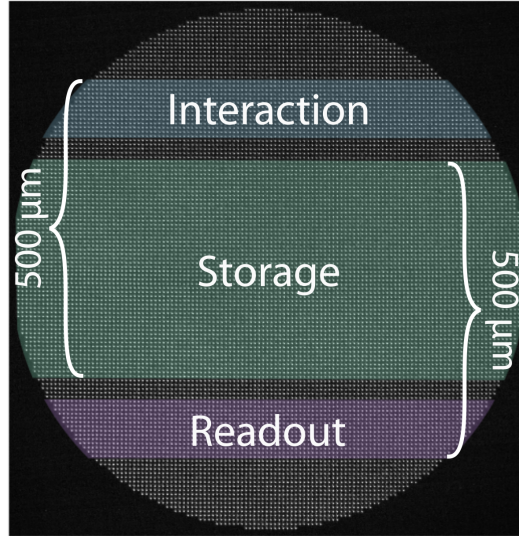


Figure 4.11: **Possible layout for a large-scale zone-based atom array quantum processor.** A proposed configuration for a zone-based architecture, shown on the 12,000-site tweezer array for scale. The storage zone contains enough sites to hold over 6,100 atoms, and the interaction and readout zones are within 500 μm of any atom in the storage zone, within the field of view of the AODs we use for movement.

can be realized without fidelity loss and with reasonable timescales. We briefly review these constraints in this section, and carefully analyze them in the following. We plan to begin an experiment by rearranging (see the following section for details) atoms into a storage zone as sketched out in Fig. 4.11. This zone is capable of holding over 6,100 atoms with the current spacing. One could additionally imagine packing qubits more closely in this zone with tighter tweezer spacing (using different wavelengths as simulated in Fig. 2.19). A diagonal configuration, as we use for a proof-of-principle characterization of pick-up and transfer across 375 μm with 47 dynamic tweezers in a large scale static array (Fig. 4.9), could also be explored.

The interaction and readout zones, where two-qubit entangling gates, and measurement could be performed respectively, are within reach of the storage zone utilizing a set of Gooch & Housego crossed acousto-optic deflectors (G&H AODF 4085) that have an angular field of view of ~ 45 mrad and an active aperture diameter of 15 mm, resulting in an effective field of view of > 500 μm after 3:2 demagnification. We have demonstrated movement of over 500 μm using these AODs, presented earlier in this chapter.

As we demonstrated in Fig. 4.5, the movement distance of ~ 600 μm can be traversed while maintaining qubit coherence during several moves. In order for the AODs to be

accessible to every part of the array, we have realigned new optical tweezer pathways since the movement presented earlier in this chapter, as pictured in Fig. 2.20. This is based on a plan that involves overlapping quadrants in the atomic plane. We have two pairs of AODs that can address adjacent quadrants simultaneously, and we additionally have a fast-steering mirror in a Fourier plane of the 4f mapping for the dynamic AOD tweezers to the atomic plane. We are able to use this mirror to switch between different sections of the array so that the two AODs can be used to move atoms in two quadrants at a time, and then switch to the other two quadrants. The fast-steering mirror in our setup is the Thorlabs FSM75, which has a specified 5 ms switching time and position repeatability of better than 20 nm in the atom plane.

The interaction and readout zones are each separated from the storage zone by ~ 40 μm , and we plan to shape the Rydberg and imaging beams that form these zones into flat-top beams utilizing SLMs. We envision additionally applying a relative light-shift between readout and other zones in order to prevent scattered photons from causing other qubits to decohere [121, 158, 159]. The readout zone, as designated in Fig. 4.11, would have over 1,400 SLM-generated sites, and for the gate spacing necessary for two-qubit Rydberg gates in the interaction zone, we would be able to perform around 500 two-qubit entangling gates in parallel.

4.3.2 Rearrangement

Atom rearrangement using optical tweezers has been demonstrated for target array sizes of a few hundred to approximately a thousand sites [55–57], with some very recent demonstrations of 2000 – 3000 atoms rearranged via either fast-SLM operation to a checkerboard goal configuration [160] or by continuous reloading [58]. However, rearranging thousands of atoms from a randomly loaded array into a compact configuration—which is necessary for the central storage zone depicted in Fig. 4.11—has not yet been achieved and remains an outstanding challenge. The field of view of commercially available AODs are currently insufficient to cover the entirety of our array size. To address this limitation, we plan to divide the 12,000-site array into four quadrants, as shown in Fig. 4.12, and perform rearrangement within each quadrant. This can be achieved either by using multiple AOD pairs or, alternatively, a pair of AODs combined with a fast-steering mirror. In our newest iteration of the experimental tweezer pathways, we incorporate an option between these extremes: two pairs of AODs that can operate in parallel, combined with a fast-steering mirror for wider array-access.

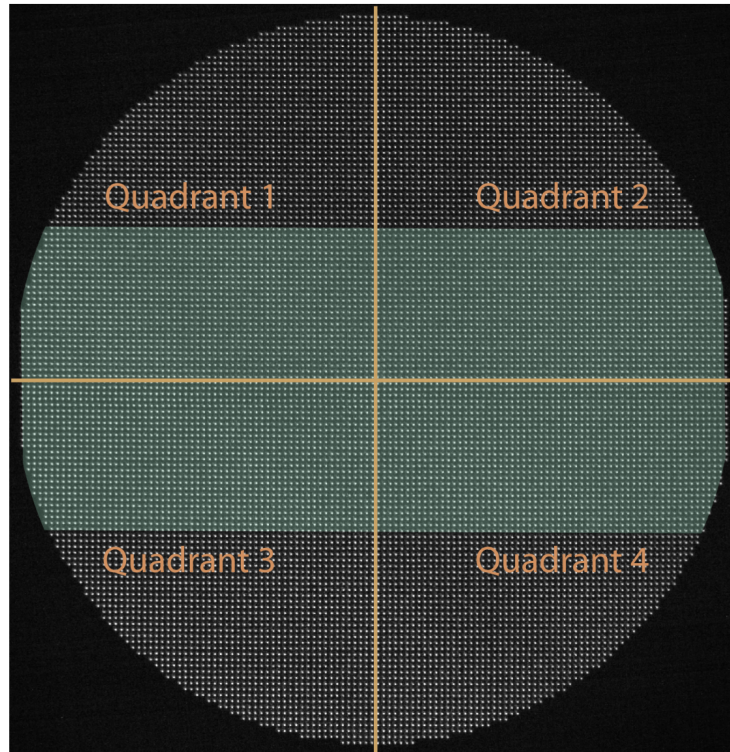


Figure 4.12: **Quadrant-based rearrangement.** A sketch of the quadrants that we propose to rearrange using large-aperture AODs. Combining pairs of AODs that cover each region will allow for parallel rearrangement. One could also rearrange quadrants in succession one after another, utilizing one set of AODs and a fast-steering mirror for switching. Our current setup (Fig. 2.20) will enable two pairs of AODs to operate at once, with a fast-steering mirror in the Fourier plane to switch to the other two quadrants for rearrangement or other operation as needed.

In this section, we estimate the time required for rearrangement based on our experimental data and simulations, demonstrating that it is both feasible and achievable within a practical timeframe.

Tetris rearrangement for each quadrant

Each quadrant requires the rearrangement of an average of $\sim 1,500$ atoms, representing a scale that has not been achieved to date for rearrangement to a compact configuration from a stochastically loaded one. A recent example of atomic rearrangement into a 828-site compact array required around 1.5 s to realize the target configuration with a single rearrangement tweezer [161]. While such a timescale would slow the experimental cycle and allow some defects to be incurred due to vacuum-induced losses, algorithms that utilize many rearrangement tweezers in parallel expedite the

process[47]. Hence, we propose to use the “Tetris algorithm” [162], a row-by-row sorting algorithm that capitalizes on parallelism and overcomes other timescale limitations through real-time calculation. In Fig. 4.13a, we show the steps of the Tetris algorithm from one iteration of our simulation for the sorting of a randomly half-filled quadrant of our 12,000 site array into 1,445 target sites.

Averaging over $\sim 1,000$ iterations of our simulation sorting one quadrant with different initial configurations, we find that the algorithm requires an average of 119 steps, with ~ 23 atoms being moved on average in each step (Fig. 4.13c). The longest distance move in the row being sorted determines the timescale of the step. In Fig. 4.13b we show the histogram of the longest distance move per step (with an average longest distance of $159 \mu\text{m}$).

Time budget for quadrant rearrangement

We can use our implementation of the Tetris algorithm in the previous section to give an estimate of the rearrangement time using conservative parameters based on the experimental data, simulation, and existing values in the literature. In Table 4.1 we present a summary of this time budget. We find that ~ 137 ms would be required to rearrange quadrants in parallel, or ~ 522 ms if the quadrants are sorted sequentially. If we choose to implement something that is partially sequential and partially parallel, such as with two AODs in parallel as we have planned in our latest upgrades, we expect the timescale to be ~ 264 ms. This timescale is compatible with the operation of a typical neutral atom experiment. Since the timescales for moves assume experimental results for straight moves, which suffer from cylindrical lensing effects, the overall time could be further reduced by mitigation strategies for cylindrical lensing outlined in section 4.3.4. In the next paragraphs we detail our rationale for each line of Table 4.1.

Image transfer and processing

We have measured a delay of 10 ms between the end of camera exposure and the determination of the atom occupation matrix. Our implementation of the image processing algorithm is written in Python, and we rely on the fast CoaXPress connection with the Hamamatsu qCMOS camera. This delay time will be the same regardless of the sorting implementation.

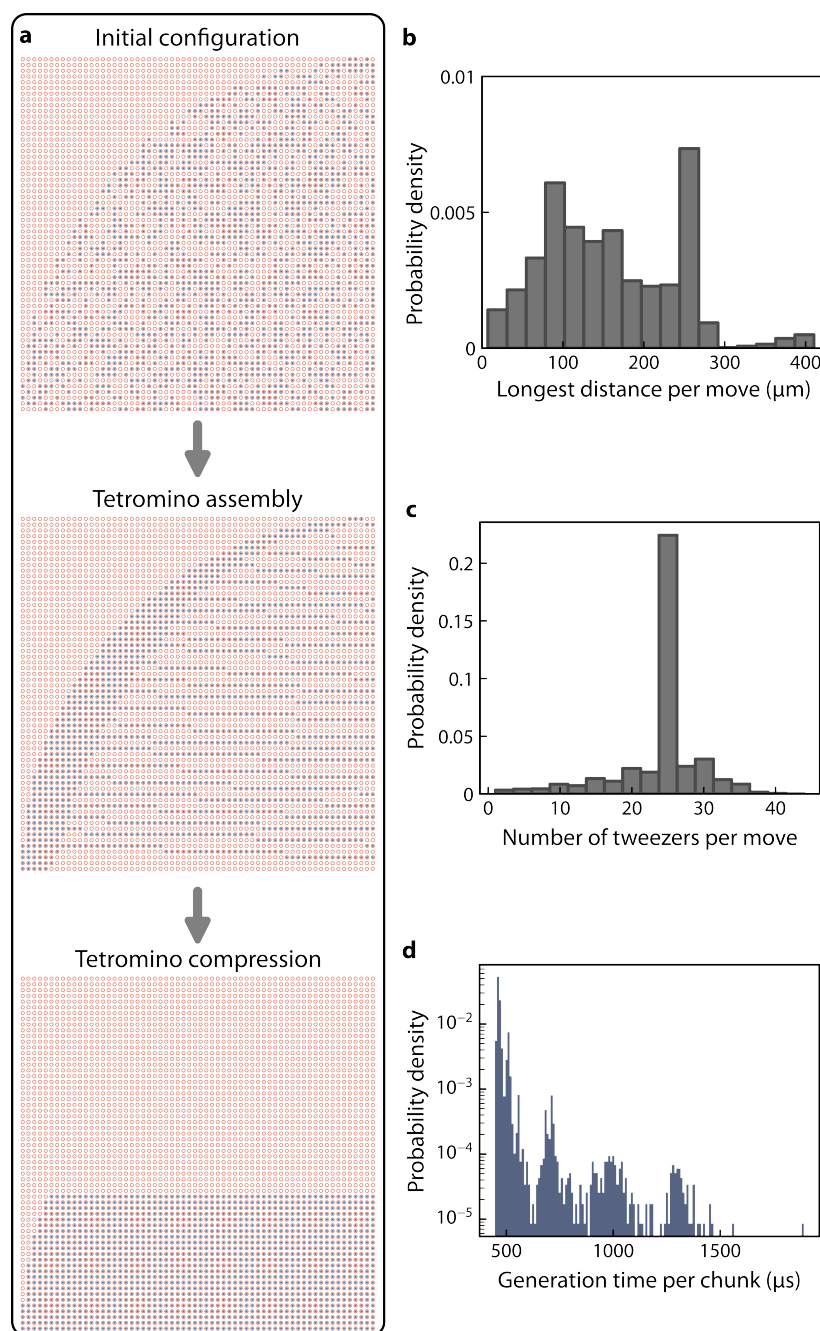


Figure 4.13: **Tetris sorting algorithm adaptation for a quadrant.** **a**, Illustration of the algorithm for a randomly selected example. **b**, Histogram of the longest distance in each row-based rearrangement move (usually composed of several atom trajectories), averaged over $\sim 1,000$ iterations of random initial configurations. This longest distance determines the move's total duration. **c**, Histogram of the number of atoms that are moved in parallel during a move. **d**, Histogram of the required computation time for each waveform chunk of 524k samples for our parallel waveform generation algorithm (with 64 parallel movements being summed). The average generation time is 488 μs , and the generated waveforms sort on average 2-3 rows.

Operation	Time	Parallel sorting	Sequential sorting	Semi-Parallel sorting
Image transfer and processing	10 ms	$\times 1$	$\times 1$	$\times 1$
Tetris algorithm computation	235 μs	$\times 4$	$\times 1$	$\times 2$
Waveform streaming latency	488 μs	$\times 4$	$\times 1$	$\times 2$
Pick-up & drop-off	24.8 ms	$\times 1$	$\times 4$	$\times 2$
Tweezer splitting & merging	24.8 ms	$\times 1$	$\times 4$	$\times 2$
Movement	74.4 ms	$\times 1$	$\times 4$	$\times 2$
Fast-steering mirror switching	5 ms	$\times 0$	$\times 3$	$\times 1$
Total		~ 137 ms	~ 522 ms	~ 264 ms

Table 4.1: **Time budget for quadrant rearrangement.** This conservative estimate, detailed in the text, is based on data presented in our work and available in the literature. We notably assume straight moves and therefore the move duration could be significantly reduced by compensating for cylindrical lensing, as detailed in section 4.3.4. We show in the table timescales for rearranging quadrants in parallel using four sets of 2D AODs (“Parallel sorting” above), sequentially using a fast-steering mirror to switch between quadrants (“Sequential sorting” above), and partially parallelized using two sets of AODs (“Semi-Parallel sorting”). The pick up and drop off timescales as well as the tweezer splitting and merging timescales are calculated as following: $(62 \text{ rows} + 62 \text{ columns}) \times 200 \mu\text{s} = 24.8 \text{ ms}$. The move time is calculated as $(62 \text{ rows} + 62 \text{ columns}) \times 600 \mu\text{s} = 74.4 \text{ ms}$.

Tetris algorithm computation

We have written an implementation of the Tetris algorithm for our array using C++. Owing to its relative simplicity, we can calculate and determine which moves to execute for rearrangement of one quadrant in about 235 μs . In the case of sequential sorting, this calculation could be done during the fast-steering mirror switching, so we only include the calculation of the first quadrant sorting in the table.

Waveform streaming latency

We turn to estimating the latency associated with waveform generation and streaming to the AODs. Results in the literature vary depending on the hardware and the implementation. Examples include synthesis from a computer-controlled arbitrary waveform generation card [31] or generation from a field-programmable gate array equipped with a digital-to-analog converter [40, 162]. For simplicity, one may store in memory all the waveforms associated with moving a single trap [33], and during the experiment sequence sum the appropriate waveforms. For the atom array introduced in this work, we estimate that we would need to store 2.53 million waveforms. Assuming a realistic sampling rate and average waveform duration,

this represents over 250 billion samples, or more than 1 terabyte of data (with single-precision floating-point numbers). One would need to sum these waveforms on the fly in addition to storing this immense amount of data, which is essentially prohibitive to the practicality of this approach.

Our approach provides an alternative to storing waveforms. We write and benchmark a real-time waveform generation algorithm from scratch on a consumer graphics processing unit (GPU, NVIDIA GeForce RTX 3090), leveraging the device's high calculation throughput using the C++ CUDA platform. The waveform samples are grouped in "chunks" for streaming, while the subsequent chunk is generated. The latency is defined as the time that it takes to calculate the first chunk before the waveform generation is simultaneous to streaming. The chunk size of 524k samples is chosen, which on average covers the waveforms for sorting 2-3 rows of a quadrant using 64 moving tweezers (slightly more than the number of rows in a quadrant). This amount of data represents about 2.6 ms of waveform streaming, and the time it takes to generate each chunk is on average $\sim 488 \mu\text{s}$ (the full histogram is shown in Fig. 4.13d).

Since the generation time is shorter than the waveform durations, one can stream the first waveform chunk to the arbitrary waveform generator, compute the second chunk and stream it before the first one ends, and so on. This enables continuous streaming with latency given by the generation time of the first chunk, on average $488 \mu\text{s}$.

In the time budget in Table 4.1, for the cases including sequential sorting, we count this latency time for the first quadrant(s) being sorted, given that after the moves are generated for the first quadrant(s), this calculation can be done during the fast-steering mirror switching time of 5 ms. For the case of parallel sorting, we count this latency four times as a conservative estimate, although this could be also implemented in parallel by utilizing additional workstations or GPUs.

Pick-up & drop-off , splitting & merging

The times needed to pick up and drop off an atom from a static to a transport tweezer, and to split and merge the static and transport tweezers are not dependent on the distance of the rearrangement move being performed. In Fig. 4.7 we present data showing atom survival for repeated (coherent transfer): we notably find that a single one-way transfer can be done in $200 \mu\text{s}$ with a survival probability of more than 98%. Moreover, higher AOD and SLM trap depths could be used to further speed up

this operation—since for rearrangement we are not concerned with qubit coherence. Hence, we pick this timescale as a reasonable assumption for rearrangement. As is typical for smaller-scale experiments, one could perform a second, much shorter rearrangement round to further increase the success probability. The Tetris algorithm first implements row-based rearrangement of atoms, with one pick-up, split, merge, and drop-off operation for every row, and then compresses atoms in columns with one pick-up, split, merge, and drop-off operation per column: hence, we multiply by (62 rows + 62 columns) in the Table.

Movement

Regarding transport itself, it would be more difficult to extrapolate from previous results since there are important trade-offs associated with long-distance moves, as discussed in the previous sections. The data that we present in Fig. 4.1 gives survival statistics for two different distances and various trap depths for straight moves. From this data, we select the average transport time, 600 μs , as the transport time required to obtain atom survival $\gtrsim 0.999$ using a tweezer with depth 0.28 mK for a distance of 200 μm , which is slightly further than the average longest distance per set of parallel moves found by simulating the Tetris algorithm. Note that this timescale could be sped up by methods for mitigating cylindrical lensing as discussed in the next section, with a deeper trap depth, or with machine learning, but we choose to base our estimate off of this non-optimized value.

Fast-steering mirror switching

As mentioned above, we have incorporated a Thorlabs fast-steering mirror in our updated optical pathway, with a switching time of 5 ms. We include this time in our estimation of the sequential sorting sequences which require the field of view of the AODs to switch between quadrants.

4.3.3 Coherent transport between zones and hyperfine control

Coherent transport of atoms [145] during the quantum processor operation is a key tool that enables all-to-all connectivity with low overhead in recent experiments that implement a universal quantum processor with an atom array [35, 51]. While there are some challenges associated with scaling up the size of the array, we demonstrated in our work that high-fidelity moves can be realized even for a distance of 610 μm as shown in Fig. 4.5. Importantly, this distance is greater than the largest distance needed to move atoms between storage and interaction or readout zones in our

proposed layout. Thus far, we found that we are able to move much more quickly with a diagonal motion when compared with a straight motion (Fig. 4.2). We have also examined coherent trap-transfer between static SLM traps and dynamic AOD traps at the scale of 195 tweezers across $504 \times 468 \mu\text{m}$ (Fig. 4.7), and combined this with a the long-distance coherent moves that we demonstrate, by characterizing the fidelity of trap-transfer for 47 tweezers, followed by a $375 \mu\text{m}$ move between the upper and lower half of a large-scale static array of tweezers (Fig. 4.9, Fig. 4.10). Further development of the fidelity and speed of this operation is expected with the use of machine learning techniques demonstrated in Fig. 4.7. These data together, we believe demonstrate the capacity to utilize coherent movement between zones in our proposed architecture, as is necessary to utilize our atomic array for quantum information processing. Our data, however, rely on diagonal moves in order to compensate for cylindrical lensing. In the next section, we discuss other methods for mitigating lensing effects as well as how diagonal moves could be used for comprehensive movements across the array.

4.3.4 Cylindrical lensing mitigation for fast coherent transport

In this section, we highlight several directions to mitigate the deleterious effects due to cylindrical lensing. Implementing such techniques would enable faster high-fidelity transport, and allow for more effective dynamical decoupling by maintaining a consistent trap depth during motion. Fig. 4.14 Essentially, there is no drawback associated with configuring transport tweezers in a diagonal way, aside from the need to conceive blocks of atoms — e.g., a block encoding a logical qubit — in a diagonal way. We have already demonstrated a preliminary such diagonal orientation of static tweezers in Fig. 4.8 which we use for characterization of the fidelity of combining coherent pick-up with a $375 \mu\text{m}$ move.

Importantly, this setup, allowing arbitrary moves for dynamic tweezers, requires three AODs. For instance, one may use two orthogonal AODs, AOD_x and AOD_y , to move in the vertical direction, such that along this direction $\dot{f}_x = \dot{f}_y$, as we have done in order to demonstrate the coherent move data. If this configuration was used to move tweezers along the horizontal direction, then $\dot{f}_x = -\dot{f}_y$, which would *maximize* detrimental effects from cylindrical lensing. To remove lensing effects for movements in the vertical-horizontal grid, another AOD is required, with a parallel opposite deflection axis with respect to one of the two other AODs.

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effects due to cylindrical lensing. Implementing such techniques would enable faster high-fidelity transport, and allow for more effective dynamical decoupling by maintaining a consistent trap depth during motion.

Diagonal motion

To mitigate the effects of cylindrical lensing in the data that we present in Figs. 4.5, 4.6, and 4.7, we have used diagonal moves (by which we mean that tweezers move with the same rate along x and y ; $\dot{f}_x = \dot{f}_y$), which turns cylindrical lensing into spherical lensing. While there is still a penalty associated with the focal shift of the tweezer, it is much reduced compared with the straight move situation. We have leveraged this strategy to realize faster high-fidelity moves, as presented in Figs. 4.2, 4.5, 4.6, 4.10.

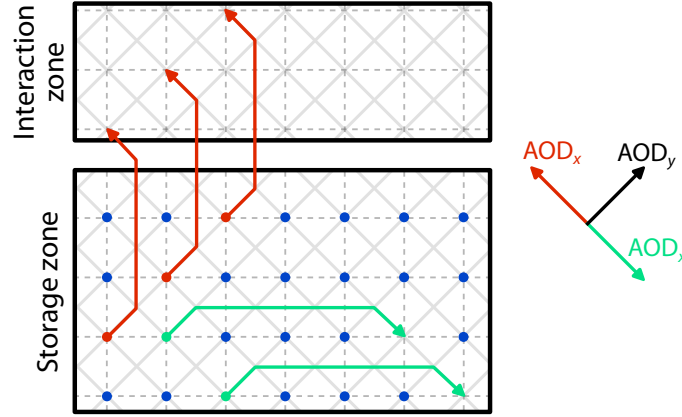


Figure 4.14: **Diagonal-move based rearrangement** Schematic of a zone layout that is compatible with diagonal motion with AODs. Each arrow associated with a distinct AOD indicates the diffraction axis in the direction of increasing acoustic frequency. To realize the vertical (red) set of moves, one would use AOD_x and AOD_y , and for the horizontal (green) set, $AOD_{x'}$ and $AOD_{y'}$.

In Fig. 4.14 we describe how a zone-based neutral atom quantum processor can be engineered with these diagonal moves. Essentially, there is no drawback associated with configuring transport tweezers in a diagonal way, aside from the need to conceive blocks of atoms — e.g., a block encoding a logical qubit — in a diagonal way. We have already demonstrated a preliminary such diagonal orientation of static tweezers in Fig. 4.9 which we use for characterization of the fidelity of combining coherent pick-up with a $375\ \mu\text{m}$ move.

Importantly, this setup, allowing arbitrary moves for dynamic tweezers, requires three AODs. For instance, one may use two orthogonal AODs, AOD_x and AOD_y , to move in the vertical direction, such that along this direction $\dot{f}_x = \dot{f}_y$. If this configuration was used to move tweezers along the horizontal direction, then $\dot{f}_x = -\dot{f}_y$, which would maximize detrimental effects from cylindrical lensing. To remove lensing effects for movements in the vertical-horizontal grid, another AOD is required, with a parallel opposite deflection axis with respect to one of the two other AODs.

Cylindrical lensing suppression with counter-propagating wave AODs

An alternative solution that has been proposed to strongly suppress acousto-optic lensing is to use two acousto-optic deflectors (for each dimension) stacked so that the acoustic wave in the second AOD is identical but counter-propagating with respect to the first AOD [163]. By diffracting the tweezer light as the respective (+1) and (−1) orders of each AOD, the diffraction angles add up but the lensing effects cancel to first order. Compared with diagonal motion, this presents the added benefit of removing any lensing (including spherical). There have been some recent works that utilize such approaches and provide demonstrations [164]. However these works, as well as recent tests that we have implemented, find that the diffraction efficiency for the pair of AODs drops fairly quickly, such that the usable bandwidth for the pair is significantly less than would be usable with just one AOD. Our understanding is that this effect arises from the birefringent nature of the AOD crystal diffraction which allows for an optimized output diffraction bandwidth, at the expense of a narrow range of input angles [165].

Acousto-optical deflectors with faster access times

As discussed in the previous section, the characteristic shift velocity for a given AOD and optical configuration is given by $v_s/w_0 \sim 5.3/T_a$, where w_0 is the tweezer waist and T_a the *access time* (or *time aperture*) of the AOD. T_a factors in a crucial figure of merit of wideband AODs, the time-bandwidth product $T_a \Delta f$, which describes up to a unit-scale factor the number of resolvable spots [149]. Therefore the parameter space to decrease T_a while preserving the number of resolvable spots and the scan angle is highly constrained. Since all commercially available AODs at near-infrared wavelengths currently use the shear mode of tellurium dioxide (which possesses an unusually low acoustic velocity), for which further improvement in the figures of merit mentioned above seem unrealistic, it appears that further research in acousto-

optic materials would be required to fabricate wide-band AODs at near-infrared wavelengths with low access time.

4.4 Qubit hiding

In order to image atoms in the readout zone without decohering qubits that are in the storage zone due to scattered photons, there are a few possible approaches. Recently developed methods utilizing dual species arrays [166, 167] will allow for readout of one atomic species, used as measurement qubits, that have different transition wavelengths and thus scatter photons that are very far off-resonant with the transitions in the data atoms. Other approaches include utilizing atoms such as Yb [168] that have states that can be used for shelving qubits while others emit photons at cycling wavelengths [169–171]. A final approach that we have explored and that has recently been demonstrated [58, 77] is to use very strongly resonant light to cause Autler-Townes splitting with enough magnitude such that atoms are no longer resonant with the imaging light or scattered photons.

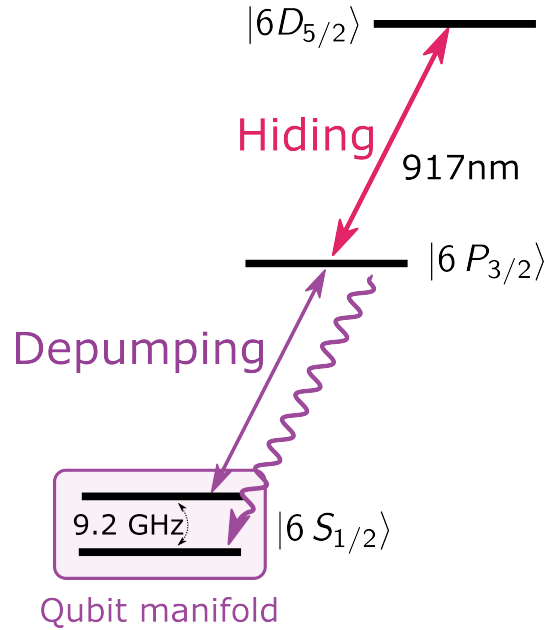


Figure 4.15: **917 nm Hiding transition.** We picture the 917 nm hiding transition, demonstrating how it shifts the D_2 imaging transitions to protect relevant qubits. There are other options for hiding transitions from $6P_{3/2}$ to other $D_{5/2}$ states including to $7D_{5/2}$, where the wavelength of interest 698 nm.

To protect qubits from depumping out of the qubit state due to incident light, we utilize the 917 nm transition from $6P_{3/2} \rightarrow 6D_{5/2}$ in order to split the $6P_{3/2}$ state. There are other transitions in Cs that could also be used for atom hiding in a similar

manner, such as the $6P_{3/2} \rightarrow 7D_{5/2}$ transition which has a wavelength of 698 nm. For each transition, there is a tradeoff between decoherence due to scattering from the hiding laser itself and protection from incident resonant photons from imaging, due to the magnitude of the induced splitting. There is, then, an optimal intensity that balances these effects in order to optimize the coherence time.

We do an experiment to explore the depumping due to direct imaging light that can be prevented, by focusing a Gaussian beam on our array such that we have a gradient of intensities that cause different splittings. When we post-select for the atoms at the center of the Gaussian beam where the intensity is around $\sim 50\text{W}/\text{mm}^2$, we observe around 1 GHz of splitting. We measure depumping for atoms across the array that are protected to various degrees from incident resonant D_2 light with intensity of $\sim 0.1 \times I_{\text{sat}}$.

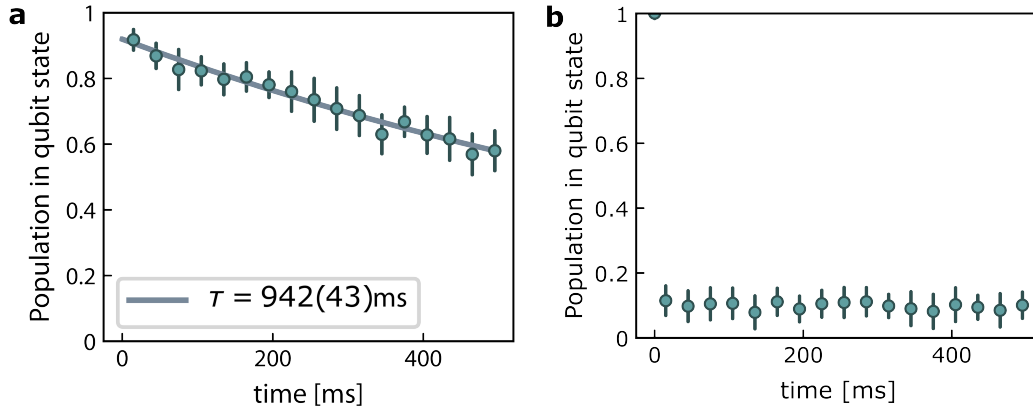


Figure 4.16: **Single-qubit hiding.** Population remaining in the $|1\rangle$ state after depumping with incident, resonant D_2 light with intensity $\sim 0.1 \times I_{\text{sat}}$ in the case where we **a**, hide several hundred qubits in the center of the array with 917 nm light with $\sim 50\text{W}/\text{mm}^2$ intensity or **b**, don't hide the atoms. We fit an effective T_1 time in this scenario of around 942(43)ms, taking into account the offset due to the average final population in the qubit $|1\rangle$ state after depumping. We note that in this data we have not properly characterized the depumping time in the case without hiding due to the difference in timescale.

4.5 Two-qubit gates

To implement two-qubit gates on the neutral atom quantum computing platform, the standard approach uses excitations to Rydberg states, or states with high principal quantum number. Such highly excited states cause the effective atomic size to be on the scale of hundreds of nanometers, thus allowing atoms that are spaced by several microns to interact. In particular, if an atom is excited to a Rydberg state, there is a

Rydberg blockade [172, 173] arising from the van der Waals interaction³ (scaling as $1/R^6$, where R is the distance between two atoms) that shifts the energy of the doubly excited state, such that a second atom within the *blockade radius* of the first cannot simultaneously be excited. This interaction then has been used to generate entanglement [29, 174, 175].

Early proposals for utilizing Rydberg-based entanglement generation to implement two-qubit gates involved individual control over two atoms [28], such that a control atom could be excited to the Rydberg state via a π pulse if it was in one of the qubit states and the excitation would be off-resonant if the control atom were in the other state. Then a 2π pulse could be applied to the target atom, and then another π pulse applied to the control. Such a scheme would apply a π phase shift to the target qubit if it had been excited, and no phase shift if it had been blockaded.⁴ Such a scheme clearly implements a CZ gate, resulting in entanglement if at least one atom begins in a superposition of states.

An excitation protocol that allows for entanglement to be engineered while simultaneously exciting multiple atoms was additionally proposed [176]. This relies on the fact that when N atoms in the ground state $|g\rangle$ are illuminated with Rydberg excitation light to excite to state $|r\rangle$ simultaneously, they undergo Rabi oscillations to an entangled excited state $|W\rangle = \frac{1}{\sqrt{N}}(|rgg\dots g\rangle + |grg\dots g\rangle + |ggr\dots g\rangle \dots + |ggg\dots r\rangle)$, since multiple atoms cannot be simultaneously excited. This excitation then has a Rabi frequency given by $\sqrt{N}\Omega$, where Ω is the Rabi frequency of the single-atom excitation.

Since early implementations, other gates have been developed that can be used with global control rather than local to engineer a CZ gate [177], and that can be implemented more quickly [178]. In these cases, both atoms are excited simultaneously, coupling to the two atom $|W\rangle$ state, but with Rydberg pulses that are engineered such that a π phase is acquired if both atoms are in the $|1\rangle$ state, and no phase is acquired otherwise, up to a global phase. These pulses can be engineered by adjusting the phase of the Rydberg laser (for a two-photon excitation only one of the lasers must be controlled).

³This interaction arises from the second-order effect of the $1/R^3$ dipole-dipole interaction, since for nS Rydberg states, there is no average dipole moment. However, atoms can exhibit first-order dipole-dipole interactions when they are prepared in Rydberg states with opposite parity (such as a resonant exchange interaction between atoms in an nS state and an $(n+1)P$ state).

⁴Note that a 2π pulse applies a π phase shift since the time evolution of the Rydberg Hamiltonian $e^{-iHt/\hbar} = e^{-i\Omega t \sigma_x/2}$, so when $t = 2\pi/\Omega$, then the time evolution is $e^{-i\pi \sigma_x}$.

Here we will discuss a zone-based implementation plan for cesium excitations in our experiment, as per the depiction in Fig. 4.11.

4.5.1 Cesium Rydberg transitions

As pictured in Fig. 2.5, and implemented across many experiments there are several approaches that can be taken for Rydberg excitation. One can choose to use single or two-photon excitation pathways, and for two-photon excitation pathways, there are additionally multiple options for intermediate states.

Single photon excitation

Single photon Rydberg excitation from the $6S_{1/2}$ ground state, as limited by selection rules, requires exciting to a P state. In cesium, the transition is around 319 nm, which is somewhat longer (less-UV) than single-photon transitions in other alkali atoms due to the heavier mass of Cs.

It is noteworthy that the oscillator strength of $6S_{1/2} \rightarrow nP_{1/2}$ transitions in Cs is around 4 orders of magnitude smaller than the oscillator strength of $6S_{1/2} \rightarrow nP_{3/2}$ transitions, due to a divergence in the core polarizability [179]. This is unique to Cs, and motivates us to only consider the $6S_{1/2} \rightarrow nP_{3/2}$ transition due to power requirements. In Fig. 4.17 wavelengths required to excite to various $nP_{3/2}$ states are recorded for reference, as calculated with quantum defect theory, using data collected in Table 1 of reference [180].

We decide to start our gate-based investigation using two-photon excitations to nS -states due to the DC electric field sensitivity of Rydberg nP -states which are enhanced due to near degeneracies with $(n - 1)D$ -states (arising from the quantum defects) [179, 181]. Such DC fields could perhaps be mitigated through RF dressing strategies [182, 183], or with a charging beam and specialized glass cell coatings [179], but we have chosen to proceed with the two-photon approach.

Two photon excitation

A two photon excitation can either take us to an nD -state or an nS -state, and since S states are less susceptible to DC electric fields, and since the interactions are isotropic, we examine this direction more closely. Wavelengths for the upper leg of two-photon excitations through the $7P_{1/2}$ and the $7P_{3/2}$ intermediate states are calculated with quantum defects and plotted in Fig. 4.19.

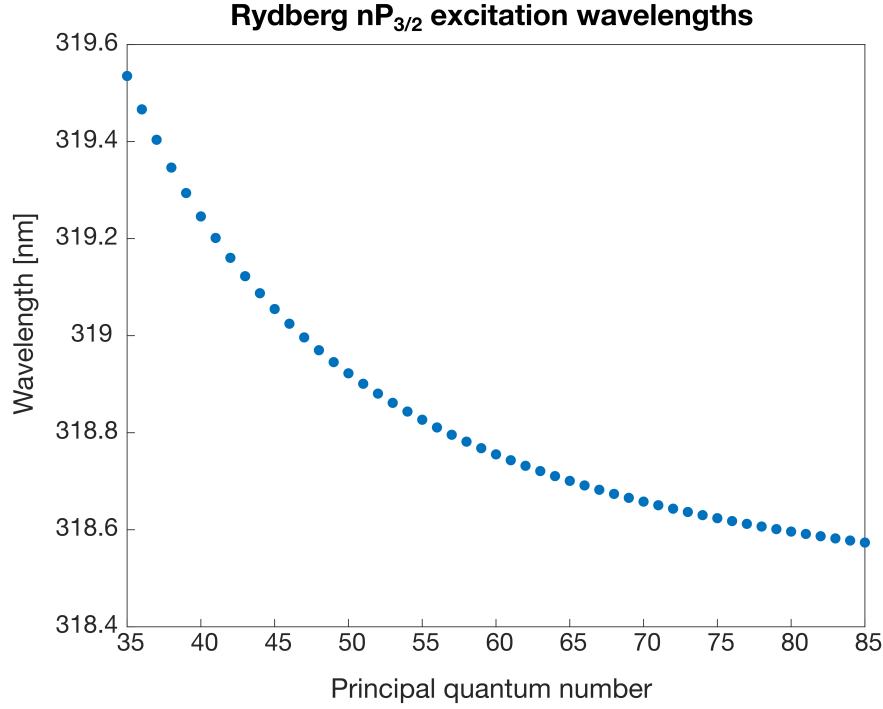


Figure 4.17: **Single-photon excitation wavelengths.** We depict a numerical calculation based on quantum defect theory for ground to Rydberg $nP_{3/2}$ transition wavelengths from $n = 35$ to $n = 85$.

The biggest downside to a two-photon excitation is introducing a fundamental limitation on fidelity through scattering off of the intermediate state as we demonstrate in the following by examining the equations of motion for the three-level system given by the ground, intermediate, and excited states.

Using a similar technique to that used by Steck for solving a Λ three-level system [62], we find the equations of motion $|\Psi\rangle = \Psi_g|g\rangle + \Psi_i|i\rangle + \Psi_r|r\rangle$ under the interaction Hamiltonian in the rotating frame given by

$$H = \frac{p^2}{2m} - \hbar\Delta|i\rangle\langle i| + \frac{\hbar\Omega_1}{2}(\sigma_1 e^{-i\vec{k}_1\cdot\vec{r}} + \sigma_1^\dagger e^{i\vec{k}_1\cdot\vec{r}}) + \frac{\hbar\Omega_2}{2}(\sigma_2 e^{-i\vec{k}_2\cdot\vec{r}} + \sigma_2^\dagger e^{i\vec{k}_2\cdot\vec{r}}) \quad (4.18)$$

We find the equations of motion

$$i\hbar\partial_t\Psi_i = \frac{p^2}{2m}\Psi_i + \frac{\hbar\Omega_1}{2}e^{i\vec{k}_1\cdot\vec{r}}\Psi_g + \frac{\hbar\Omega_2}{2}e^{i\vec{k}_2\cdot\vec{r}}\Psi_r - \hbar\Delta\Psi_i \quad (4.19)$$

$$i\hbar\partial_t\Psi_g = \frac{p^2}{2m}\Psi_g + \frac{\hbar\Omega_1}{2}e^{-i\vec{k}_1\cdot\vec{r}}\Psi_i \quad (4.20)$$

$$i\hbar\partial_t\Psi_r = \frac{p^2}{2m}\Psi_r + \frac{\hbar\Omega_2}{2}e^{-i\vec{k}_2\cdot\vec{r}}\Psi_i \quad (4.21)$$

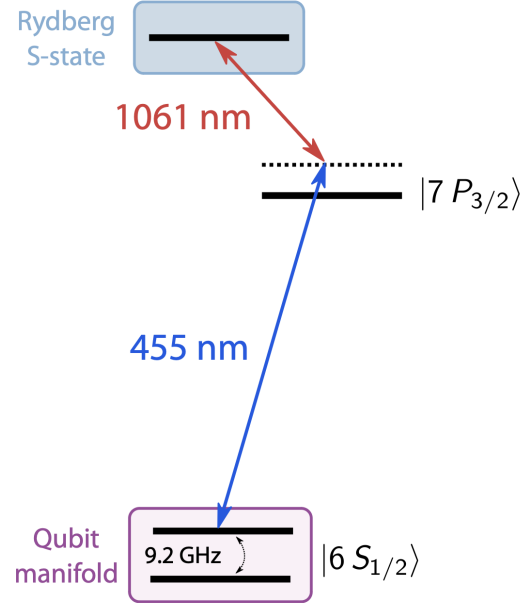


Figure 4.18: **Rydberg excitation pathway.** We picture the two-photon excitation pathway for Cs through the $7P_{3/2}$ state, consisting of a 455 nm and 1061 nm photon. We have available 120 W of power at 1061 nm and 2.5 W of power at 455 nm.

and using the approximation that $\partial_t \Psi_i = 0$ because it contains the quickly oscillating Δ term, we find

$$\Psi_i = \frac{\Omega_1}{2\Delta} e^{ik_1 \cdot r} \Psi_g + \frac{\Omega_2}{2\Delta} e^{ik_2 \cdot r} \Psi_r \quad (4.22)$$

which we substitute in to find an effective two-level system with Rabi frequency given by equation 2.47.

To obtain the scattering rate from the intermediate state with decay rate Γ , we recognize that

$$R_{sc} = \Gamma \rho_{ii} \quad (4.23)$$

which we can rewrite using equation 4.22 as

$$R_{sc} = \Gamma \frac{\Omega_1^2}{4\Delta^2} \rho_{gg} + \Gamma \frac{\Omega_2^2}{4\Delta^2} \rho_{rr} + \Gamma \frac{\Omega_R}{2\Delta} e^{i(k_2 - k_1) \cdot r} \rho_{gr} + \Gamma \frac{\Omega_R}{2\Delta} e^{i(k_1 - k_2) \cdot r} \rho_{rg} \quad (4.24)$$

Because $\rho_{gg} + \rho_{rr} \approx 1$ and taking $\Omega_1 = \Omega_2$, we can write

$$R_{sc} \approx \Gamma \frac{\Omega_2^2}{4\Delta^2} \quad (4.25)$$

For the scattering rate for the Rydberg states, we use tabulated values given in [184] for zero temperature. These values are calculated using the semi-empirical formula

$\tau_0 = \tau_s n_{\text{eff}}^\delta$, where τ_0 is the expected lifetime of the relevant Rydberg state, based on scaling with the effective principle quantum number n_{eff} and a known state lifetime for another state τ_s . The scaling coefficient $\delta \approx 3$ for high Rydberg states, as found empirically through experimental measurements.

We initially anticipated choosing the detuning, based on equalizing the scattering from the intermediate state and from the Rydberg state, which represent the two fundamental limits on Rydberg gate fidelity. In fixing a two-photon frequency to be around 5 MHz, and with anticipated laser powers, doing so would result in requiring a detuning of around 20 GHz.

However, recent full error-model simulations ⁵ have revealed that including the relevant states light shifts induced from the Rydberg lasers themselves (self-light shifts) have a large impact on resulting fidelities. While such energy offsets can be compensated in order to maintain the desired detunings during gate operations, they are proportional to intensity, and thus beam inhomogeneity and laser RIN, which result in differences of intensity across space and time, have a potentially strong impact which impacts the optimal choice of parameters. In particular, anticipated beam inhomogeneity that is possible due to beam shaping techniques has revealed that variation in light shifts across the array from the upper photon in the two-photon excitation will contribute strongly to Rydberg gate infidelity with larger detuning, if Rabi frequencies are maintained. The variable light shift induced by this beam both shifts the Rydberg and the ground state (due to D₁ and D₂ lines) strongly. This error source then suggests that closer detuning to the intermediate state would be optimal, while the error source from scattering off the intermediate state would suggest that increasing the detuning would be optimal. While the inhomogeneity of the Rydberg beam is not a fundamental limitation on the atomic side, engineering very flat potentials with currently existing techniques is a known challenge, with potentially increasing impact when scaling. To mitigate this effect, one could choose to optimize the detuning to balance error sources (2 GHz seems optimal after assuming about 1% flat-top inhomogeneity), such that the power needed in the upper photon of the two-photon transition is less (resulting also in less light shift) while maintaining a given Rabi frequency.

⁵See future thesis from the group by Gyohei Nomura for a more in-depth description.

4.5.2 Rydberg lasers

We utilize Titanium-Sapphire (Ti:Sa) lasers at 910 nm and 1060 nm for our Rydberg wavelengths. We double the 910 nm to 455 nm using a doubling crystal in a bow-tie cavity, and can obtain a maximum of around 2.5W output power. We then have two AOMs, one for intensity stabilization and a second which is double passed and used for frequency tuning and phase modulation. We use the 1060 nm output to seed a fiber amplifier that can output 130W of optical power. From error model simulations, we believe that the measured phase noise from the Ti:Sa lasers will be low enough to not be limiting in our error budget.

We additionally have a crossed cavity that we use to lock the 910 nm and 1060 nm lasers using a Pound Drever Hall (PDH) offset sideband locking scheme. The cavity is chosen to be crossed such that the Finesse in the two directions of the cavity can be chosen differently while maintaining some common mode fluctuations.

4.5.3 Rydberg zone

To execute two-qubit gates in parallel during quantum processor operations, we design an interaction zone as depicted in Fig. 4.11. We plan to utilize the two-photon excitation pathway pictured in the level scheme diagram in Fig. 4.18 through the $7P_{3/2}$ intermediate state. In order to achieve uniform two-photon Rabi frequency across the zone and to minimize the cross-talk between the Rydberg beam and qubits in the storage zone, we plan to employ spatial light modulators to shape the Rydberg beams into a flat-top profile [47, 185, 186], instead of Gaussian-shaped beams. We plan to use beam dimensions of approximately 100 μm in the atom-plane and 30 μm along the axis perpendicular to the plane. Our setup includes a laser system with an output power of 120 W at 1061 nm sourced from a fiber-amplified laser, and an output power of 2.5 W at 455 nm.

We take into account laser power losses along the optical path from components such as acousto-optic modulators (AOMs) and SLMs [186] assuming that $\sim 20 - 25\%$ of the output from the lasers reaches the atoms (different for the two paths taking into account wavelength-dependent diffraction efficiencies). For an intermediate state detuning of -2 GHz, we estimate achieving an effective two-photon Rabi frequency of approximately 6.9 MHz using the $60S_{1/2}$ Rydberg state.

For the implementation of the two-qubit gate using Rydberg blockade, we plan to position two interacting atoms 2.5 μm apart using two tweezers created by lasers of different wavelengths to avoid interference (see Fig. 2.19). Each pair will be spaced

by $11.4\text{ }\mu\text{m}$ from adjacent pairs. We note that at this large spacing of tweezers from the same laser source, out-of-plane interference is minimized. This configuration is anticipated to yield a blockade shift of around 300 MHz for the paired atoms and a residual interaction strength of ~ 0.03 MHz on neighboring sites, comparable to the residual interaction strengths in other contemporary experiments with state-of-the-art two-qubit gate fidelities [34, 187]. With these settings, we expect that state-of-the-art results for two-qubit gate fidelities can be reproduced at the scale of ~ 500 two-qubit gate sites, in the $\sim 110\text{ }\mu\text{m}$ by $\sim 900\text{ }\mu\text{m}$ zone.

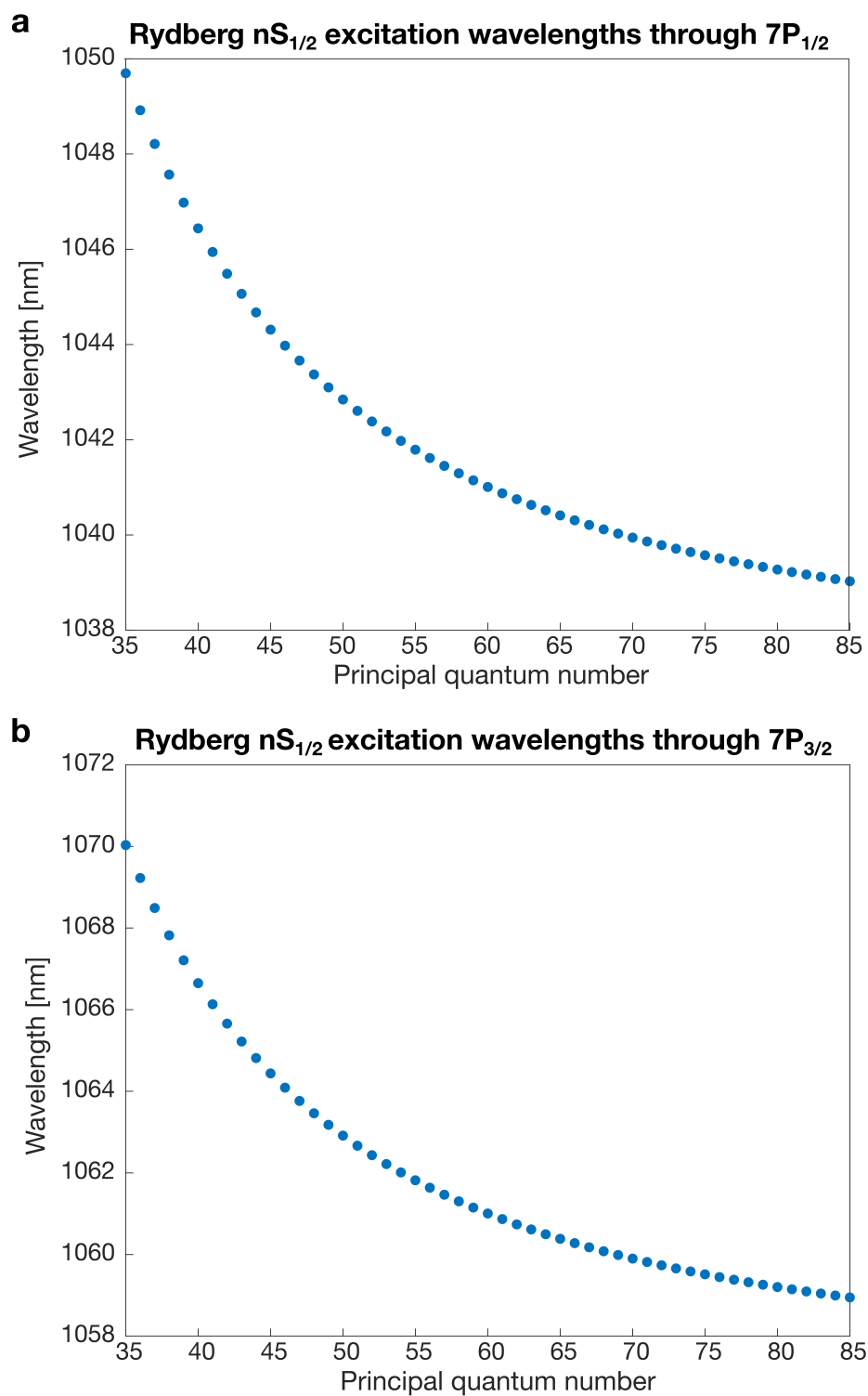


Figure 4.19: **Two-photon excitation wavelengths.** We depict a numerical calculation based on quantum defect theory for the wavelengths of the upper leg of a two-photon transition in Cs $6S_{1/2} \rightarrow$ intermediate state $\rightarrow nS_{1/2}$ state, from $n = 35$ to $n = 85$. **a**, Wavelength of upper transition through $7P_{1/2}$ state and **b**, wavelength of upper transition through $7P_{3/2}$ state.

Chapter 5

OUTLOOK AND SCALING TO UTILITY-SCALE

This work has not only demonstrated the largest (to this point) register of qubits alongside record imaging fidelities, coherence times, and long-distance moves, it also has allowed us to probe some of the limitations that currently exist when it comes to scaling neutral atom quantum systems.

The current scale of 6100 qubits demonstrated here, already presents a new frontier that will enable the first implementations of large scale experiments in the areas of quantum simulation, metrology, and computing.

However, for specific problems such as cryptographic applications or quantum chemistry applications, where resource estimates have been mapped out, scaling by a factor of at least two to twenty more will be necessary for utility-scale operation [14, 188]. Thus, in the following chapter, we discuss approaches and insights that we have for scaling further to ‘utility scale’ tweezer numbers for atom trapping (5.1)). We then discuss achieving connectivity at such larger scales (5.2)). We finally give an outlook perspective for near-term use of the platform presented in this dissertation for quantum science at scale (5.3)) and a conclusion (5.4)).

5.1 Trapping atoms at even larger scale

Our work indicates that further scaling of the optical tweezer array platform to tens of thousands of trapped atoms should be achievable with current technology, while essentially preserving high-fidelity control. In our present apparatus, several factors limit the number of sites. We explore these factors here, focusing on the most prominent limitations that we have run into: path thermalization and SLM pixel number.

As it stands, in our qubit register, we have pushed the limits for utilizing the pixel numbers of existing spatial light modulators, covering almost the whole field of view that the SLMs can access with traps in this work. We are currently combining two SLMs with polarization (Fig. 2.20), but only fully utilizing one of the two fully for tweezer creation, due to the thermal lensing that we incur on our objective. In the case where this thermal lensing can be mitigated, we could imagine interleaving more traps with different wavelengths or polarizations to increase the number of

atoms that can be trapped (see Fig. 2.19) ¹.

5.1.1 Thermalization

We first discuss the limitations arising due to thermalization, as this was the most directly limiting factor in the results displayed in this thesis, despite being something that can be readily addressed in our setup as well as easily adjusted in future experiments.

Objective lensing

In particular, we observe worsening optical aberrations at tweezer powers greater than that in the present study due to thermal heating of the objective. This is the main limitation on atom number for the results in this work, even after aberrations were mitigated using the SLM. Our understanding was that this effect was caused by the housing material for our objective expanding at higher power due to absorption. While our observations lead us to believe that there is not too much hysteresis at the power at which we now operate (i.e. with power cycling between high and low power, we have been able to obtain the same results consistently), seeing increasing aberrations led us to stop after making essentially as many tweezers as we could with a single SLM.

We have a few hypotheses for the cause of the heating. We first note that we did take precautions to minimize the heating and power incident on the objective housing by including an aluminum aperture that is the same diameter behind the objective aperture (not in thermal contact) in order to minimize the intensity clipped directly by the objective aperture.

From Zemax simulations that we implemented after these experiments, it is clear that if the alignment of the tweezer array to the back focal plane of the objective is inaccurate (by even a couple millimeters), then rays that create tweezers on the edge of the objective FOV would be clipped within the objective. It is certainly possible that our alignment is off by enough such that there is some clipping even though we are not utilizing the full objective FOV. Another possible contributing factor to the heating of our objective is the absorption of the lenses within the objective that have approximately a total transmission of 97.3% at 1060 nm only accounting for loss in

¹We note that spacing considerations may need to be revisited in the context of array operation due to observed atom loss during long-distance moves, but loading atoms into this configuration, and rearranging them to a distance that can be used for array operation may work well for experiments that are not continuously re-loaded

Material	Thermal Expansion ($10^{-6}/K$)	Thermal Conductivity (W/mK)	Hardness (GPa)	Young's Modulus (GPa)
Ultem	85	0.17	0.22	3
Macor	9	1.46	2.3	66.9
Shapal Hi M Soft	4.8	92	3.8	176
Aluminum	16-24	80-220	0.25-1.25	68-88.5
Titanium	8.5-9.3	16.3-18	1.15-1.25	100-105

Table 5.1: **Possible objective housing materials** Note that all values are quoted for room temperature. The data for the metals was found in [189] and [190], and data for the ceramics was found from [191].

the lens material, not due to interfaces. Such thermal lensing could be additionally contributing to the aberrations that we find with increasing power, and which we struggle to correct for with the SLM. Finally, the coatings that are used have a standard loss of 0.5%, with transmission that is dependent on angle of incidence. Extra power from back-reflections then can be reflected into the objective housing, and for an objective with around ten lenses, this is a significant amount of power that is then absorbed by the objective housing. Additionally, the coatings are less effective at higher incident angle, contributing to exacerbating this problem.

These contributing factors mostly contribute to incident power on the objective housing. This is a particular problem because our objective is made of Ultem, which has a high coefficient of thermal expansion, and a low coefficient of thermal conductivity. In table 5.1 we compare this material to other possible material choices for these and other relevant metrics. We initially chose this material because of its non-magnetic properties, because it has been used in other experiments including those in our lab, and because it is lighter, thus making for an objective that is easier to mount. However, for experiments, such as those which are continuously loaded where magnetic fields are not changing in the science chamber, and for the case where specialized mounts can be designed or obtained, these factors need not be weighted.

We find that Ultem in particular has a large coefficient of thermal expansion and a very low coefficient of thermal conductivity (it is an insulator). We believe that this combination, and the factors that contribute to power absorption by the housing lead to heating that causes expansion of the Ultem housing, thus shifting the placement of lenses relative to one-another, and leading to increasingly strong aberrations with

optical power. Due to the insulating properties of Ultem, the heating cannot be dissipated to any surrounding housing or added elements for active cooling.

To mitigate the issues that we have described, one could choose to design an objective with housing made of, for example *Shapal Hi M Soft*, which has low thermal expansion, and high thermal conductivity, thus allowing for less impact with high power, as well as active cooling if necessary². To address issues related to absorption and reflection from lens coatings, one could design an objective for a particular wavelength, and use entirely fused silica with AR or nanotextured coatings that have very high transmission and flexibility with angle of incidence.

We believe that addressing this issue would allow us to almost double the number of tweezers that we create using the two fiber amplifiers that we already have in our system. We further anticipate the potential to switch from polarization combination to wavelength-based array combination, opening up further avenues for increasing tweezer number with similar techniques to those utilized in this work by stacked dichroics for combining many SLM-fiber laser paths.

Thermalization and Heating of SLM

Another limiting factor is that at higher powers, we also observe the SLM diffraction efficiency decrease. We re-optimize the optimal SLM pixel-voltage/phase relationship (by updating the look-up-table referenced by the SLM) at high power using MLOOP [155], but we do not see significant improvement. It is our understanding that at high power, the approximately 2π modulation depth that we have may decrease due to the effects of temperature change on the liquid crystal properties. Thus, with decreased modulation depth, we cannot reach the same diffraction efficiency. Starting with SLMs that have higher modulation depth, such that there is leeway when heated is a good first step. There are also additional techniques that can be implemented with SLMs that have higher modulation depth in order to achieve more efficient power usage [90].

5.1.2 Limitations in tweezer generation

In order to utilize the full field of view of the objective, we have proposed stacking SLMs on different pathways with different lasers that are wavelength combined. While this can potentially be effective, if available laser power increases while the

²We note that we don't fully understand the impact of dielectrics near the glass cell on electric fields which are currently a difficult source of error in many Rydberg tweezer experiments.

number of tweezers that can be made with SLMs remains the same (i.e. because pixel numbers don't increase), this will not be optimal. Additionally, dichroic combination with arrays that have large incident angle requires separation by wavelengths that are $\sim 15 - 20$ nm apart.

In section 2.4, we showed the linear relationship between SLM pixel number and the number of resolvable spots in the array. Thus, to increase the FOV of the SLM and the subsequent number of tweezers that can be created, more pixels are required. The diffraction efficiency of tweezers across the field of view is also related to the pixel number, and thus, even for creation of the same array size, one would increase power efficiency by having a phase pattern created with more pixels. The pixels don't need to be smaller (in fact, for SLMs, pixel cross-talk increases with smaller sizes), but we need more of them.

To solve this problem, there have been proposals and demonstrations of large numbers of optical tweezers [89] for trapping more atoms, in particular using metasurfaces, where more pixels can be readily created. We explore a few considerations in this direction.

Metasurface arrays

Metasurfaces are 2D surfaces, fabricated with lithography and nanoimprinting methods, allowing for wavefront shaping through engineering phase discontinuity on an interface to steer light [192]. Metasurfaces can be utilized for holography and for creating lenses among other applications and have been engineered for both situations requiring polarization insensitivity, or for situations requiring different holograms for opposite polarizations [193, 194].

Metasurfaces have been used to demonstrate the creation of tweezer arrays with hundreds of thousands of tweezers [89]. These tweezer arrays can either be created in a configuration where they are focused directly using the metasurface (as in the demonstration in [89]), or in a holographic configuration where they apply a phase pattern in a Fourier plane that is then focused by an objective.

Both options pose challenges. In the case of utilizing the metasurface in a Fourier plane (holographic configuration), the requirements on optics for relaying tweezers with strong incident angles become extreme for increasingly large tweezer arrays with a large tweezer FOV in the atom plane. In particular, we discuss the cone half-angle (CHA) and optical element size requirements for optics for particular

tweezer FOV coverage in the next section.

In the case of direct focusing, where the metasurface is in a lens plane, one can use relay optics to map the array to the science chamber, in which case, the same problems arise as for the Fourier plane configuration regarding requirements on relay optics. Another option would be to use the metasurface to focus directly into vacuum, either with a metasurface directly outside or directly inside the science chamber. The option of direct focusing without relay into a vacuum requires creating a metasurface with a long focal length, such that atoms are not too close to the glass cell. Such an approach could also have an impact on the electric field that the atoms are exposed to, which has a strong impact on the stability of the Rydberg excitation wavelength, particularly at high principal quantum number.

The approach of directly focusing does have the benefit of separating background light, which remains unfocused, from the rest of the array. In the holography configuration, this light would be focused with the array, similar to a 0th order in an SLM phase-modulation setup. Thus, for the directly focusing case, approaches such as filtering would not be required. Using a metasurface in such a configuration would still require an objective for atom imaging and for any other wavelength that would be needed opposite the metasurface, since designing the metasurface to operate as an objective for other wavelengths simultaneously is not something which has been demonstrated.

Relay optics requirements

For the case where optics are used to relay a tweezer pattern to an objective that focuses the array into a science chamber, with increasing objective FOV that is covered, there arise more significant requirements for relay optics when it comes to both their size and CHA. There additionally are tradeoffs that can be made between the two— if the array is magnified, then the CHA incident on optics in the pathway will be less, but the optics must be larger to compensate. We explore this tradeoff in the following and discuss array sizes that would cover up to a 6mm objective FOV, equating to over 500,000 trapped atoms with our current array spacings.

We depict in Fig. 5.1a a relay lens configuration between an optical element that imprints a holographic phase onto the light, which is relayed with some magnification M onto an objective³. We depict in the rest of Fig. 5.1 the ensuing requirements on

³Typically, the back Fourier plane of the objective (or back focal node) which will be the image of the metasurface, as labeled here, will be a few centimeters inside of a large custom objective.

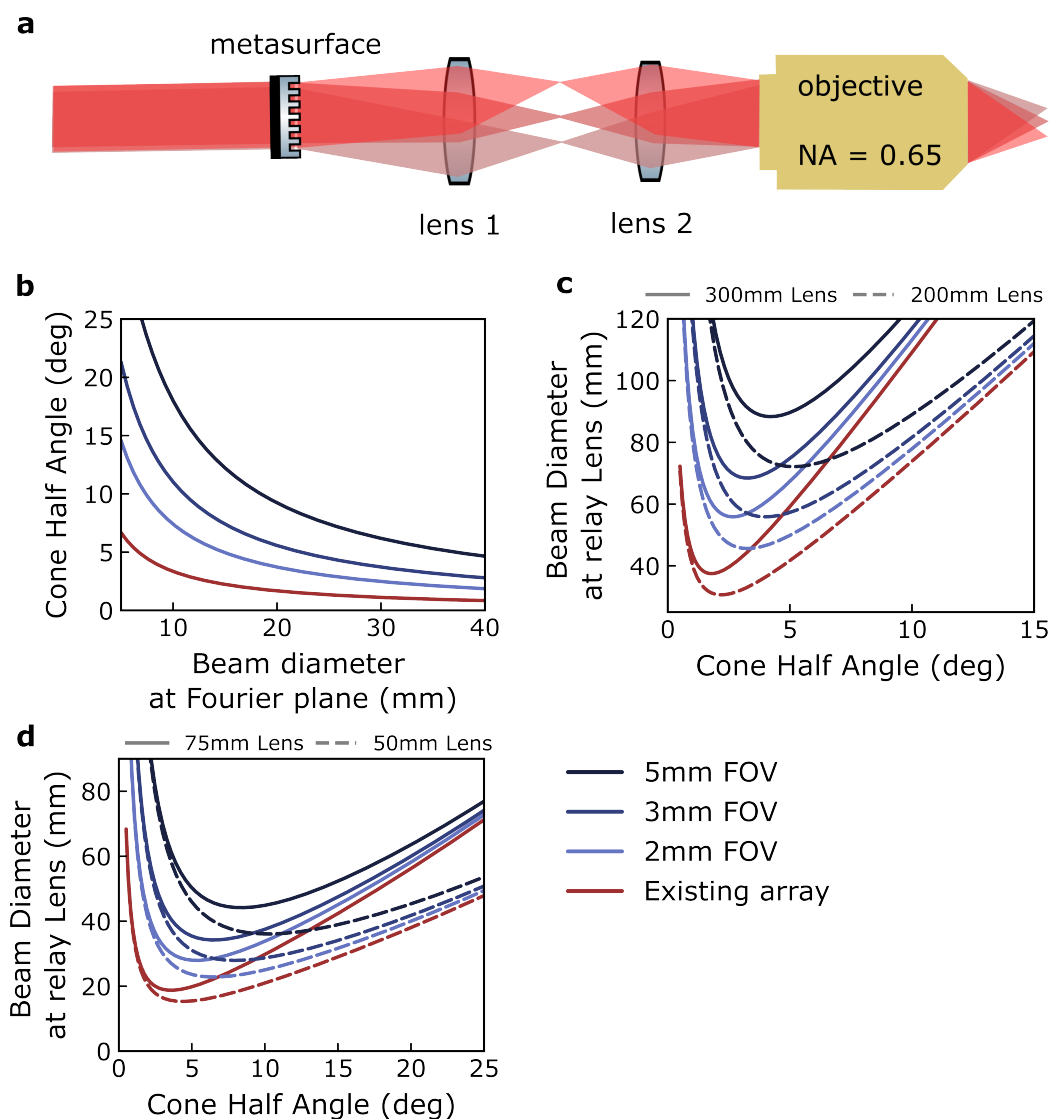


Figure 5.1: Relay Lens Requirements. For the case of using a diffractive element such as a metasurface to imprint a holographic phase, we show **a**, a schematic for a relay configuration to a high NA objective, and the requirements on these relay optics when it comes to **b**, Cone half-angle (CHA) as a function of beam diameter at the Fourier plane, and **c**, Lens size requirements, as a function of the required CHA, determined by beam size. These requirements could be converted to being independent of NA (through utilizing resolvable spots, although this depends on where the beam is clipped).

the relay optics that must be imposed given a particular FOV choice, assuming an objective with $NA = 0.65$, and wavelength of 1060 nm.

First in Fig. 5.1b, we examine the CHA requirements of the lenses as a function of the beam diameter in the Fourier plane of the lenses, for various FOV covered in

the objective focal plane. For the case of a 4f telescope the requirements will be the same for the two lenses, but for the case of a magnifying telescope, they will be different. The beam size in the back focal plane of the objective determines the necessary CHA for lens 2, and the beam size in the metasurface plane (or the plane of another diffractive element creating a holographic phase) determines the CHA requirement for lens 1. The existing array which is plotted is the array with 6100 atoms and 900 μm diameter.

The CHA and beam diameter trade-off in the Fourier plane then have an impact on the beam size at the lens itself, which is also dependent on the focal length of the lens. We demonstrate this tradeoff in Fig. 5.1c, for different FOV arrays (as determined in the objective's focal plane). As can be seen, for each FOV in the objective focal plane, and for each focal length of the relay lenses, the beam diameter at the given relay lens is a function of the incoming CHA. These taken together provide constraints for designing custom lenses in order to relay large arrays, which is required if not using a metasurface to directly focus into a glass cell, which would potentially cause difficult-to-control electric fields due to the distance of the atoms from the glass and from dielectric surfaces.

To make concrete a particular scenario, if a 1cm diameter metasurface is used to generate traps for an array that spans a 5 mm FOV in the focal plane of the objective, the CHA from the metasurface, incident on 'lens 1' would be about 20° , and for a magnification of $4\times$ to the objective back focal plane with 40 mm diameter using relay lenses with focal lengths of 75[50]mm and 300[200]mm, the lenses would need to account for beam sizes of around 60[45]mm and 90[75]mm respectively (See Fig. 5.1c,d). The CHA that the 'lens 2' would need to accommodate in this scenario would be about 5° , per Fig. 5.1b.

5.1.3 Increased Loading

It is worth noting in terms of scaling up atoms in optical tweezer arrays, that we have created about 12,000 optical tweezers and we are only using about half of the tweezers to trap atoms in each shot due to the stochastic nature of atom loading. With the current experiment, there are a few options with which we could use the tweezers we have created more effectively. In particular, we could pursue the option of enhanced loading, or, following recent demonstrations, we could upgrade the experiment to be continuously re-loaded.

Enhanced Loading

The stochastic nature of the tweezer-loading process rests on the short-lived bi-atomic molecular bound states that atoms are being excited to during the light-assisted collision (LAC) stage of atom loading, resulting in pair-wise loss. In our experiments, LAC occur during the 3D Molasses stage since the cooling configuration already performs this action, and results in stochastic loading. However, there have been many systematic studies that demonstrate significantly more control over the excitations in order to expel atoms one at a time, rather than in pairs and achieve larger array filling fractions. This has resulted in single tweezers loading above 90% and arrays that are loaded up to 80% [195] in alkali atoms and above 90% array loading in alkaline earth atom systems [196]. In alkali atoms, the approach to enable enhanced loading has been through Λ -grey molasses [195, 197, 198], with optimized detunings relative to the tweezer trap depths. While observations of slightly above half filling for Cs have been mentioned in some places [199], careful studies of enhanced loading in Cs have not been published, to our understanding.

In order to implement this in our experiment with high levels of control over the detuning in order to optimize the loading carefully, we have prepared the following locking scheme for the D_1 line. We maintain the offset-lock, as described in Chapter 2.3.2 for the optical pumping beam, and phase lock a second DBR laser to this one using a heterodyne beat detector referenced to a DDS, as in the case of the D_2 line. This second laser we use to prepare both the cooling and repump for Λ -grey molasses, by using an electro-optic modulator (EOM) to create a sideband 9.2 GHz from the carrier. This setup allows for careful optimization of Λ -grey molasses and enhanced loading, as the relative intensities of the carrier and sidebands can be tuned via the EOM RF power, and the Λ -grey molasses detuning can be changed with a large accessible range by changing the DDS and the relative phase lock with the laser that is offset-locked to the vapor cell.⁴

Continuous Re-loading

There have been several recent demonstrations of continuously loaded arrays of neutral atoms in optical tweezers [56–58]. Such a configuration allows for atoms to be re-filled when they are lost, but also for all tweezers used for trapping to be filled.

⁴Note that this setup was implemented as part of Shrishti Kulkarni's undergraduate summer research project, mentored by Kon Leung and me.

While our experiment was not originally designed for continuous reloading, we do have ideas for upgrading the experiment such that it can operate in a continual fashion. In particular, the 2D MOT could potentially be repurposed into a 3D MOT that is directly loaded from background vapor by removing the PICAS optics and building 3D optics around the glass cell. A lattice could then be aligned with the 750 μm differential pumping hole in this cell, and a hand-off could be performed in the 3D MOT cell, with hiding beams co-aligned with parts of the lattice as necessary. Such an upgrade could enable the system to operate continuously at moderate flux, and to utilize all the tweezers being created, effectively doubling the number of qubits. In the most simple case, this could be done without ever breaking vacuum. Additionally, since the FOV of the objective extends beyond the 12,000 tweezer array, there would be space to implement other necessary operations for refilling traps in a quantum processor.

Dark-state re-loading techniques [166, 200] also present a promising approach to achieving higher loading probabilities in the array without a full continuous-reloading upgrade. Here, hiding beams may be used to enable the application of such techniques in a single-species array of alkali-atoms.

5.2 Controlling atoms

Beyond simply trapping more atoms, maintaining control of the atoms for operating a quantum processor at scale is a significant challenge. While AODs can operate in parallel in the ways described in Chapter 4 for the array size described in this work, when scaling array sizes several times further, the amount of multiplexing that is necessary may pose significant difficulties. Additionally, the overhead of the number of AODs and pathways necessary when considering how to compensate for cylindrical aberrations that arise in AODs is significant.

Particularly when it comes to operations for quantum computing and syndrome extraction circuits requiring long-distance connectivity, it becomes clear that desired timescales of around and under 1 ms will require technological innovation. Ideas related to multiplexing devices in order to overcome these limitations present promising paths forward that deserve careful investigation.

5.3 Near term outlook

Beyond increasing the system size of our array further, there is much that can be done with the existing system in the directions of quantum simulation and computation—both with and without error correction. In particular, with the addition of Rydberg

gates, and the capacity for local single-qubit Raman-based addressing (whether zone-based [51] or through single-site addressing [201]), our 6, 100 qubits will be usable in a quantum processor at scale, heralding a new era of quantum science.

This new generation of neutral atom quantum processors based on several thousand qubits is particularly relevant for QEC [12, 14, 202]. Near-term outlooks in this direction involve memory experiments for quantum error correcting codes with high encoding rates, as well as experiments that probe the exponential decay of logical error with code distance for the Rydberg-gate based neutral atom platform (similar to experiments implemented by Google on a superconducting qubit platform [20]).

Additionally, large-scale programmable devices enabling advances in quantum metrology [5, 38–40, 54, 203] and simulation [6–8] are made accessible through this work. For example, our platform—with the demonstrated qubit numbers—could be used for verifiable quantum advantage with low-depth evolution [8, 9]. Tweezer clocks could be scaled using near-infrared, high-power tweezers for loading and imaging [204] before transferring atoms to magic-wavelength traps for clock operation [38–40, 54]. We also foresee applications in quantum simulation for problems where boundary effects play an important role [6, 41, 45–47], which can be minimized with the large system sizes demonstrated here.

5.4 Conclusion

In this thesis, we demonstrated the design of and results achieved utilizing a new platform for trapping and control of an array of 6, 100 atomic qubits in far-detuned optical tweezers. The platform was specifically designed for scalability of tweezer number while maintaining coherence times, as well as the ways the experimental apparatus was optimized for long vacuum lifetimes. A hybrid-momentum Weighted-Gershberg Saxton algorithm led also to improvements in large-scale array generation.

We further demonstrated high imaging fidelity and imaging survival using many-image characterization. Uniformity was demonstrated in the spatial distribution across the array for parameters such as loading, scattering rates, and trap depth across the array. State-of-the-art long vacuum lifetimes which minimize un-avoidable atom loss were measured.

We went on to demonstrate record-long T_1 and T_2^* coherence times for an alkali atom array. Dynamical decoupling methods were further applied to extend the T_2^* time to a T_2 time of over 12 s. These coherence times contributed to high single-qubit gate

fidelities which were characterized with randomized benchmarking.

The control over the connectivity of atoms in this array was further examined, including limitations arising from long-distance atomic movement such as cylindrical lensing and ways to mitigate these impacts. We utilized interleaved randomized benchmarking to investigate the fidelity of long-distance moves, of atomic trap-transfer between AOD and SLM traps, and of the combination of trap-transfer and long-distance moves. Improvements to fidelities or equivalently to move-times were achieved by utilizing machine-learning techniques. We further discussed details regarding Rydberg excitations for two-qubit gates, which is current work in-progress.

Finally, we have discussed necessary considerations for scaling neutral atom arrays even further by an order of magnitude, including considerations regarding thermalization, regarding relay lens requirements, and regarding achieving higher loading probabilities in the existing array.

The results in this thesis therefore have presented a novel platform with 6, 100 highly-coherent atomic-qubits, and a direct pathway towards computation and quantum error correction utilizing this number of qubits. This thesis further examined the limitations regarding scalability to even larger system sizes and presented approaches to circumventing these limitations, anticipating near-term scaling by at least twice the array size while maintaining controllability and fidelity.

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Appendix A

VACUUM CHAMBER AND BAKING

A.1 Chamber design

The glass cell design was optimized to allow for 2cm MOT beams passing through adjacent to our microscope objective, and has a thickness of 6.35 mm in order to minimize bending and resulting aberrations due to vacuum deformations. Based on Molflow simulations, and calculations for maximizing the conductance for pumping particles in the glass cell, we minimized the distance between the science cell and the main chamber (note that due to manufacturing issues we ended up adding a conical reducer that had not been a part of the the original design).

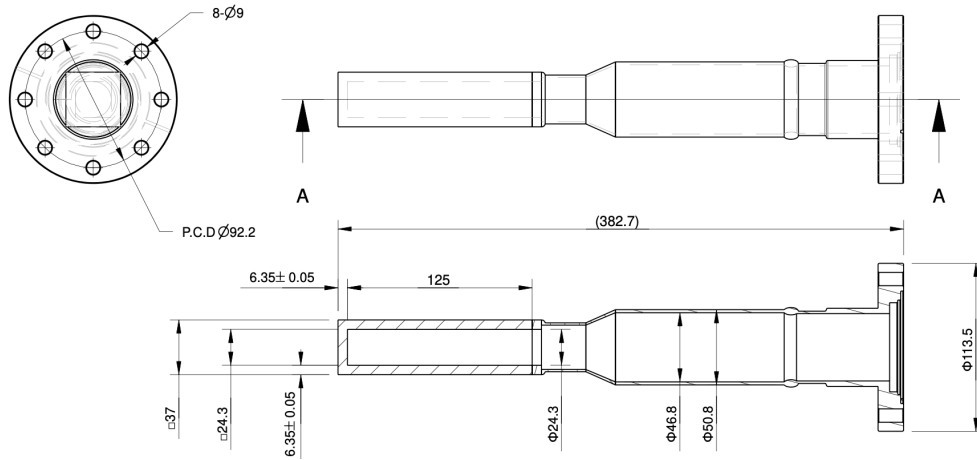


Figure A.1: Pictured is the glass cell design for the experiment that was manufactured by Japan Cell. Our cell is made of uncoated quartz, for the sake of wavelength flexibility. It was designed to maximize NA, enable large MOT beams around a large-diameter objective, and minimize cell deformation under vacuum. The labels in the figure are in units of millimeters.

A.2 MolFlow Simulations

We simulate our expected vacuum pressure utilizing MolFlow, a Monte Carlo based vacuum simulator in which one can input a simplified chamber design and expected outgassing rates and pump rates, and calculate expected partial pressures of various gases.

We implement this simulation, taking into account the following sources of back-

ground pressure: H_2O from the glass cell outgassing, Cs from the 2D MOT Inflection cell which is line-of-sight to the tweezer location, and He permeation through the glass cell from external atmospheric partial pressure. We find that the contribution from the line-of-sight to the 2D MOT is likely dominant in our setup.

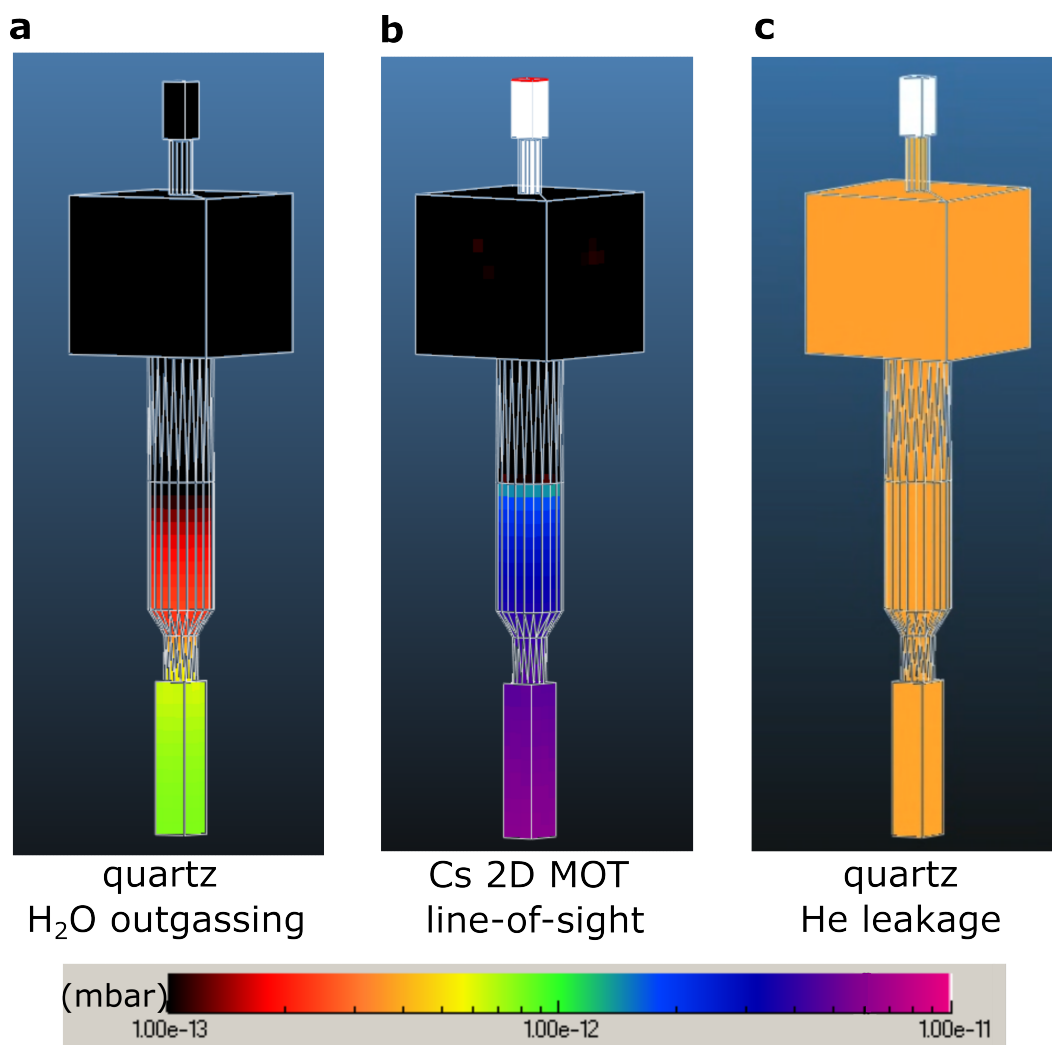


Figure A.2: Vacuum-limited Lifetime MolFlow Simulations We depict simulations for the equilibrium partial pressure from three different sources. We use a simplified schematic for estimation purposes that still maintains the correct distances and sizes of our apparatus. The square box is our steel chamber that has two attached TiSub pumps and an ion pump, the top cell is our 2D MOT, and the bottom is our glass science cell where we trap atoms. Contributions to background pressure are depicted from **a**, outgassing of H_2O from the quartz glass cell, **b**, line-of-sight to Cs 2D MOT source cell, **c**, leakage of He through our 6.35mm thick quartz glass cell. The pressure throughout the chamber appears more uniform for He leakage due to the lack of pumping via the sputtered titanium.

Element	Leak Rate (mbar*L/s*cm ²)	Pumping: Ti-Sub sticking coeff Ion pump rate (L/s)	Science Cell Partial Pressure (mbar)
H ₂ O	1×10^{-13} [205]	0.5 [206]	9×10^{-13}
Cs	*	0.95	6×10^{-12}
He	3×10^{-15} [207]	7.5 (L/s) [208]	4×10^{-13}

Table A.1: **Expected Contributions to Vacuum Background** Leak rates and pumping rates that were used for MolFlow Monte Carlo simulations for various elements that we expect to contribute to the background vacuum pressure. We consider H₂O outgassing from the glass cell, Cs background due to line-of-sight to the 2D MOT and He permeation through the glass cell. We note that the H₂O outgassing rate that we expect is between values given in the cited literature due to the amount of time and temperature at which we baked [205]. We note that in the simulation, the sticking rates are applied to all areas that are coated by the TSPs, while the ion pump rate is applied to only one side of the chamber.

* We choose an outgassing rate of Cs in the source cell, such that the equilibrium pressure in the cell is 5×10^{-7} mbar, which is central to the range of 2D MOT pressure quoted by Infleqion.

A.2.1 Water Outgassing

The source we have for outgassing rates from glass gives the outgassing rate of glass that is cleaned, and the outgassing rate for glass that has been baked at 400°C for 24 hours. These cases however do not represent our scenario given that we baked at 225 – 250°C for significantly longer (16 days).

For a material with thickness d that has been baked (or ‘de-gassed’) for time t at temperature T , the outgassing rate will be [206]

$$\dot{Q} = \frac{4c_0 D_r}{d} e^{-\frac{\pi^2 D_T t}{d}} \quad (\text{A.1})$$

where c_0 is the initial concentration of the gas in the material, and D_r and D_T represent the temperature-dependent diffusion coefficient at room temperature and the baking temperature respectively. $D = D_0 e^{-E_d/RT}$ where E_d is the activation energy for diffusion, and D_0 is the diffusivity at infinite temperature. Because the thermal energy of our bake is significantly lower than the activation energy barrier for water in glass ($RT \ll E_d$), we are operating in a regime where the diffusion coefficient increases steeply with temperature. Consequently, while the outgassing rate decays exponentially with baking time, the rate of that decay is exquisitely sensitive to the baking temperature due to this Arrhenius dependence.

Thus for our simulation, we choose to use a value for outgassing that is about an order

of magnitude higher than that quoted for the outgassing over the higher temperature bake. We find a sticking coefficient of water with titanium of 0.5 [206] and thus use this across all surfaces that are coated in the simulation.

A.2.2 Cesium 2D MOT line-of-sight

We estimate the partial pressure in the science cell due to the line-of-sight with the 2D MOT cell. We choose the outgassing rate of Cs in the 2D MOT source cell to reach an equilibrium pressure in the cell of 5×10^{-7} mbar. This cell then has a differential pumping aperture to the main chamber of 0.75 mm. We only account for the impact of the hot room temperature atoms and don't take into account the slight increase in pressure that would be caused by the cooled atoms in the push beam during the 3D MOT loading.

Cs, being an alkali atom, is highly reactive with titanium, and the work function of titanium is ~ 4.33 eV and the ionization energy of Cs is ~ 3.89 eV, so when the titanium surface is fresh, the sticking coefficient of Cs is very close to 1 (we use 0.95 in the simulation). With very high surface coverage of Cs, this can decrease, but we expect it to remain high still due to surface charge migration. We also can choose to refire the titanium sublimation pumps to refresh the pumping rate, which we have not yet found necessary.

A.2.3 Helium permeation through Glass

We take into account helium permeation through the glass cell, because although helium accounts for a very small percentage of the atmosphere, it is small enough to penetrate through glass (unlike H_2 for instance), and it is not well pumped by the TSPs given that it is an inert gas. We estimate the leak rate with the following equation,

$$q = \frac{K \Delta P}{d} \quad (\text{A.2})$$

where $K = D_r \times S$ is the permeability constant, equal to the product of the diffusion coefficient at room temperature D_r and the solubility S [207]. $\Delta P = 1$ atm is the pressure difference between the outside and inside of the glass cell, and d mm is the thickness of the glass cell. We include both outgassing from the science cell where $d = 6.35$ mm as well as from the 2D MOT source cell where $d \approx 1$ mm.

We note that helium is more likely to permeate the glass cell due to its molecular size and shape when compared with other molecules, including H_2 . Given that helium is a noble gas, the titanium deposited on the chamber walls does not pump it

well, but the ion pump can contribute to a pumping rate. We reference [208] as well as information provided by the manufacturer to estimate a pumping rate for He with our 75S Gamma Vacuum (CX-TiTan CVX, XHV/SEM) ion pump. In particular, we estimate that the pump rate of helium is $\sim 10\%$ that of N_2 .

A.2.4 Titanium Sublimation pumps

Titanium-sublimation pumps consist of filaments that outgass titanium when current is run through them, which coats any surface that is line-of-sight. The deposited layers of titanium then provide a pumping surface that has very high sticking coefficients for many atomic and molecular species [206] (though not for noble gases). The surface coverage of this method in comparison to other pumps provides significant pumping rates.

We have a TSP with 3 filaments in the top nipple, and a shorter folded single filament TSP in the bottom nipple, to fit within the spatial constraints but still provide titanium coverage for the top half of the glass-to-metal transition

We optimize the length of the 6in diameter nipples attached to the top and bottom of our chamber in order to optimize line-of-sight sputtering from the titanium sublimation pumps to coat all parts of the chamber except for the science cell (but including the glass-to-metal transition). We picture the initial design in Fig. A.3, though we note that an additional conical reducer between the glass cell and the chamber extends the distance from the atoms to the source and extends slightly the amount of titanium-free glass in the glass cell arm. We have a 45° elbow to protect the ion pump from direct line of sight to the TSPs, but to still provide surface coverage for part of the elbow itself.

A.3 Baking Timeline

We include the timeline for our baking procedure in Fig. A.4. We implemented an initial bake at 400°C of all components that could sustain this temperature, and then a final bake around $225 - 250^\circ\text{C}$ for the whole chamber. We used sonication in 99.99% purity acetone and methanol¹ to clean vacuum parts that were suspicious for various reasons (visible contamination or non vacuum-safe packaging). Several bakes were performed due to various unexplained leaks and the necessity to replace various vacuum parts as depicted in the baking timeline in Fig. A.4. In particular, due to various circumstances described in the timeline, two bakes at 400°C were

¹Note that safety precautions should be taken when using chemicals, and particularly methanol due to toxicity.

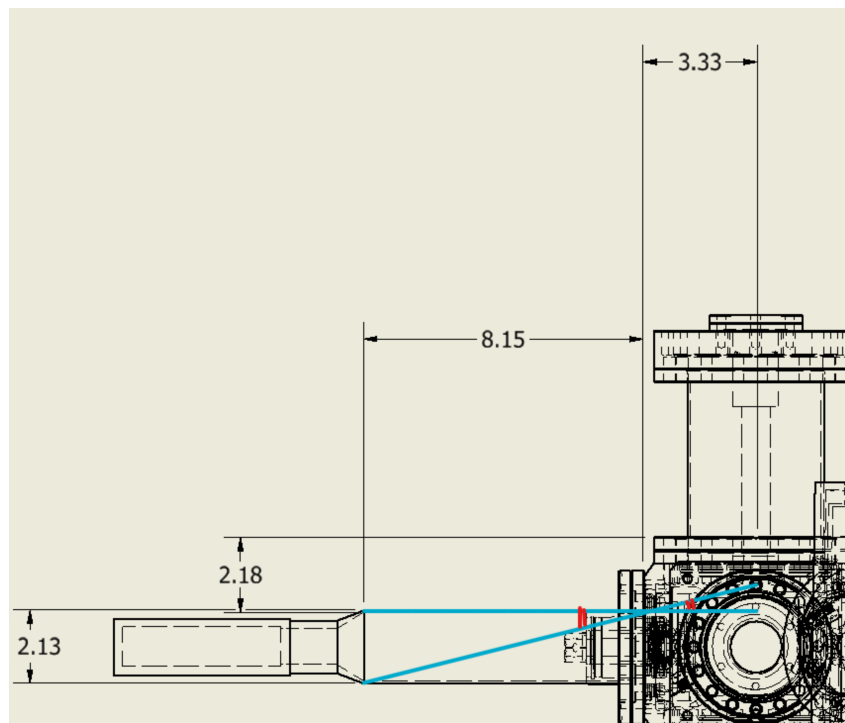


Figure A.3: **Titanium sublimation pump line-of-sight** A sketch of the designed line-of-sight sputtering from the TSPs to the glass cell. The TSPs are mounted inside of custom nipples, attached to the 6" steel chamber top and bottom (the top nipple and filament are pictured here, but the same procedure is implemented for the bottom TSP). The length for the custom nipples is chosen in order to optimize line-of-sight sputtering to coat the glass-to-metal transition in the arm of the glass cell. The diagram denotes the choice for the minimum scattering angle for the chosen line-of-sight sputtering. We note that an additional conical reducer was added later between the glass cell and chamber due to manufacturing issues, which slightly increases the distance between the sputtering line-of-sight and the science cell in our experiment. Labeled distances are in units of inches.

performed for a total of around 25 days, and five bakes at $\sim 250^{\circ}\text{C}$ were performed for a total of around 86 days. The number of bakes that we performed may contribute to the long vacuum lifetimes that we measure.

Our bakes were performed primarily under vacuum, following the procedure outlined in this thesis [33]. We pumped down with a roughing/turbo pump and then applied a layer of UHV-compatible aluminum foil, followed by fiberglass insulated heat tape. We placed thermocouples along the assembly to monitor the uniformity of the temperature during the bake, and to ensure that various parts did not exceed the maximum temperature rating. This layer was then covered in UHV foil again, followed by another layer of fiberglass insulation, and then a final layer of UHV

foil. When ramping up the temperature, we increased the temperature by around 2°C/min. At the end of the final bake, we degassed each filament in both titanium sublimation pumps and hi-potted the ion pump. We then fired all filaments at 50A for 120s to coat our vacuum chamber after the temperature was brought down, after which we sealed off the turbo pump.

We note that several leaks that were limiting to our final vacuum pressures prior to our troubleshooting were not apparent on a residual gas analyzer (RGA), and only could be found when utilizing a helium leak-checker. When using the helium leak-checker, one should be aware of the expected time frame between spraying helium and seeing an increase in value to signal a leak. In particular, we took note of the time between the helium being sprayed and a rise in the reading to ensure differentiation between true leak sources (visible as an increase in reading within tens of seconds), and helium traveling through the air outside the vacuum to the leak checker (closer to around 1 – 1.5mins before appearing).

We found it necessary to replace all silver-coated gaskets that were used in the high temperature pre-bake before the final lower temperature bakes even for parts that remained attached (we replaced them with standard copper gaskets). We saw leaks particularly at locations where gaskets had not been replaced.

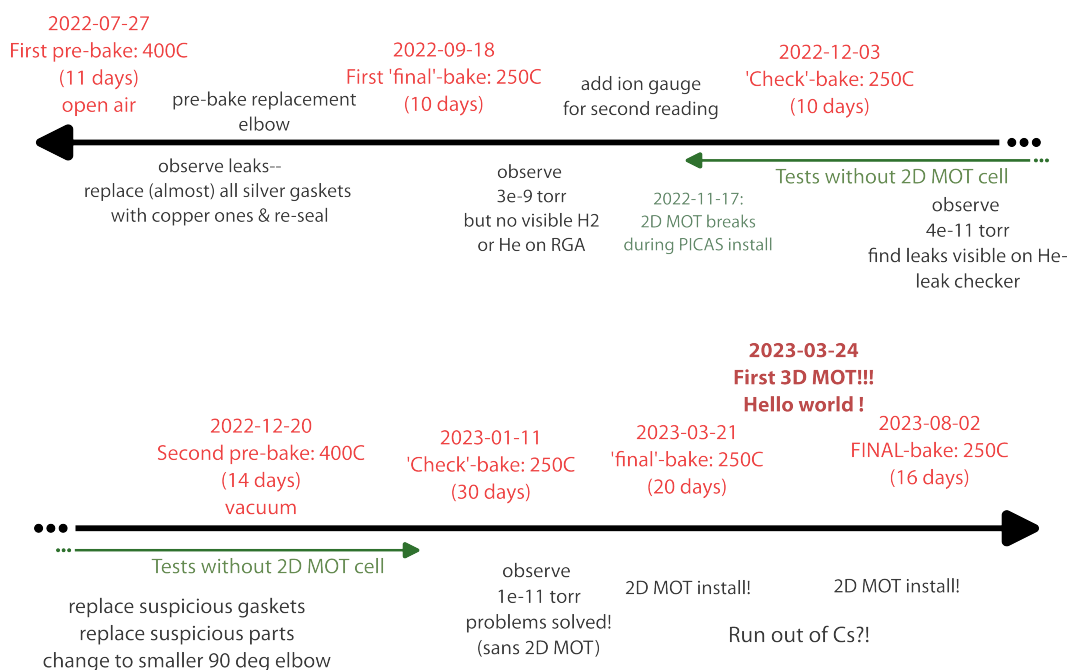


Figure A.4: **Baking timeline.** We include a timeline of the baking procedure that was implemented in the experiment. We note that much of the process was not planned and several bakes were implemented in order to troubleshoot unexplained leaks. We note that some of our long vacuum lifetimes may be attributable to this bake process, despite the unplanned nature.

Appendix B

MAGNETIC COIL DESIGN

We design our magnetic field coils to allow for high field in any direction while maintaining high optical access as well as spatial uniformity of the field across the objective FOV. As such, we utilize water-cooled coils that can be operated at high currents in three dimensions (we label the axes as ‘MOT’, ‘Objective’, and ‘Cell’), in addition to a low-field ‘offset’ coil in the MOT direction, which is used to apply a Helmholtz field to cancel external fields when the water-cooled coil is used for generating the Anti-Helmholtz MOT field.

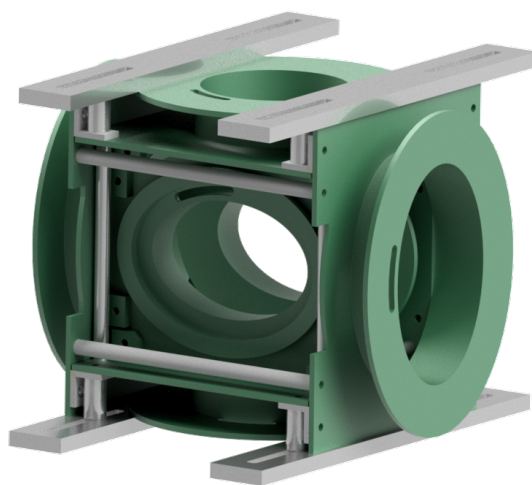


Figure B.1: **Assembled Coil Mount Design** Pictured is the assembled mount structure on which the coils were wound. The mount material is Garolite-10, and the mount configuration can be adjusted relative to the science cell in three axes, via custom aluminum attachments anchored to both horizontal breadboard layers.

We initially were motivated to design for high fields in all directions for a few potential applications. In particular, given that cesium has many hyperfine magnetic

Coil	Number of turns	Inner Major Diameter (mm)	Inner Minor Diameter (mm)	Closest separation (mm)
Low Field	12 (4×3)	187.5	136.7	122.0
MOT Coil	15 (5×3)	162.1	106.2	73.5
Objective Coil	40 (8×5)	130.8	130.8	192.4
Cell Coil	36 (6×6)	191.0	130.8	243.2

Table B.1: **Coil Geometry.** For each magnetic field coil, we include the number of turns and the parameters for the major and minor diameters of the elliptical shape of the inner-most feature. We additionally include the minimal separation between the two coils in the pair. We note that the copper wire used for the low-field offset-field coil pair are not water cooled, but all of the other axes are water cooled and wound with coils whose cross-section is depicted in Fig. B.3.

sublevels (due to the $I = 7/2$ nuclear spin), separating these states to allow for energy-selective operations such as state preparation [177] seemed potentially useful. Cesium additionally has several Feshbach resonances at accessible fields [209], and having the capacity to access these for applications regarding quantum simulation is another benefit with access to high fields.

We have come to realize after working with the coils, that mitigating the noise imparted due to using high current sources to achieve necessary fields is paramount. Due to the water-cooled nature of the coils, we operate in a floating configuration, ensuring also that the analog modulation for the current is passed through an isolation amplifier. In particular, one must pay attention to the stability of the current supply. We recently have been using air-cooled current 3A sources from Wavelength Electronics (with current noise of $\sim 35\mu\text{A}$ rms in constant current operation) for the offset coils, the objective coil axis, and the cell coil axis, where we don't require high current for our recent operations.

Along the quantization axis, our original configuration used an identical H-Bridge switching mechanism as presented in this thesis [33]. However, this has been updated to a new design utilizing MOSFETs in place of IGBTs to mitigate noise and allow for switching between different current sources for the Anti-Helmholtz and Helmholtz configurations¹. For the high current required in the Anti-Helmholtz configuration, a Keysight 90A (with current noise of $\sim 70\text{mA}$ rms in constant current operation) supply is used, and for the Helmholtz configuration, a Wavelength Electronics 15A supply is used.

¹A detailed explanation and design will be found in future thesis from the group by Elie Bataille.

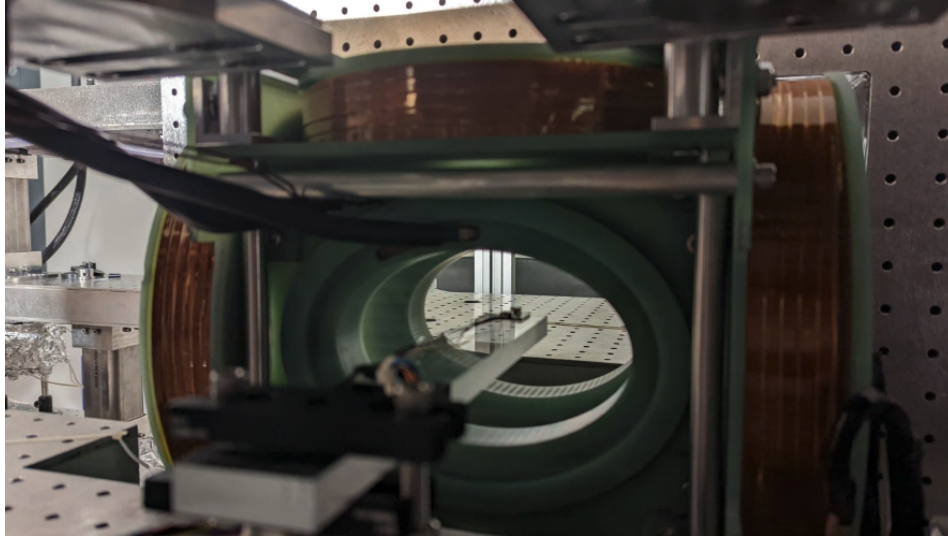


Figure B.2: **Measuring fields** Pictured is the location of the coils in the experiment, during an operation to zero the magnetic fields via measurements made using an Arduino. The glass cell is removed from the center of the fields for this measurement, via the translation stage.

Coil	Helmholtz Field Dis-uniformity across FOV	Helmholtz field strength (G/A)	Anti-Helmholtz field gradient (G/cm/A)
Low Field	9.86×10^{-5}	0.9386	0.1726
MOT Coil	3.89×10^{-5}	1.5268	0.2555
Objective Coil	4.1905×10^{-5}	1.3727	0.2231
Cell Coil	1.10×10^{-4}	0.8622	0.1240

Table B.2: **Coil Specifications for Helmholtz configuration.** For each field direction we include the Helmholtz field strength that can be obtained, the anti-Helmholtz field gradient that can be obtained, and the maximum dis-uniformity of the Helmholtz field across the objective FOV=1.5mm diameter plane. This dis-uniformity is the difference between the maximum and minimum field strength across the FOV normalized by the maximum field strength from the given set of coils.

We include in table B.1 the details of the coil winding geometry of the coil mounts for reference, where each mount is elliptically shaped with the inner major and minor diameters mentioned in the table. We include in table B.2 the magnetic field that we expect when each set of coils are operated in a Helmholtz configuration as well as the ratio of the difference in field across the objective FOV in that axis to the total field. We measure a resistance in the MOT coils of around 19mOhms, and measure a switch time of about 5ms after implementing adaptive methods to avoid overshoot

and damping. The geometry of the bore of the copper water cooled coil is pictured in Fig. B.3. The low-field coil is made of 16 WGS insulated round copper wire.

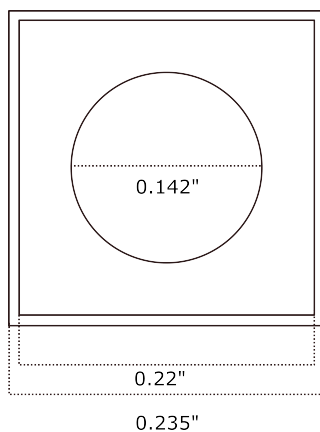


Figure B.3: **Water-cooled coil cross-section.** The dimensions of the cross-section of our copper water-cooled coils for reference. The inner bore is for water cooling, and the dimensions of the copper wire with and without insulation are labeled. The wires are coated in DDG insulation, and the water supply is de-ionized.

Appendix C

OBJECTIVE ALIGNMENT

We include here a step-by-step description of the steps that we have used for objective alignment relative to a glass cell in the science chamber. Two objectives are needed for alignment, but both do not have to be high-NA and optimized for tweezer trapping or glass corrected. Glass plates of the thickness of the cell, as well as of the thickness of the correction of the second objective are also required. One should do this alignment process at low power. We assume that we have an incoming beam from the tweezer pathway, as well as a ‘reference beam’ which we use for the alignment. The process begins with these two counter-propagating beams that are both aligned perpendicularly to the glass cell, and also centered on the center of the array.

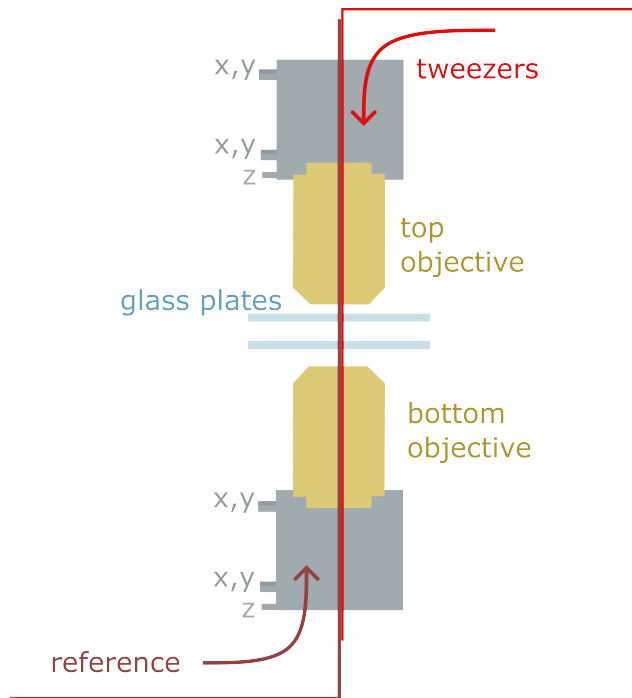


Figure C.1: **Objective alignment.** A schematic of the objective alignment including the reference beam, tweezer beam, objectives, glass plates, and alignment stages.

We note the significant importance of properly aligning the objective perpendicularly to the glass cell surface. In particular, we calculate that a slight misalignment of any more than a tenth of a degree, with our 6.35 mm thick glass and 0.65 NA objective

will result in tweezers whose depth is significantly compromised, as displayed in Fig. C.2 where we plot the Strehl ratio as a function of the tilt between the objective and the glass cell. The Strehl ratio is given by the ratio between the highest intensity of beam relative to the case of no aberration (in this case, then no tilt gives a Strehl ratio of 1.0), and here we calculate it numerically by utilizing ray tracing to propagate beams through the glass, and then a Debye-Wolf integral to calculate the beam interference near the focal plane (where we search over a 3D grid to find the maximal intensity).

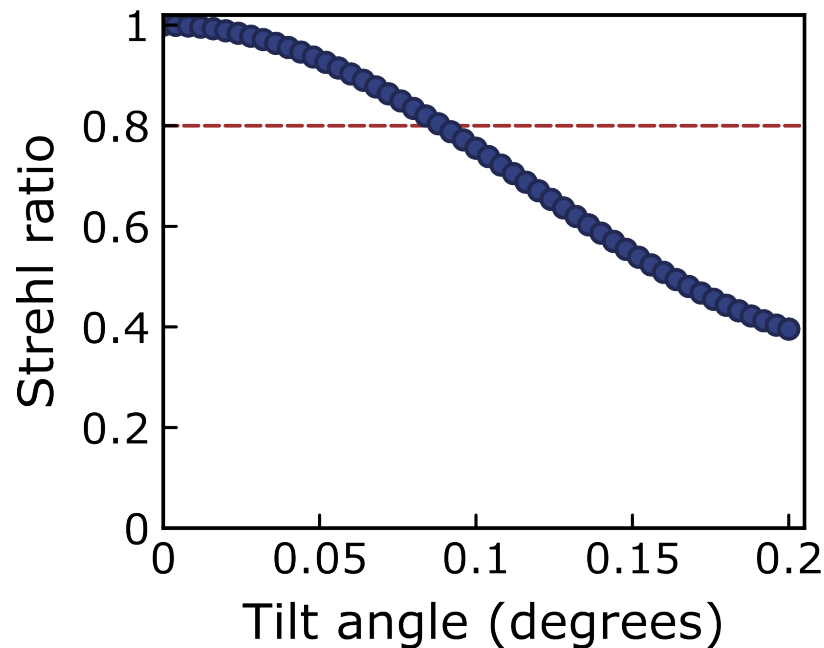


Figure C.2: **Impact of objective tilt misalignment.** The Strehl ratio (numerically calculated via ray tracing and a Debye-Wolf integral for each angle) as a function of tilt angle between the objective and the glass cell, presumably due to misalignment. The maximum intensity of the tweezers created drops below 80% (The Marechal Criterion: marked in dashed red) after around a tenth of a degree. This motivates extremely careful alignment of the objectives relative to the glass cell.

Step 0: Roughly align reference beams

Roughly align two reference beams to the anticipated optical axis for the objectives (without the objective being in place). One should be along the tweezer pathway and the other from the opposite direction.

Step 1: Align one reference beam perpendicular to glass plates.

With the glass cell in place, block one of the reference beams and use the other to

find the direction perpendicular to the glass plates. We do this in our experiment by looking at the reflection from the plates, and then optimize the angle of the reference beam such that the beam and its reflection are copropagating when we look very far from the glass cell along the incoming optical path.

Step 2: Remove glass cell and input glass plates.

Remove the glass cell now that the beam is aligned in angle to this reference. A glass plate should then be mounted on adjustable mounts such that it can be aligned to be perpendicular to the reference beam. Align the plate such that the reflection from the glass plates is aligned to be counter-propagating to the incoming beam (align by adjusting the plate angle, not the incoming beam). Ensure that the plate is mounted such that a magnetometer can fit to find the field center (next step). A second glass plate corresponding to the glass correction of the second objective (if it is the same as the first objective, then the glass will be also the thickness of the cell) can either be inserted and aligned now or else after the field center is found.

Step 3: Find field center.

In a tweezer experiment that is directly loaded from the MOT, we would like to center our array on the zero point of the quadrupole field used to make the MOT from which we will load the array. Thus, we use a magnetometer to measure the field and then mark the zero field position using a mounted thin metal needle. We note that we have already worked to ensure that the glass cell was fairly centered in the magnetic field when we aligned the coils.

Step 4: Align first reference beam.

After the field center is marked, and the glass plate is aligned, walk the beam such that it is both centered on the field center as well as co-aligned with its reflection. This may take some iteration.

Step 5: Co-align the second reference beam.

Next, co-align the incoming beam from the tweezer pathway to be co-propagating with the reference beam. The light should now be aligned both to the position and angle that we expect.

Step 6: Align objective mounts to reference beams.

Now the objective mounts can be aligned to the reference beams. Ideally, each mount is on a stage with two irises that, when closed, mark the center of the optical axis that the objectives will be mounted along. The stage should be a 5-axis stage, with movement of at least a few mm in each direction. Having this mount on picomotors is especially useful if possible. In our experiment, we do not have such remote control unfortunately, due to the weight of the objectives and the stages that we use. In any case, each stage should be adjusted such that the two irises are centered on the reference beams, which will provide a rough radial alignment for the objectives. Prior to setting the axial alignment of the stages/mounts, be sure to measure and account for the objective length and working distances, such that the proper distance is within the accessible range of the stage.

Step 7: Insert and align one objective

Mount one of the objectives in the now roughly-aligned mount (we'll call it the 'top' objective). Now, send the reference beam through the objective backwards *at low power and narrow*— use an iris as needed¹. (7a) Adjust the two knobs on the stage closer to the entrance aperture (top knobs) to overlap the beams above/before the objective. (7b) Then adjust the two knobs on the stage closer to the atom plane (bottom knobs) to overlap the beams on that side of the objective, near the focus. Iterate until beams are overlapped before and after the objective.

Step 8: Insert second objective

Mount the second objective ('bottom' objective) on the mount which was already previously centered on the reference beam in step 6. Use a power meter after the two objectives to optimize the axial distance between the objectives by optimizing transmission. (8a) Then adjust the two knobs on this stage that are further from the focal plane between the objectives (bottom knobs) to co-align the reference and tweezer beam after it has just passed through both objectives. (8b) Use the two knobs closest to the focal plane (top knobs) to align the reference beam after it has transmitted through both objectives and much further down the path align with the incoming beam. After this step, one should see both beams on both sides of both

¹Note that sending too much power can result in beams being focused on lenses, as objectives are not designed for collimated light to be sent from this direction at significant powers. A few milliwatts should be fine, a few watts will likely not be. Mistakes were made in this regard during the tenure of this work.

objectives.

Step 9: Fine tuning

Return to the top knobs on the top objective (from step 7), and use them to overlap the beams underneath both objectives close to the output. Use the bottom knobs of the top objective to overlap the beams significantly further down the reference pathway. Iterate between Steps 8 and 9 until convergence.

Step 10: More wavelengths

More wavelengths may be needed for alignment such as for an imaging pathway which should be aligned through the objective, but is split at a dichroic where the reference beam will not transmit. Additional beams can then be added. It is preferable that when setting up the reference beam to begin with, a cube is already included for combination of reference beams, such that adding additional wavelengths does not shift the existing reference beam.