

MEASURING THE RATIO OF MAGNETIC DIPOLE TRANSITION ($M1$) TO
VECTOR TRANSITION POLARIZABILITY (β) IN FRANCIUM-211

by

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Dedicated to my grandparents
Laxmi Chand Kapil, Santosh Rani

SUMMARY

Atomic parity violation (APV) in the weak interaction provides a pathway to test the Standard Model at a uniquely low momentum scale. APV measures the strength of highly forbidden atomic transitions induced by the parity-violating (PV) exchange of Z bosons between electrons and quarks in heavy atoms. An alkali atom is a suitable choice because of its simple atomic structure with a single valence electron, which allows for straightforward atomic physics calculations. The APV effect scales as $\sim Z^2N$, implying that heavy alkali atoms are ideal candidates due to their combination of simple structure and high atomic number. The best APV measurement so far has been performed in Cs, with an uncertainty of 0.35%. The APV signal in Fr is 18 times larger than in Cs.

Our group is working toward APV measurement in radioactive francium ($Z = 87$), the heaviest alkali, at ISAC (a radioactive beam facility) at TRIUMF, where we capture Fr atoms in a magneto-optical trap (MOT). We trap and cool francium (Fr) atoms in a dual MOT and perform atomic spectroscopy by driving the highly forbidden 7S–8S transition in francium. In the 7S–8S transition in Fr, the parity violation observable will be the interference between a parity-conserving amplitude—a Stark-induced E1 amplitude created by applying a DC electric field, E , to mix the S and P states—and the vastly weaker PV amplitude. The quantity of interest is the ratio of the parity-violating to parity-conserving amplitudes. These transitions are accompanied by a magnetic dipole (M1) transition, which vanishes in the non-relativistic limit.

The hyperfine interaction and relativistic effects give rise to a very faint M1 transition, which is 13 orders of magnitude weaker than allowed transitions. Understanding the contributions of the Stark-induced E1 and M1 amplitudes is crucial for optimizing the apparatus and improving the signal-to-noise ratio, enabling detection of small transition rates in APV measurements.

To drive the M1 transition, we intensified the 507 nm light in the interaction region 4000-fold using a power build-up cavity. We improved our detection efficiency by driving atoms that have undergone the 7S–8S excitation on a cycling transition up to ~ 5000 times. Overall, we have increased the signal rate by more than a million-fold. This puts us, in terms of statistics, within reach of taking APV data.

In this thesis, I will shine light on our francium trapping facility, highlight all the key features of the trapping apparatus, and describe the techniques and upgrades we have implemented in recent years to understand the contributions from distinct transition amplitudes in the 7S–8S forbidden transition rate. Furthermore, I will present recent results for the $M1/\beta$ measurement, which not only reveal interesting physical effects but also serve as a sensitive test of theoretical predictions.

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The journey of my PhD has been unforgettable, filled with invaluable life lessons and experiences I will always treasure. I am thankful to Timothy Hucko, from whom I gained practical knowledge related to FrPNc lab techniques and the handling of apparatus. I wish to express my gratitude to Luis Orozco for his guidance on numerous occasions, to Andrea Teigelhoefer for being consistently present from casual conversations to more serious discussions and for her assistance, and to John Behr for his insightful suggestions and for exemplifying dedication and passion in one's career. I also thank Alexandre Gorelov for his help with the neutralizer.

This journey also allowed me to build connections with many fellow colleagues, including Michael Kossin, Louis Croquette, Tasanul Morshed, Erik Freiling, Iris Halilovic, Liang Xie, Mukut Kalita, Gordon Arrowsmith, Ohene-Kofi Budu and others. I appreciate the staff at TRIUMF, especially the procurement team, and the staff at the University of Manitoba, particularly Susan Beshta, for their diligent efforts behind the scenes to support us.

I am deeply thankful to my family, including my parents, Trymbak Kapil and Anita Sharma, and my siblings, Garima Sharma and Abhoo Kapil, who have continuously inspired and supported me emotionally, financially, and in countless other ways. I also express my appreciation to my relatives, my uncle, aunts, and cousins, for their encouragement throughout this journey. Last but not the least, I want to honor Jagjit Singh, the late ghazal singer, whose music served as a source of calm and focus during challenging times.

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STATEMENT OF PERSONAL CONTRIBUTIONS

The research presented in this thesis is the result of several years of collaborative work carried out by many skilled researchers, including young researchers like me and other students. The FTF effort initially involved a team of about ten people (onsite and offsite). During my time, the core onsite team consisted of 3–5 members (a postdoctoral fellow, PhD students, and Master’s students), each working directly with the full experimental apparatus and contributing during beamtime.

The work in this thesis reflects techniques, skills, lessons, and infrastructure developed over many years, along with new additions made during my involvement in the project. I enrolled in the PhD program in 2018 and, after completing coursework, began continuous laboratory work in May 2019. Below are my contributions, learnings, and efforts toward the results presented in this thesis.

Summary of Personal Contributions

■ Experimental Leadership and Beamtime Responsibilities

- ☐ Served as team lead during the 2023, 2024, and 2026 beamtimes, including setting up the measurement scheme and running the full experiment to measure $M1/\beta$.
- ☐ Recorded and analyzed data to extract the relativistic and hyperfine-induced components of $M1$.
- ☐ Participated in and contributed to multiple beamtimes conducted at the francium trapping facility at TRIUMF.
- ☐ Trained a master’s student in core experimental laboratory skills and also super-

vised him in setting up and measuring the Stark shift in ^{87}Rb .

■ Major Experimental Operations Performed Independently

- ☐ Ran the full experiment independently: trapped atoms in the MOT, transferred the MOT, aligned the 507 nm beam in the PBC to the MOT, and aligned the depumper, burst beam, and reset beam.
- ☐ Ensured optimal MOT visibility on the CCD, reliable SiPM data acquisition with low background, and stable locking of the 507 nm beam in the PBC while chopping.
- ☐ Operated all components simultaneously to observe a transition thirteen orders of magnitude weaker than the allowed transition.

■ Optical Cavity Recovery, Alignment, and Mode Matching

- ☐ Addressed degradation of the Fabry–Perot cavity (PBC) in vacuum by maintaining lock during air introduction, recording degradation data, and assisting collaborator John Behr in controlled air leakage into the science chamber.
- ☐ Performed full alignment of the Fabry–Perot cavity with a build-up factor of 4000 from scratch.
- ☐ Mode-matched the 507 nm beam to the TEM_{00} mode, locked the cavity to the transmission peak, and ensured reliable relocking in a vibration-sensitive environment at a 20% duty cycle ($400\ \mu\text{s}$).

■ Laser Systems: Alignment, Optimization, and Upgrades

- ☐ Replaced two of the existing pump lasers and aligned them to the Ti:Sapphire laser, optimizing alignment and bringing the system to lasing.
- ☐ Installed a new HeNe laser into the existing trapping-laser locking setup.

- Optimized the burst beam (cycling transition $7S_{1/2} \rightarrow 7P_{3/2}$) with proper retro-reflection to minimize its impact on the MOT using live signal feed and maintain beam overlap. Installed intensity stabilization setup on burst beam with a colleague.

■ Vacuum System Work and Infrastructure Development

- Assisted in pumping down the system—guided by John Behr—by operating residual gas analyzers and activating the non-evaporable getter pump.
- Designed adaptor components for the optical pumping setup in SOLIDWORKS and designed the optical pumping tower which will be used to launch the push and optical pumping beams.
- Coordinated with TRIUMF onsite staff, including electrical services, 3D printing personnel, buyers, procurement staff, and technicians from Coherent who serviced the lasers onsite.

■ Scientific Output and Publications

- Performed the first $M1/\beta$ measurement for both $\Delta F = \pm 1$ hyperfine transitions in 2023 and analyzed the data. The thesis presents results from the 2024 beamtime, which show a discrepancy with theoretical predictions for the $\Delta F = -1$ transition, indicating the need for improved understanding of the underlying physics or systematics.
- Results obtained for $\Delta F = +1$ from the 2023, 2024, and 2026 beamtimes (all of which I led) toward measuring $M1_{\text{rel}}$ will form the basis of my first-author publication; this work is in progress.
- A second prospective publication, currently in manuscript preparation, concerns the method and results of repeated recovery of Fabry–Perot cavity degradation with air.

CHAPTER 1

ATOMIC PARITY VIOLATION

1.1 Motivation

History: Symmetry is something that fascinates our minds, especially our eyes. One example is parity (or mirror) symmetry, which implies that the system will behave in an identical way across a reversal of spatial coordinates. The belief in the conservation of parity in all physical interactions persisted until the mid-1950s, when the result of an experiment posed what came to be known as the ‘Tau-Theta’ puzzle, where a particle, now called the K^+ meson, had apparently two different decay modes, with opposite final state parities. Assuming parity conservation, this suggested two different particles forming a parity doublet of nearly identical mass. In 1956, Lee and Yang [2] proposed that the non-conservation of parity in the weak interaction could explain a single initial particle decaying into different final states of opposite parity. Within a year, definitive experimental evidence for parity violation (PV) in the weak interaction was found by Wu et al. [3] and others.

C. S. Wu performed one of the experiments suggested by Lee and Yang, the beta decay of polarized nuclei. They studied beta decay $^{27}\text{Co}_{60} \rightarrow ^{28}\text{Ni}_{60} + e + \bar{\nu}$ and observed the direction of the electrons emitted relative to the magnetic moment (or spin) of the parent nuclei. In Fig. 1.1, suppose that the upper image depicts the original decay where a beta particle emerges under an angle θ against the (negative) spin axis. If we observe instead the mirror image of this process, shown in the lower half, the spin (axial vector) points in the same direction, but the beta trajectory is reversed, corresponding to an angle of $180^\circ - \theta$. If mirror symmetry is valid (i.e. the image and its mirror image are indistinguishable), this implies that the rate of beta emission for these two directions must be identical. However, this was not observed by Wu and colleagues. They found that the beta particle is preferable

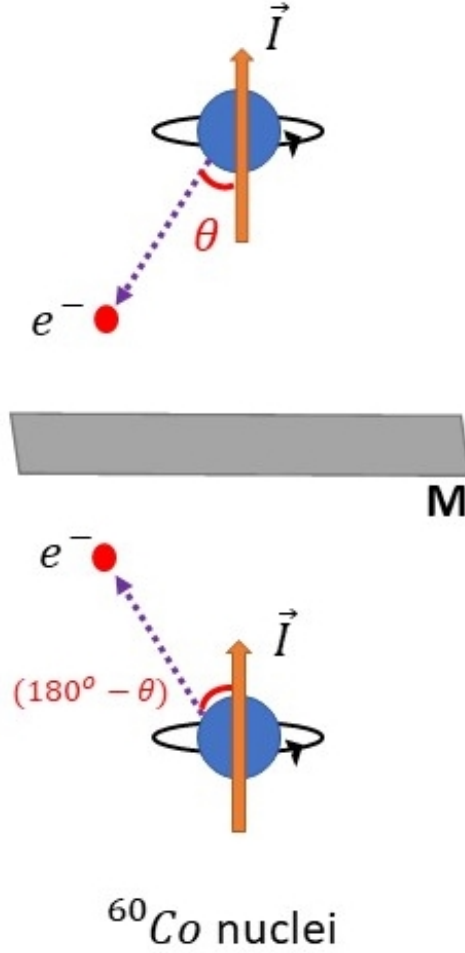


Figure 1.1: Effect of a mirror reflection on the beta decay of ^{60}Co nuclei.

emitted in the direction opposing the spin, requiring PV. Furthermore, it was found that parity is violated maximally.

As the name suggests, weak processes are often masked by the presence of electromagnetic and strong interactions. PV was unsurprisingly first observed in situations where the weak interaction is dominant, such as beta and charged pion decays.

Fermi's theory of the weak interaction (WI): Fermi in 1932 proposed a theory for the WI in analogy to the electromagnetic (EM) interaction. Consider the Lagrangian density, L_{EM} , for a simple case of an interaction between an electron and the radiation field:

$$L_{\text{EM}} = -e\bar{\psi}\gamma_\mu\psi A^\mu, \quad (1.1)$$

where ψ is the electron field, $e\bar{\psi}\gamma_\mu\psi$ is the EM current density j_μ^{EM} associated with ψ , γ_μ are the Dirac matrices, and A^μ is the four-vector potential of the radiation field. In this current-current interaction, Fermi substituted hadronic and leptonic currents in lieu of the electron and photon currents, and replaced the strength of the EM interaction, e , with that of the weak interaction, $G_F/\sqrt{2}$, where G_F is known as the Fermi coupling constant:

$$L_\beta = \frac{G_F}{\sqrt{2}} (\bar{\psi}_p \gamma_\mu \psi_n) (\bar{\psi}_e \gamma_\mu \psi_{\nu_e}). \quad (1.2)$$

Fermi treated the WI as a four-point interaction occurring at a single point in space-time, due to the absence of a propagator in the above expression (1.2). In the late 1950s, the concept of parity violation in β -decay was incorporated into the old Fermi theory of β decay, and Feynman and Gell-Mann (1956) [4] proposed that the factor responsible for particle transformation (γ_μ) could be a mixture of vector and axial-vector quantities. Under space inversions, a (polar) vector reverses direction, whereas an axial vector does not.

A combination of γ_μ and $\gamma_\mu\gamma_5 = \gamma_\mu i\gamma_0\gamma_1\gamma_2\gamma_3$ terms will introduce handedness and hence can describe parity non-conservation. This is known as ‘V-A theory’ [5]. Expression (1.2) can now be rewritten by replacing γ_μ by $\gamma_\mu(1 - \gamma_5)$. Fermi’s theory worked well for low-energy weak interactions, but the linear rise of cross sections with energy made this approach untenable at higher energies.

Introduction of the weak propagator to the theory: Fermi’s theory drew on an analogy with electromagnetism, where interactions occur via photon exchange between charged particles. By extension, the weak force should then also be mediated by vector particles. Since weak interactions are short-ranged, the mediators must be massive. Moreover, because β -decay alters the nuclear charge, the mediator must itself carry charge. Consequently, the weak force was understood to be mediated by charged vector bosons ($W^\pm$), giving rise to

what is termed the weak charged current. The Lagrangian for weak interaction becomes:

$$L_W = \frac{G_F}{\sqrt{2}} J_\mu^+(x) J^\mu(x), \quad (1.3)$$

where $J^\mu(x)$ is the charged current and can be utilized as a source of a charged vector boson $W^\pm$.

In 1954, C.N. Yang and R. Mills introduced a renormalizable theory for interacting massless particles known as a *gauge theory*. This theory could represent photons and W bosons if they were massless. *Spontaneous symmetry breaking* was introduced by Peter Higgs (1964) and others, now generally referred to as the *Higgs mechanism* in particle physics, permitting the introduction of massive boson mediators such as the $W^\pm$ into gauge theories.

Neutral weak interaction currents: In 1967-68, Salam, Weinberg, and Glashow developed a model now known as the *Standard Model* where they unified the electromagnetic and weak interactions into electroweak interactions by giving a single coupling constant, the elementary electric charge, e and the Weinberg angle, θ_W . There were four massless mediator particles: W^+ , W^- , W^0 , and B . W^0 and B mix to produce two physical particles: the photon and the Z^0 boson. The photon doesn't have any mass, and the Higgs mechanism can give comparable mass to the Z^0 boson as to the $W^\pm$ bosons. Thus, the neutral heavy vector Z -boson associated with *neutral weak currents* came into the picture. Having no charge, this boson acts in a way like a heavy photon, and indeed, in the presence of electromagnetic interactions, its presence is often entirely masked. Some processes like neutrino scattering with the exchange of Z^0 bosons is one of the examples where the weak neutral current effect is dominant because the neutrino does not couple to a photon.

1.2 PV induced by weak neutral currents in atomic physics

There are two possible different lepton-hadron processes that do not change the charge of the particles: (i) the exchange of a photon gives us an electromagnetic amplitude A_{EM} , and

(ii) the exchange of a Z^0 boson gives us a weak amplitude, A_W . The existence of neutral weak currents in atoms results in the coupling of a heavy neutral vector boson Z^0 to an atomic electron and the nucleons in two parts: a part that conserves parity by behaving as invariant under mirror reflection, A_W^{even} , and another part A_W^{odd} that creates parity-violating effects. The transition probability is:

$$P = |A_{\text{EM}} + A_W|^2, \quad (1.4)$$

where

$$A_W = A_W^{\text{even}} + A_W^{\text{odd}}. \quad (1.5)$$

The weak amplitude A_W is dramatically smaller than A_{EM} , making the former unobservable on its own. However, the interference between A_{EM} and A_W^{odd} provides the signature of parity violation, which can be picked out even in the presence of an overwhelming electromagnetic, parity-conserving, background.

Justification for calling weak interaction ‘weak’: We know that $A_{\text{EM}} \propto \frac{e^2}{q^2}$ and $A_W \propto \frac{g^2}{q^2 + M_{Z_0}^2 c^2}$ where q is the four-momentum transfer between the electron and the nucleus. The mass of the Z^0 boson ($\sim 90 \text{ GeV}/c^2$) [6] is heavier compared to zero mass photons making the range of EM interactions infinite, and the range of the WI, $\sim \frac{\hbar}{M_{Z_0} c} \sim 10^{-18} \text{ m}$, exceedingly short. Thus, an interaction with a Z^0 boson is very weak and of short range at low energies compared to the EM interaction.

Difficulty of Detecting Electroweak Processes in Atoms: The interference of the dominant A_{EM} amplitude with the parity violating A_W^{odd} amplitude gives rise to a left-right asymmetry, A_{LR} that can be defined as the normalized difference between the left and right-handed transition probabilities, P_L and P_R , arising from two different left and right experimental configurations posing as mirror images of each other. From expressions (1.4)

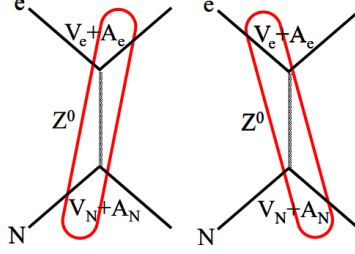


Figure 1.2: Vector electron-axial nucleon interaction, nuclear spin independent and dependent parts.

and (1.5), we find:

$$A_{\text{LR}} = \frac{P_L - P_R}{P_L + P_R} \approx 2 \left(\frac{A_W^{\text{odd}}}{A_{\text{EM}}} \right). \quad (1.6)$$

The smallness of $A_{\text{LR}} \sim \frac{\alpha^2 m_e^2}{M_{Z^0}^2} \sim 10^{-15}$ [6] makes it difficult to detect. The chance of its detection can be increased by reducing A_{EM} by using a highly forbidden transition.

1.2.1 Parity Violating Potential

The Lagrangian L_{e-N} from equation (1.3) for an interaction between the electron and the nucleus with the exchange of a heavy neutral vector boson Z^0 can be written as:

$$L_{e-N} = \frac{G_F}{\sqrt{2}} J_e^\mu J_\mu^N, \quad (1.7)$$

where the Fermi constant, G_F [7], is

$$G_F \approx 3 \times 10^{-12} m_e c^2 \left(\frac{\hbar}{m_e c} \right)^3, \quad (1.8)$$

where J_μ^e and J_N^μ are the neutral weak currents due to the fields of the electron and nucleus respectively and can be written as

$$J_e = V_e + A_e, \quad J_N = V_N + A_N, \quad (1.9)$$

where V is the vector current and A is the axial-vector current. The Hamiltonian H_{PV} has

to be pseudoscalar if it violates parity. Thus, we can write it from equation (1.7) as

$$H_{PV} = A_e V_N + V_e A_N = H_{PV}^{(1)} + H_{PV}^{(2)}. \quad (1.10)$$

Expanding the two parts of equation (1.10) by putting wave functions of electron, proton, and neutron as ψ_e , ψ_p , and ψ_n , respectively:

$$H_{PV}^{(1)} = A_e V_N = \frac{G_F}{\sqrt{2}} \sum_{p,n} \bar{\psi}_e \gamma_\lambda \gamma_5 \psi_e \left[C_{1p} \bar{\psi}_p \gamma^\lambda \psi_p + C_{1n} \bar{\psi}_n \gamma^\lambda \psi_n \right], \quad (1.11)$$

$$H_{PV}^{(2)} = V_e A_N = \frac{G_F}{\sqrt{2}} \sum_{p,n} \bar{\psi}_e \gamma^\lambda \psi_e \left[C_{2p} \bar{\psi}_p \gamma^\lambda \gamma_5 \psi_p + C_{2n} \bar{\psi}_n \gamma^\lambda \gamma_5 \psi_n \right], \quad (1.12)$$

where C_{1p} , C_{1n} , C_{2p} , and C_{2n} are the parity-violating vector and axial-vector coupling coefficients (C_1 , C_2) for protons and neutrons with electrons. At tree level they can be written as functions of the Weinberg angle θ_W as

$$C_{1p} = \frac{1}{2}(1 - 4 \sin^2 \theta_W), \quad C_{1n} = -\frac{1}{2}, \quad (1.13)$$

$$C_{2p} = \frac{1}{2}g_A(1 - 4 \sin^2 \theta_W), \quad C_{2n} = -\frac{1}{2}g_A(1 - 4 \sin^2 \theta_W), \quad (1.14)$$

where $\sin^2 \theta_W \approx 0.23$ and $g_A \approx 1.2754$ is the axial-vector coupling constant.

Now, substituting these values into equation (1.11) to cast $H_{PV}^{(1)}$ into a suitable form for the atomic calculations, we get

$$H_{PV}^{(1)} = \frac{G_F}{2\sqrt{2}} \bar{\psi}_e \gamma_\lambda \gamma_5 \psi_e \left[(1 - 4 \sin^2 \theta_W) \sum_p \bar{\psi}_p \gamma^\lambda \psi_p - \sum_n \bar{\psi}_n \gamma^\lambda \psi_n \right]. \quad (1.15)$$

Substituting $\sum_p \bar{\psi}_p \gamma^\lambda \psi_p = Z \bar{\psi}_N \gamma^\lambda \psi_N$ and $\sum_n \bar{\psi}_n \gamma^\lambda \psi_n = \sum_n N \bar{\psi}_N \gamma^\lambda \psi_N$ in the above equation, we get

$$H_{PV}^{(1)} = \frac{G_F}{2\sqrt{2}} Q_W \left(\bar{\psi}_e \gamma_\lambda \gamma_5 \psi_e \right) \left(\bar{\psi}_N \gamma^\lambda \psi_N \right), \quad (1.16)$$

where Q_W is known as the weak charge which is equal to $(1 - 4\sin^2\theta_W)Z - N$. We now simplify the expression by using the non-relativistic approximation for nucleons, which offers a more intuitive form. The u_{N_L} and u_{N_S} represent the large and small components of spinors respectively, as follows:

$$\psi_N = \begin{pmatrix} u_{N_L} \\ u_{N_S} \end{pmatrix} = \begin{pmatrix} u_N \\ \frac{\vec{\sigma}_N \cdot \vec{p}_N}{2Mc} u_N \end{pmatrix}. \quad (1.17)$$

The $\vec{\sigma}_N$ is the Pauli matrix spin operator acting on the spin of the nucleon, and the $\vec{p}_N$ is the three-momentum operator acting on the nucleon wavefunction. Starting from the normalization condition $\psi_N^\dagger \psi_N = \rho_N$, where the latter is the nucleon density, and using the approximation $|u_S| \approx \frac{1}{2} \frac{V}{c} |u_L|$, where V is the speed of the electron we see that the small component contributes only a relativistic correction to the total density. Expanding equation (1.15) in a suitable expression by putting $\gamma_\lambda = (\gamma_0, \gamma_k)$ where $k = 1, 2, 3$ and $\psi^\dagger = \bar{\psi}\gamma_0$, we get

$$H_{PV}^{(1)} = \frac{G_F}{2\sqrt{2}} Q_W \left[\bar{\psi}_e^\dagger \gamma_5 \psi_e \rho_N(\vec{r}) + \bar{\psi}_e \gamma_k \gamma_5 \psi_e \bar{\psi}_N \gamma^k \psi_N \right]. \quad (1.18)$$

We can neglect γ^k , $k = 1, 2, 3$ in the sum over λ when mixing u_L and u_S , and the small value of V can be neglected. To first order, we can take the nucleus as point-like, $\psi_N^\dagger \psi_N = \rho_N(\vec{r}) = \delta(\vec{r})$. Thus, $H_{PV}^{(1)}$ is left with only the time component of $(A_e V_N)$ and equation (1.18) becomes

$$H_{PV}^{(1)} = \frac{G_F}{2\sqrt{2}} Q_W \left[\bar{\psi}_e^\dagger \gamma_5 \psi_e \right]_{r=0}. \quad (1.19)$$

$H_{PV}^{(1)}$ can be written in a useful form as a potential, V_{PV} , by treating the electron as non-relativistic for low values of Z as:

$$\langle a | H_{PV}^{(1)} | b \rangle = \frac{G_F}{4\sqrt{2}} \frac{Q_W}{m_e c} \left[u_a^* (\vec{\sigma}_e \cdot \vec{p}_e) u_b + (\vec{\sigma}_e \cdot \vec{p}_e) u_a^* u_b \right]_{r=0}. \quad (1.20)$$

The exchange of the heavy neutral vector boson Z^0 gives rise to the parity-violating electron-nucleus potential, V_{PV}^{e-N} , and this potential can be taken as a sum of two terms: one dependent on nuclear spin and the other independent of nuclear spin in the non-relativistic limit. The former term is relatively small for sufficiently large atomic number, $Z \gtrsim 50$. Thus, nuclear spin and nuclear momentum are neglected to get V_{PV}^{e-N} with zero-range which is invariant under time reversal and can be written as

$$V_{PV}^{e-N} = \frac{G_F}{\sqrt{2}} \frac{Q_W}{4m_e c} \left(\vec{\sigma}_e \cdot \vec{p}_e \delta^3(\vec{r}) + \delta^3(\vec{r}) \vec{\sigma}_e \cdot \vec{p}_e \right), \quad (1.21)$$

where m_e , σ , p , and r are the mass, spin, momentum, and position of the electron, respectively. We neglected the small contribution from the short-range electron-electron PV potential. The weak charge Q_W reveals the coherent contribution from all the nucleons in $V_{PV} \approx H_{PV}^{(1)}$. The weak charge is a model-independent, dimensionless quantity that is a function of (Z, N) . The presence of $\vec{\sigma} \cdot \vec{p}$ in the above equation introduces a handedness to the system.

Parity Mixing in Single-Particle Wave Functions: The V_{PV}^{e-N} potential mixes atomic wave functions of different parities, specifically between s - and p -states. The $ns_{1/2}$ wavefunction is

$$\psi_{ns_{1/2}}(r) = \frac{1}{\sqrt{4\pi}} R_{n0}(r) \chi, \quad (1.22)$$

and for $n'p_{1/2}$ we have

$$\psi_{n'p_{1/2}}(r) = \frac{1}{\sqrt{4\pi}} R_{n'1}(r) Y_l^\mu(\theta, \phi) \chi, \quad (1.23)$$

where $R_{n0}(r)$ and $R_{n'1}(r)$ are the radial functions of the s and p states respectively, χ is a two-component Pauli spinor describing the spin state of the electron and Y_l^μ are spherical harmonics with $l = 1$. The mixing of these states results in a new *mixed wave function* that involves both the s - and p -states.

1.3 Matrix Elements and Enhancement Mechanism

The second term in equation (1.21) gives zero contribution to the matrix element because at the origin, the electron in the s-wave state has zero momentum and the (non-relativistic) p-wave function vanishes at $r = 0$. The Z -scaling is unaffected by the angular momentum distribution, so we can substitute $\tilde{\sigma} \cdot \tilde{\mathbf{p}} = \frac{\hbar \mathbf{p}_r}{2}$. The matrix element of the V_{PV}^{e-N} potential then is

$$\langle ns_{1/2} | V_{PV}^{e-N} | n'p_{1/2} \rangle \approx \frac{G_F}{\sqrt{2}} \cdot \frac{Q_W}{4m_e c} \cdot \frac{1}{4\pi} \int \psi_{ns_{1/2}}(r) \left[\delta^3(r) \cdot \frac{\hbar}{i} \frac{\partial}{\partial r} \right] \psi_{n'p_{1/2}}(r) d^3r. \quad (1.24)$$

We know $\chi^2 = 3$ and using expressions (1.22) and (1.23), we get

$$\langle ns_{1/2} | V_{PV}^{e-N} | n'p_{1/2} \rangle \approx \frac{3i}{16\pi m_e c} \frac{G_F}{\sqrt{2}} Q_W(Z, N) R'_{n'1}(0) R_{n0}(r), \quad (1.25)$$

where $R'_{n'1}(r) = \frac{d}{dr} R_{n'1}(r)$. The matrix element is imaginary to maintain the time-reversal invariance (T). In 1974, Bouchiat and Bouchiat [7] pointed out the approximate Z dependence of APV effects using the Fermi-Segre semi-empirical formula (non-relativistic),

$$R_{nl} \sim r^l Z^{l+1/2} \quad (\text{large } Z). \quad (1.26)$$

Only R_s and $\frac{d}{dr} R_p$ do not vanish at the origin, implying that mixing occurs only between s- and p-states.

(i) In heavy atoms, the wave function of an s-electron is more concentrated near the nucleus because the increasing nuclear charge Z exerts a stronger Coulomb attraction, and s-electrons, having no angular momentum barrier, can penetrate the inner regions of the atom where electron shielding is minimal. Non-relativistically, the value of the wavefunction at the origin scales as $Z^{1/2}$, while relativistic effects enhance this further.

(ii) The second contributing factor is that the electronic orbital near the nucleus is approxi-

mately Coulombic, which implies that the electron velocity scales as $v_e \propto Z$. The derivative of the $p_{1/2}$ wavefunction at the origin contributes $Z^{3/2}$.

(iii) The weak nuclear charge is approximately $Q_W \approx -N$, since the proton contribution is suppressed by the small factor $1 - 4\sin^2 \theta_W \approx 0.08$. For most stable nuclei, the neutron-to-proton ratio satisfies $N/Z \gtrsim 1$, so $Q_W \propto Z$, providing an additional factor of Z in the scaling of the weak interaction amplitude.

This results in a Z^3 enhancement of PV effects, making heavy atoms the optimal systems for detecting atomic parity violation (APV).

Possible Case to Look for APV: It is challenging to perform the atomic physics calculations required to establish the connections between the electroweak theoretical parameters and experimentally measured transitions. This constraint limits the exploration to heavy atoms with the simplest electronic structure, such as heavy alkali atoms like cesium, which has only one valence electron outside closed shells.

Parity Mixing Effects in Atomic States: In the absence of parity-violating effects, atomic states have definite parity, and electric dipole (E1) transitions between states of the same parity are strictly forbidden. However, the mixing of s and p states due to H_{PV} implies a small E1 transition amplitude between such states, e.g. between two s states:

$$|\widetilde{nS}\rangle = |nS\rangle + \sum_{\alpha} \frac{|\alpha P\rangle \langle \alpha P | H_{PV} | nS \rangle}{E_{nS} - E_{\alpha P}}. \quad (1.27)$$

H_{PV} is assumed to be small and a treatment in first order perturbation theory is sufficient. The electric dipole transition $E1_{PV}$ between the mixed atomic states is

$$E1_{PV} = \langle \widetilde{n'S} | -e\vec{c} \cdot \vec{r} | \widetilde{nS} \rangle, \quad (1.28)$$

$$E1_{PV} = \sum_{\alpha} \left[\frac{\langle n'S | H_{PV} | \alpha P \rangle \langle \alpha P | \hat{E}1 | nS \rangle}{E_{n'S} - E_{\alpha P}} + \frac{\langle n'S | \hat{E}1 | \alpha P \rangle \langle \alpha P | H_{PV} | nS \rangle}{E_{nS} - E_{\alpha P}} \right]. \quad (1.29)$$

The first and second terms represent the perturbation of the final and initial states by H_{PV} respectively. Both H_{PV} and $\hat{E}1$ are P -odd, meaning that for a non-zero value of $E1_{PV}$, the unperturbed states $|n'S\rangle$ and $|nS\rangle$ must have the same parity.

1.4 Parity Violation Observables

There are two observables: one measuring interference in transition amplitudes, and the other measuring optical rotation, which also arises from interference effects. The former probes the transition probability directly, while the latter measures the phase shift between polarization components of the transmitted light, details of both methods are discussed below.

1.4.1 PV-Interference Effects in Atomic Transitions

Bouchiat suggested to observe the interference of the PV transition amplitude with another, larger, parity conserving (PC) transition amplitude. This PC amplitude could be a magnetic dipole (M1) and/or a Stark-induced E1 amplitude. Thus, the rate of transition can be written as

$$I \sim |A_{PC} + E1_{PV}|^2 = |A_{PC}|^2 \pm 2|A_{PC}||E1_{PV}| + |E1_{PV}|^2. \quad (1.30)$$

The interference term will change sign upon a parity reversal, allowing us to separate out the PV signature.

Interference of $E1_{PV}$ and M1: Before discussing parity violation, it is useful to consider the allowed magnetic dipole (M1) transition between states of the same parity. The matrix element for the M1 amplitude is:

$$\langle n'S | \hat{M}1 | nS \rangle \sim -\frac{e\hbar}{2m_e c} \langle n'S | \mathbf{L} + 2\mathbf{S} | nS \rangle = \mathcal{M}(n'S, nS) \hat{o}_{m_{n'S} m_{nS}}, \quad (1.31)$$

where $\mathcal{M}(n'S, nS)$ is the reduced matrix element. When a small parity-violating electric dipole transition $E1_{PV}$ is present, it can interfere with the allowed M1 amplitude, the in-

terference produces circularly polarized fluorescence from excitation with linearly polarized light, or a rotation of transmitted plane-polarized light. It leads to a degree of circular polarization P_c , even if the incoming radiation is linearly polarized:

$$P_c = \frac{2 \operatorname{Im}(E_{PV}\mathcal{M})}{|E_{PV}|^2 + |\mathcal{M}|^2} \approx 2 \operatorname{Im}\left(\frac{E_{PV}}{\mathcal{M}}\right). \quad (1.32)$$

The interference of $E1_{PV}$ and M1 can also be interpreted as introducing different transition probabilities for right- and left-handed circularly polarized light.

1.4.2 Interference of $E1_{PV}$ and Stark-Induced E1 electric dipole transition

An external static electric field $\mathbf{E}_S$ mixes the unperturbed atomic states with different parities, leading to an electric dipole transition between nS and $n'S$ states via a parity conserving interaction. The electric dipole transition amplitude $E1_{ST}$ due to the Stark-induced interaction is:

$$E1_{ST} = \langle n'S | -e\vec{\mathbf{E}}_S \cdot \vec{\mathbf{r}} | nS \rangle. \quad (1.33)$$

$$E1_{ST} = \sum_{\alpha} \left[\frac{\langle n'S | \hat{E}1_S | \alpha P \rangle \langle \alpha P | \hat{E}1 | nS \rangle}{E_{n'S} - E_{\alpha P}} + \frac{\langle n'S | \hat{E}1 | \alpha P \rangle \langle \alpha P | \hat{E}1_S | nS \rangle}{E_{nS} - E_{\alpha P}} \right]. \quad (1.34)$$

The interference of the small $E1_{PV}$ amplitude with a Stark-induced electric dipole amplitude $E1_{ST}$ can be measured in Stark interference experiments. Bouchiat and Bouchiat [8] showed that absorbing a circularly polarized photon perpendicular to a static electric field induces a non-zero electronic angular momentum $\langle \mathcal{J} \rangle$ in the excited state, revealing parity-violating (PV) effects.

Optical Rotation Effects: Light of different handedness propagates with different velocities in the optically active medium resulting into different refractive indices for the left and right-handed polarized light. This is referred to as circular birefringence. It causes the rotation of the polarization plane due to retardance of right-handed circular component by different amount than the left-handed circular component or vice versa. The polarization

plane of polarized light of wavelength λ is rotated in a distance l by an angle ϕ because of differing phase velocities c/n is [9]

$$\phi = \frac{\pi l}{\lambda}(n_R - n_L) \quad (1.35)$$

or

$$\phi = \frac{\pi l}{\lambda}(n_R - n_L) = \frac{4\pi l}{\lambda}(n - 1)Im\frac{\mathcal{E}_{PV}}{\mathcal{M}}. \quad (1.36)$$

The PV optical rotation where the observable quantity is ϕ follows the dispersive lineshape of $(n - 1)$ and can be distinguished by the dispersive shape of $(n - 1)$ as a function of frequency. The experiments exploiting this observable are referred to as optical rotation experiments. These were performed in bismuth by groups working in Seattle, Oxford, Novosibirsk, and Moscow in the late 1970s. More details are given below.

1.5 Experimental Observation of PV Effects

1.5.1 Optical Rotation Experiments

In 1978, the first PV signal was observed in optical rotation experiments by the Novosibirsk group in Bi [10]. The rotation of the polarization plane by very minute angles, $\Delta\phi_{PV} \sim 10^{-7}$ rad. Similar experiments had started after 1974 in Seattle [11], Oxford [12], and later in Moscow [13]. All the experiments initially chose Bi in search for optical activity because of the presence of allowed M1 absorption lines that were accessible with tunable lasers available at the time. Later, experimenters also used lead and thallium.

Experimental Scheme for Optical Rotation Experiment The simplest scheme for the experimental setup consists of a laser beam which passes through a setup of a polarizer P and crossed analyzer A, with an atomic vapor sample (say Bi) in between. The optical rotation of the sample by an angle ϕ_P and transmission intensity I are measured while sweeping the laser frequency across the M1 absorption line.

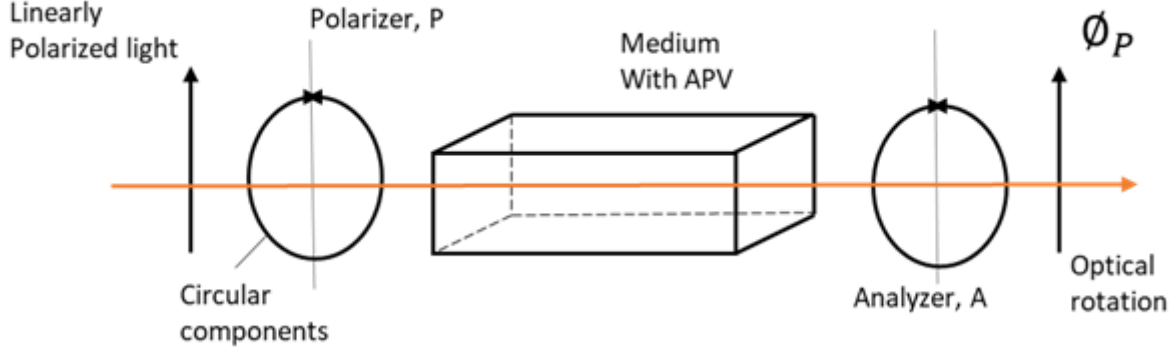


Figure 1.3: Basic setup for optical rotation method to explore PV effects.

Measurement of Small Angle of Rotation: A very small rotation of the plane of polarization of light between P and A will produce a large fractional change in I transmitted by A. For $\phi_P \ll 1$,

$$I = I_0 \sin^2(\phi_P) = I_0 \phi_P^2, \quad (1.37)$$

where I_0 is the intensity of incident light on the A. It is very difficult to detect a very small change in the light intensity, so, a reference angle, ϕ_M , is added such that $I = I_0 \sin^2(\phi_P + \phi_M)$, where ϕ_M is modulated with a Faraday cell which rotates the plane of linearly polarized light as a function of an applied magnetic field. Thus, the intensity of transmitted light is

$$I = A(\phi_P + \phi_M + \phi_R)^2 + B, \quad (1.38)$$

where A and B are constants, and ϕ_R is any residual angle caused by the misalignment of polarizers etc. The I_M component of I reverses sign with ϕ_M and depends on $\phi_M(\phi_P + \phi_R)$. The effect of ϕ_P can be separated from that of ϕ_M because ϕ_P changes rapidly and in a characteristic way as the frequency approaches resonance. The observation of a small rotation angle needs a long accumulation time of signal, thus drifts of the optical elements and electronics make it difficult in practice. This puts a limitation on resolving the small rotation in a sufficiently short time.

Overview of One of the Optical Rotation Measurements: S J Phipp et al. [14] measured parity violating optical rotation in atomic lead on the $1.270\ \mu\text{m}$ magnetic dipole transition. The basic setup of the experiment was the same as discussed above, with an oven and a Faraday modulator placed between polarimeters. A semiconductor laser source with an external cavity was used. The rotation signal was obtained by phase-sensitive detection on the low noise InGaAs detector by recording the transmitted laser light which has the fundamental and second harmonic of the modulation frequency. The other sources of frequency-dependent rotation came from the inhomogeneous and time-varying vapour distribution inside the oven, atomic absorption of off-mode laser light. These undesirable effects were taken care of by computer modelling the effects with additional theoretical outlines in the fitting procedure. They found an optical rotation of

$$R_{Pb_{208}} = (-9.80 \pm 0.33) \times 10^{-8} \text{ rad.} \quad (1.39)$$

1.5.2 Stark-Interference Experiment

The first generation of Stark-interference experiments were carried out in Paris [15] and Berkeley [16] in the 1980s where the basic setup of the experiment was to send a circularly polarized laser beam into an atomic vapour cell and transitions of interest were excited in the presence of a dc electric field. The transition rate was monitored by observing the fluorescence of atoms that decay back to the ground state. The polarization of the excited state was detected by looking at the degree of circular polarization of the fluorescence light and thus, parity violation was measured with fractional uncertainty of $10 - 20\%$. The drawback of these experiments was the low signal to noise ratio due to background signals arising from collision induced transitions, scattered light, and blackbody radiation.

Principle of the experiment: The atoms from the $6S_{1/2}$ state of Cs were excited to the $7S_{1/2}$ state, using $540\ \text{nm}$ laser light. The detection used the allowed electric dipole

$7S_{1/2} - 6P_{1/2}$ transition at 1360 nm. A controllable static electric field, $\vec{E}_S$, was applied. The Cs vapour was excited by a circularly polarized laser beam with a wave vector $\vec{k}$ orthogonal to $\vec{E}$, where the sign of projection of the angular momentum $\vec{\Lambda}_i$ of the photon along $\hat{k}$ that is helicity $\zeta = \pm 1$ characterized the circular polarization: $\Lambda_i = \zeta \hbar \vec{k}$. Suppose a mirror M perpendicular to $\hat{k}$ of incident beam with ζ reflects the incident beam and thus, due to reflection, the sign of $\hat{k}$, a polar vector, has changed whereas the $\vec{\Lambda}_i$, an axial vector does not change. The helicity, a pseudo scalar, changes sign during reflection. Thus, the mirror reflection of incident laser beam resulted into a new laser beam characterized by $-\hat{k}$ and $-\zeta$, while $\vec{\Lambda}_i$ remains unchanged. In a reflection through the plane xOy perpendicular to $\hat{k}$, the new laser configuration formed by two laser beams remain unchanged. So, if parity is conserved, the fluorescence light emitted at 1360 nm remains unchanged in this reflection. But the fluorescence light emitted along a direction orthogonal to $\hat{k}$ and $\vec{E}$ is circularly polarized revealing that the emitted photon along this direction has angular momentum $\delta\vec{\Lambda}_f$. This circularly polarized fluorescence is a chiral effect because the sign of $\delta\vec{\Lambda}_f$ changes in a reflection through the xOy plane, and this chiral effect gives the proof of an atomic parity violation.

Stark-interference experiment with the best result so far: At JILA in Boulder, Colorado, in the 1980s a new experimental approach was thought of to increase the signal-to-noise ratio by achieving high detection efficiency and low background. They moved from a vapor cell approach to atomic beams and lifted the degeneracy of m magnetic sub-levels by applying an external magnetic field.

The A_{PV} amplitude is $\propto m$ whereas the A_{ST} amplitude $\propto |m|$. Summed over all m states, the interference term, $A_{PV}A_{ST}$ vanishes, but if individual m levels can be probed, PV becomes observable. This approach was used in Colorado in the 1988 experiment [1] by applying a large magnetic field ~ 74 G to resolve the magnetic states. In their subsequent work described below, they used a smaller field but optically pumped the atoms into a specific m level.

Why Cesium (Cs) and the 6S - 7S transition? The Z^3 dependence of PV mixing makes Cs, the simplest heavy atom with only one valence S-state electron outside closed shells, a favourable case for an APV experiment. The simplicity of atomic structure makes the atomic calculations simpler, accurate and more reliable than heavy atoms with several valence electrons. After choosing Cs, one needs an atomic transition between atomic states with suitable quantum numbers, and for Stark-interference method, the transition from the ground state, $6S$ to the first excited S state is the most natural candidate.

The transition between a particular ground state ($6S, F, m$) and the excited state ($7S, F', m'$) level was excited with a laser beam in the presence of mutually perpendicular static electric and magnetic fields. There are three relevant transition amplitudes which can give rise to a $6S - 7S$ transition as follows:

Electric dipole transition amplitude, A_{PV} : $A_{PV}(F, m \rightarrow F', m')$ arises due to the presence of weak neutral currents which contaminate the $6S$ and $7S$ states with states of opposite parity, P . This mixing in the presence of the oscillating laser beam with electric field polarization ϵ results into a small, parity-violating, amplitude A_{PV} between the two S states:

$$\begin{aligned}
A_{PV}(F, m \rightarrow F', m') &= \sum_n \left(\frac{\langle 7S F' m' | H_{PV} | nP \rangle \langle nP | -e \epsilon \cdot \mathbf{r} | 6S F m \rangle}{E_{7S} - E_{nP}} \right. \\
&\quad \left. + \frac{\langle 7S F' m' | -e \epsilon \cdot \mathbf{r} | nP \rangle \langle nP | H_{PV} | 6S F m \rangle}{E_{6S} - E_{nP}} \right) \\
&= i \operatorname{Im}(E1_{PV}) \epsilon \cdot \langle F' m' | \boldsymbol{\sigma} | F m \rangle, \tag{1.40}
\end{aligned}$$

where H_{PV} is the PV Hamiltonian, and σ is the Pauli spin operator. The constant $E1_{PV}$ contains the information about the radial integrals.

Stark-induced amplitude, A_{ST} : A dc electric field, E_S , is applied to mix states of opposite

parity leading to a parity-conserving $6S - 7S$ electric dipole amplitude A_{ST} :

$$A_{ST}(F, m \rightarrow F', m') = \sum_n \left(\frac{\langle 7S F' m' | e \mathbf{E}_S \cdot \mathbf{r} | nP \rangle \langle nP | -e \boldsymbol{\varepsilon} \cdot \mathbf{r} | 6S F m \rangle}{E_{7S} - E_{nP}} + \frac{\langle 7S F' m' | -e \boldsymbol{\varepsilon} \cdot \mathbf{r} | nP \rangle \langle nP | e \mathbf{E}_S \cdot \mathbf{r} | 6S F m \rangle}{E_{7S} - E_{nP}} \right). \quad (1.41)$$

The Bouchiat [17] showed that this amplitude can be parametrized as two components:

$$A_{ST}(F, m \rightarrow F', m') = \alpha \mathbf{E}_S \cdot \boldsymbol{\varepsilon} \delta_{F, F'} \delta_{m, m'} + i\beta (\mathbf{E}_S \times \boldsymbol{\varepsilon}) \cdot \langle F' m' | \boldsymbol{\sigma} | F m \rangle, \quad (1.42)$$

where the coefficients α and β are the scalar and vector transition polarizabilities, respectively. The α and β coefficients are given by:

$$\alpha = \frac{1}{6} \sum_n \left[\langle 7S || r || nP_{1/2} \rangle \langle nP_{1/2} | r | 6S \rangle \left(\frac{1}{E_{7S} - E_{nP_{1/2}}} + \frac{1}{E_{6S} - E_{nP_{1/2}}} \right) - \langle 7S || r || nP_{3/2} \rangle \langle nP_{3/2} | r | 6S \rangle \left(\frac{1}{E_{7S} - E_{nP_{3/2}}} + \frac{1}{E_{6S} - E_{nP_{3/2}}} \right) \right], \quad (1.43)$$

$$\beta = \frac{1}{6} \sum_n \left[\langle 7S || r || nP_{1/2} \rangle \langle nP_{1/2} | r | 6S \rangle \left(\frac{1}{E_{7S} - E_{nP_{1/2}}} - \frac{1}{E_{6S} - E_{nP_{1/2}}} \right) + \langle 7S || r || nP_{3/2} \rangle \langle nP_{3/2} | r | 6S \rangle \left(\frac{1}{E_{7S} - E_{nP_{3/2}}} - \frac{1}{E_{6S} - E_{nP_{3/2}}} \right) \right]. \quad (1.44)$$

The α term drives the $\Delta F = 0, \Delta m = 0$ transitions with $\mathbf{E}_S$ orthogonal to $\mathbf{E}$, and the β term excites the $\Delta F = 0, \pm 1, \Delta m = 0, \pm 1$ transitions with $\mathbf{E}_S$ parallel to $\boldsymbol{\varepsilon}$, respectively. The α term is about 10 times larger than the β term in Cs.

Magnetic dipole transition amplitude, A_{M1} : The oscillating magnetic field of the laser, $\mathbf{B}_{ac}$, induces a magnetic dipole transition between the $6S$ and $7S$ states, given below:

$$A_{M1}(F, m \rightarrow F', m') = \mu_B \langle 7S F' m' | \boldsymbol{\sigma} \cdot \mathbf{B}_{ac} | 6S F m \rangle, \quad (1.45)$$

where μ_B is the Bohr magneton, and can be written with the corresponding oscillating electric field ϵ and wave vector k of the laser as $B_{ac} = \hat{k} \times \epsilon$. Thus, the above equation can be written as

$$A_{M1}(F, m \rightarrow F', m') = M_1(\hat{k} \times \epsilon) \cdot \langle F' m' | \sigma | F m \rangle, \quad (1.46)$$

where $M1$ is a constant incorporating radial integrals stemming from relativistic effect (M) and the hyperfine interaction, M_{hf} :

$$M1 = M \pm M_{hf} \delta_{F, F' \pm 1}. \quad (1.47)$$

The total transition rate, R , can be written in terms of the sum of these three transition amplitudes:

$$R(F, m \rightarrow F', m') \propto |A_{ST} + A_{PV} + A_{M1}|^2. \quad (1.48)$$

The sign of the interference term between the Stark-induced and PV amplitudes will change with parity reversals, e.g. flipping the direction of the electric field, and this change is the experimental PV signature.

Coordinate Configuration The direction of propagation of the laser, $\mathbf{k}$, the electric field, $\mathbf{E}_S$, and the magnetic field were along the y , x , and z axes, respectively, as shown in Fig. 1.4a. The Wieman group has chosen elliptically polarized light of the form $\epsilon = \epsilon_z \hat{z} + \epsilon_x \hat{x}$, where ϵ_z is the component orthogonal to $\mathbf{E}_S$, whereas ϵ_x is the component parallel to $\mathbf{E}_S$. The ϵ_z and ϵ_x are defined to be real and imaginary, respectively, such that the A_{PV} and A_{ST} have the same phase. Thus, the rate of transition for $F \neq F'$ while neglecting the terms that involve only A_{PV} and A_{M1} is

$$R(F, m \rightarrow F', m') = \left(C_{F, m}^{F', m'} \right)^2 \left[\beta^2 E^2 \epsilon_z^2 \mp 2\beta E \epsilon_z \text{Im}(E 1_{PV}) \epsilon_x \right] \delta_{m, m' \pm 1}. \quad (1.49)$$

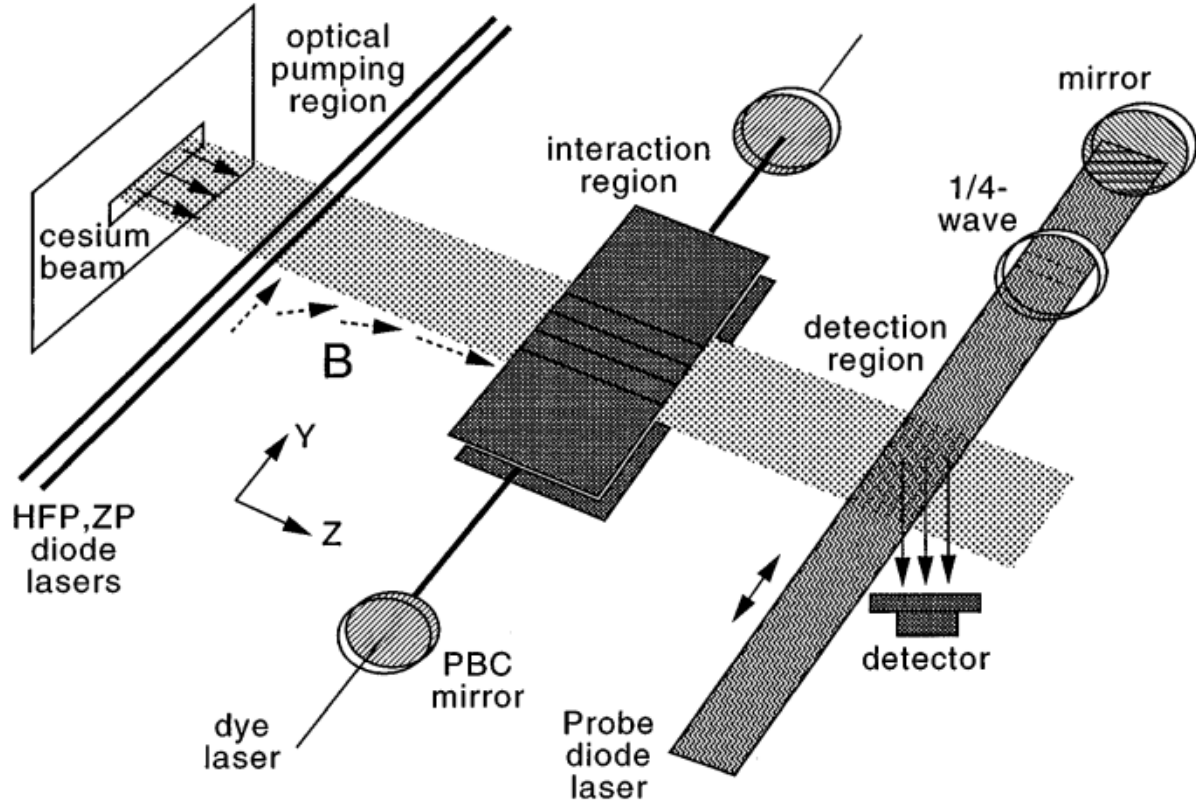
The first term in the above expression is the pure Stark-induced transition rate. The second term represents the $A_{ST} - A_{PV}$ interference, which changes its sign with electric field reversal,

magnetic (m) sublevel, and polarization (p) reversal of the laser beam. The quantity iE_{PV} is the PV electric-dipole reduced matrix element, and $C_{F,m}^{F',m'}$ are related to $\langle F'm'|\sigma|Fm\rangle$. The signal of interest is the fractional modulation of the total signal rate with reversals of E , B , m , and p . This fractional modulation is given by:

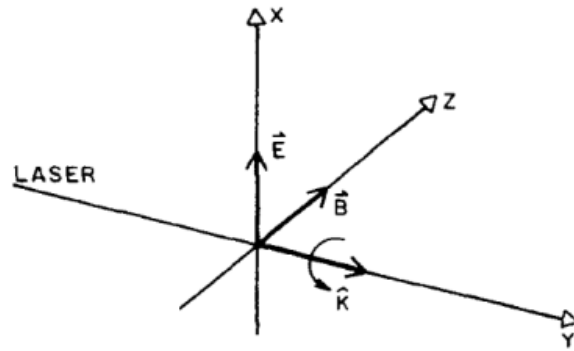
$$\frac{\Delta R}{R} = \frac{R_+ - R_-}{R_+ + R_-} \propto \frac{\text{Im}(E1_{PV})}{\beta E}, \quad (1.50)$$

where R_+ and R_- refer to the observed rates for a given handedness of the coordinate system. If the relative rate change is measured and β is known, the PV amplitude can be deduced.

Layout of the Cesium Experiment: Fig. 1.4b shows that the atomic beam of cesium comes out of a heated oven inside a vacuum chamber. In the optical pumping region, the atoms are placed in a single hyperfine state by using the hyperfine pumping (HFP) laser and Zeeman pump (ZP) laser to put all the atoms in a single m level. The spin-polarized beam enters the interaction region of mutually perpendicular magnetic, electric, and laser fields. A standing wave is created in a high-finesse Fabry-Perot etalon which increases the power of the laser beam, thus called the power build-up cavity (PBC). The laser beam at 540 nm is provided by a tunable dye laser, which is kept on the $6S - 7S$ resonance by locking it to an ultra-stable external reference cavity. The PBC is in turn adjusted in length to be in resonance with the laser, using the Pound-Drever-Hall scheme [1]. The DC electric field is produced by a pair of parallel plates that are divided into five isolated segments aiding to apply small uniform and gradient electric fields for greater control and for diagnostic tests. The magnetic field is provided with coils of various shapes to give the required fields and gradients. The atomic beam enters the detection region where the probe diode laser scatters photons from the atoms that repopulated the depleted hyperfine state after making the $6S - 7S$ transition. A photodiode with a large area collects the scattered photons, and the photocurrent is proportional to the number of atoms that got excited on the $6S - 7S$ transition.



(a) Cesium experimental setup.



(b) Field configuration for Cs experiment.

Figure 1.4: Overview of the Cs experiment and field configuration [1].

The measured value of $\text{Im}(E1_{PV})/\beta$ is $-1.6349(80)$ mV/cm for $(6SF = 4 \rightarrow 7SF' = 3)$ and $-1.5576(77)$ mV/cm for $(6SF = 3 \rightarrow 7SF' = 4)$. This experiment achieved a fractional uncertainty of 0.35% in measuring the PV amplitude for the $6S - 7S$ transition in Cs with a spin-polarized Cs beam, the most accurate APV measurement to date.

1.6 Implication for the Standard Model

The weak charge Q_W can be extracted from $E1_{PV}$. This involves atomic theory, essentially the behaviour of the electron's wavefunction close to and inside the nucleus, and also requires nuclear knowledge, in particular the nuclear charge radius as well as its neutron radius. Q_W is given by:

$$Q_W = ZQ_W(\text{proton}) + NQ_W(\text{neutron}) = (2Z + N)Q_W(u) + (Z + 2N)Q_W(d), \quad (1.51)$$

where N and Z are the number of neutrons and protons, respectively. The precise determination of Q_W will help determine the fundamental constants C_{1u} and C_{1d} in expressions (1.11) and (1.12) describing the PV electron-hadron interaction, which can be related to Q_W as

$$C_{1u} = -\frac{1}{2}Q_W(u), \quad C_{1d} = -\frac{1}{2}Q_W(d). \quad (1.52)$$

The empirical determination of Q_W defines an allowed band in the (C_{1u}, C_{1d}) plane. To determine the basic parity-violating electron quark couplings individually, other experiments are needed that probe a different combination of up and down quarks. The recent determination of the proton's weak charge using PV electron scattering on hydrogen at Jefferson Lab [18] provides the most stringent, complementary information. Together these two measurements constrain C_{1u} and C_{1d} , in agreement with Standard Model predictions, as shown in Fig. 1.5.

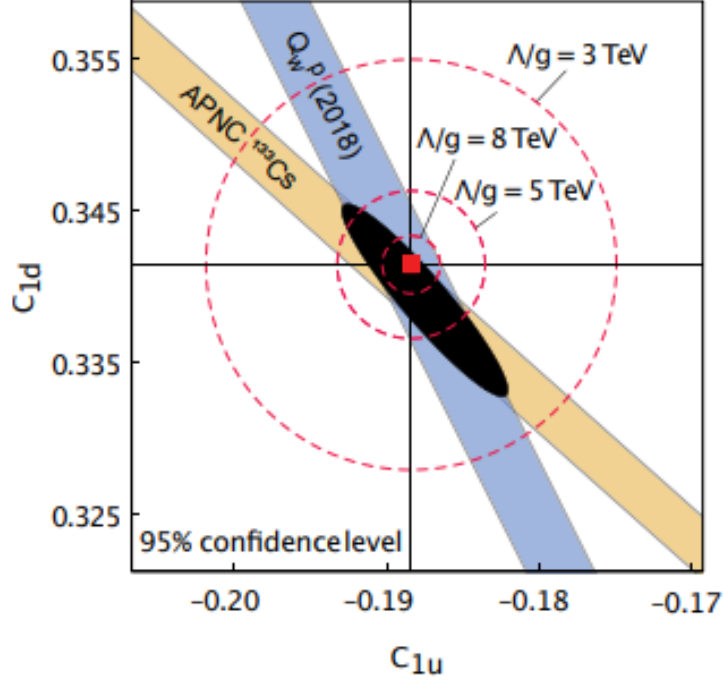


Figure 1.5: Constraints on C_{1u} and C_{1d} from parity violating electron scattering and atomic parity violation. The red dot is the Standard Model prediction. Taken from [18].

1.7 Ongoing and Upcoming APV Efforts

Experimental efforts are being made to improve the precision of APV measurements in Cs or other heavy atoms by different groups [19].

Cesium: Elliot et al. [20] at Purdue University are pursuing an improved measurement in cesium. They use a new two-pathway coherent control technique to measure very weak highly forbidden optical transition amplitudes. The two optical excitations, directed by two different mutually-coherent fields, start from the same initial state $6S_{1/2}$ to the final state $7S_{1/2}$. The allowed two-photon E1 amplitude is driven by two photons, one provided by a laser resonant with the $6S_{1/2} \rightarrow 7S_{1/2}$ transition and the other by a laser at half the resonant frequency. The interference term between a sum of Stark-induced and PV amplitudes and the two-photon E1 amplitude gives the absorption rate which depends on the relative phase of the applied laser fields. One can observe the oscillating modulation of the transition rate

by experimentally varying the relative phase.

Ytterbium, Yb [21] — Yb has two valence electrons and it is being investigated [22] with a goal of exploring the isotopic dependence of Q_W .

Dysprosium, Dy [23] — Another candidate is Dy where progress is being made towards searching for the nearly degenerate states of opposite parity, and later for APV.

Francium, Fr: The Z^3 enhancement of PV effects in atoms and the need for structural simplicity makes francium, the heaviest alkali atom, an interesting candidate for APV measurement, despite the fact that it has no stable isotope. An effort to this end is underway in our group and is presented in more detail in later chapters.

Spin-dependent APV: In addition to nuclear spin-independent APV one can study the nuclear spin *dependent* part, which in heavy atoms is dominated by the nuclear anapole moment [24]. This is accomplished by measuring the *difference* in $E1_{PV}$ between the the $\Delta F = \pm 1$ hyperfine components of the $ns \rightarrow n's$ transition.

Molecules: De Mille et al. [25] laid out a plan to use Stark-interference technique to measure spin dependent APV effects by determining the mixing between opposite-parity hyperfine/rotational levels of ground state molecules. Efforts with BaF [26] are ongoing.

1.8 Summary

Atomic parity violation (APV) arises from the weak interaction, which does not conserve mirror symmetry. Parity was believed to be universal until the mid-1950s, when the tau–theta puzzle and the beta-decay experiment of Wu et al. demonstrated clear violation of parity in weak processes. These results led to the electroweak theory, where weak interactions are mediated by massive vector bosons; in atoms, the relevant mediator is the neutral Z^0 boson whose exchange between electrons and nucleons produces a tiny parity-violating amplitude. In atoms, the electromagnetic amplitude A_{EM} dominates, while the weak amplitude A_W is

extremely small, yet their interference yields measurable PV effects. The weak interaction introduces a parity-violating potential that mixes atomic states of opposite parity, mainly s and p states, generating a small electric dipole amplitude $E1_{PV}$ between states of the same parity—a transition otherwise forbidden by parity selection rules.

APV effects scale roughly as Z^3 due to increased electron density near the nucleus, relativistic enhancements, and the weak charge $Q_W \approx -N$. Heavy alkali atoms such as cesium and francium, with simple electronic structure, are therefore ideal systems for APV studies.

Experimentally, APV is detected through (i) interference between $E1_{PV}$ and a larger parity-conserving amplitude (M1 or Stark-induced E1), or (ii) optical rotation arising from weak-induced mixing of polarization components. Both methods probe extremely small weak neutral-current effects.

In cesium, APV was measured via interference between the tiny $E1_{PV}$ amplitude and a Stark-induced amplitude, yielding the most precise determination of the weak charge. Experiments in other heavy atoms (Yb, Tl, Pb, Dy, Fr) continue to investigate APV through optical rotation and transition-interference techniques, with future work aiming at higher precision, isotope comparisons, and enhanced sensitivity in radioactive alkali systems such as francium.

CHAPTER 2

TECHNICAL APPROACH

Atomic parity violation (APV) is enhanced by a factor of $\approx Z^3$ [27], making heavy atoms ideal candidates for APV. However, performing the necessary atomic physics calculations to connect electroweak theoretical parameters with experimentally measured weak amplitudes in APV experiments is challenging. This difficulty has been limiting the exploration of heavy atoms to those with the simplest electronic structures, typically single-valence s-electron alkali atoms. We seek a suitable candidate that carries a unique combination of a high Z and a single electron in the outermost shell.

2.1 Francium — the heaviest alkali atom

Francium (Fr) is a radioactive alkali metal with $Z = 87$. It is the heaviest alkali metal and the least stable of the first 103 elements. It has 35 observed isotopes, ranging from $A = 193$ to 231 [28], which provides an advantage to testing fundamental symmetries with different isotopes. The longest-lived isotope has a half-life of only 23 minutes. Only isotopes 221 and 223 occur naturally in minute amounts through uranium and neptunium decay chains, and at any given time, only amounts on the order of grams exist on Earth. Therefore, francium is produced artificially with fusion-evaporation reactions, fragmentation, spallation, and other ways [28]. Being radioactive, Fr nuclei are unstable and decay into daughter nuclei by emitting radiation such as alpha, beta, or gamma rays. We obtain francium isotopes for our work at the ISAC radioactive beam facility at TRIUMF, described below.

2.1.1 What and where is TRIUMF?

TRIUMF is Canada’s national laboratory for nuclear and particle physics, operated by an association of Canadian universities. TRIUMF was named the ‘Tri-University Meson Facil-

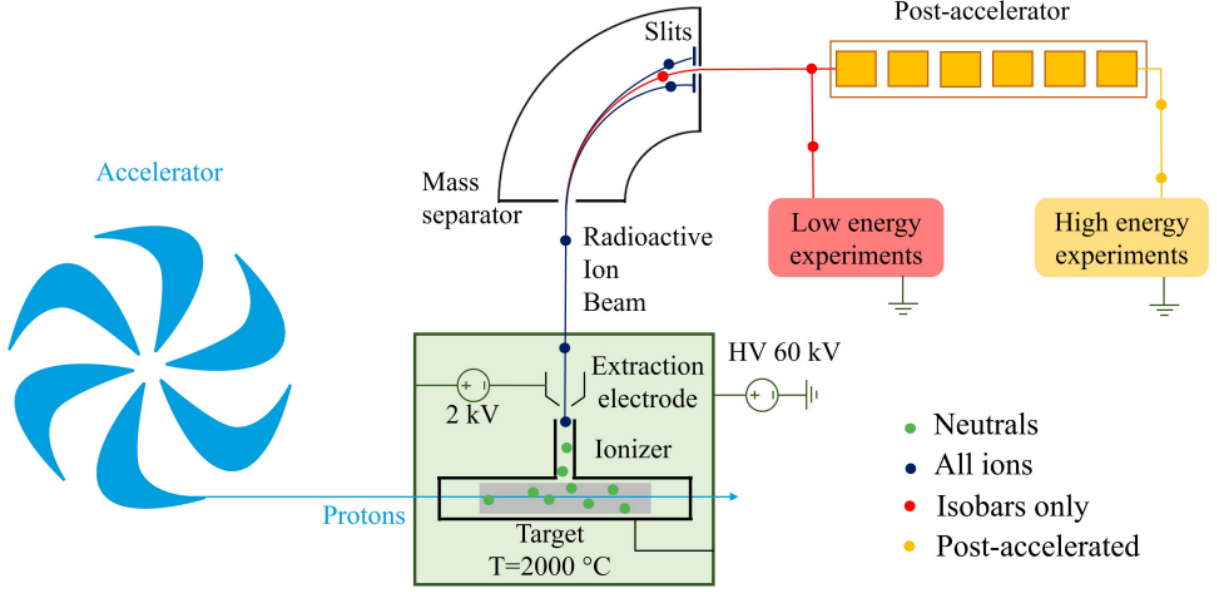


Figure 2.1: Schematic of radioactive ion beam production and delivery with a brief description of the process in section 2.1.1. The francium trapping facility is located in the low-energy experimental area. Image taken from [30].

ity’ and started as a joint venture of three universities: the University of British Columbia (UBC), the University of Victoria, and Simon Fraser University. It now has 21 member universities with the University of Manitoba as one of them. The TRIUMF facility is located on an 11 acre site on the south campus of UBC. It houses one of the world’s largest cyclotrons, capable of delivering high-intensity proton beams (initially accelerated as negative hydrogen ions) at energies up to 520 MeV. Its uniqueness lies in its ability to extract several proton beams for different energies and intensities [29]. The proton beam is extracted by intercepting the H^- beam with a thin carbon foil that strips off the two electrons. The change in charge bends the beams out of the cyclotron into beamlines that deliver the beams to users.

Production and delivery of Francium: At TRIUMF, francium isotopes are produced and delivered by the ISAC (Isotope Separator and ACcelerator) radioactive beam facility, using the Isotope Separator OnLine (ISOL) method. In this method, a $\lesssim 100\mu A$ beam of 520 MeV protons produced from the TRIUMF H^- cyclotron hits a uranium carbide (UC_x) production target, initiating a nuclear reaction releasing highly energetic reaction products.

The nuclear reaction is a spallation reaction in which an incoming particle strikes the target nucleus with enough energy to knock out multiple small fragments, resulting in the formation of lighter isotopes. The target material is heated, and the reaction products diffuse to the surface and get ionized for extraction of a radioactive ion beam (RIB) [29]. The RIB is steered and focused with electromagnetic lenses and sent towards a mass separator dipole magnet to select a required mass. From here, the RIB comes to our facility. The beam delivery requires many components, such as lenses and magnets, to facilitate the beam tuning. To date, ISAC is a single user facility; hence, only a single experiment can be served at a time. The francium experiment can expect about 5 days of dedicated beam per year. The Advanced Rare IsotopE Laboratory (ARIEL), currently under construction, calls for three concurrently operating targets. This will become critical for the feasibility of APV measurements with francium.

Choosing a Francium Isotope: Our choice of francium isotope is guided by three main requirements: **(1)** the isotope must be suitable for accessing the highly forbidden $ns_{1/2} \rightarrow (n+1)s_{1/2}$ transition, and its near-spherical nuclear shape, which results in simpler nuclear parameters, is advantageous for accurate theoretical nuclear structure calculations [31]; **(2)** it must have sufficiently high production rates for our experimental needs; and **(3)** its decay chain must not lead to significant long-lived residual activity in the facility. While the presence of long-lived nuclei is a concern for long-term contamination, the dominant radiation-safety consideration during beamtime is the short-term activity arising from rapid decays of the parent and daughter nuclei, particularly gamma radiation, since the emitted α and β particles do not escape the experimental apparatus.

We explore the highly forbidden atomic $nS_{1/2} \rightarrow (n+1)S_{1/2}$ transition, which has an oscillator strength of approximately 10^{-10} to 10^{-17} [32, 33], depending on the specific transition type discussed in Chapter 1. For reference, allowed atomic transitions typically have an oscillator strength of 0.1–1. This requires operating francium beams at the highest available intensities to obtain sufficiently large signals. We therefore work with ^{211}Fr , which has a

half-life of $t_{1/2} = 3.10$ minutes. This nuclide decays into ^{207}At by emitting an α particle with a branching ratio greater than 80%, and into ^{211}Rn via electron capture with a branching ratio below 20%. The decay chain of ^{211}Fr , shown in Figure 2.2, proceeds through ^{207}At and

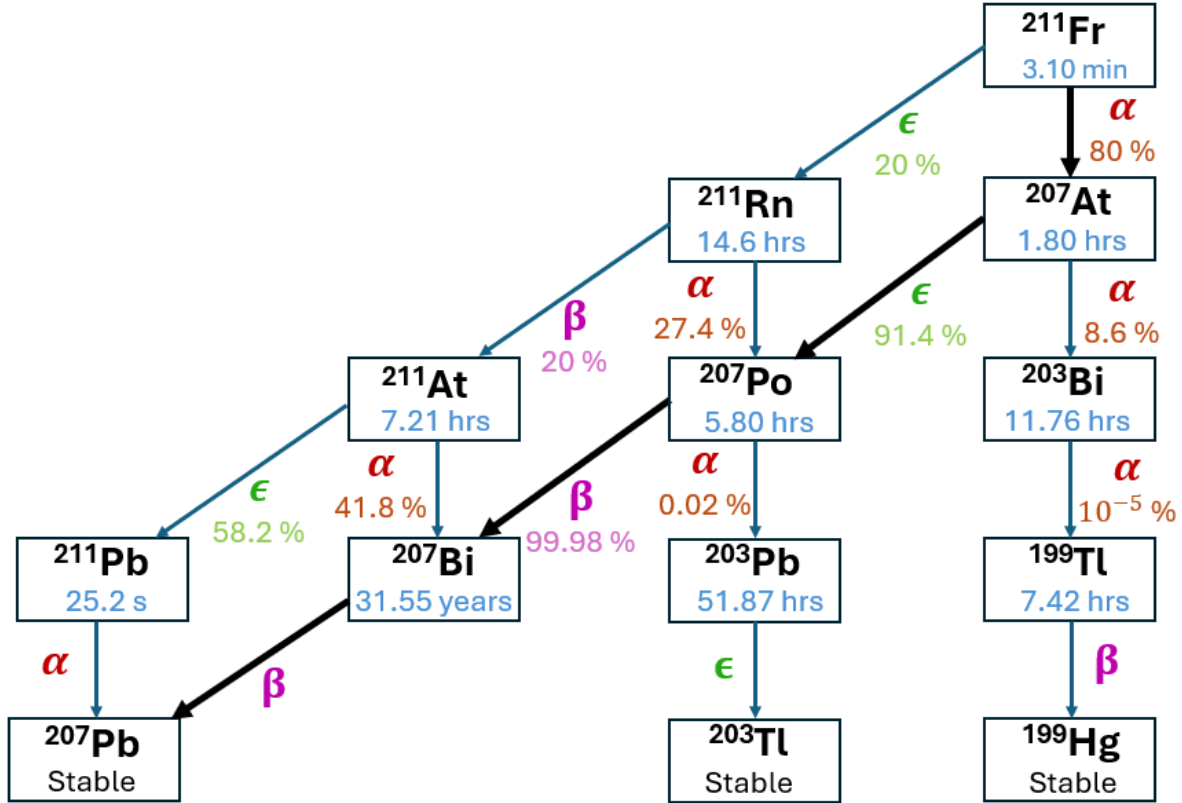


Figure 2.2: Decay chain of $^{211}_{87}\text{Fr}$ with their half lifetimes. The bold thick line shows the most probable decay route with the highest relative percentage of decays.

terminates largely at stable ^{207}Pb via a sequence of α and β decays. Although ^{223}Fr has a longer half-life ($t_{1/2} = 22$ min), ^{211}Fr is better suited for our experiment because its decay characteristics result in lower operational radiation risk and reduced accumulation of long-lived activity in the experimental area. The primary daughter of ^{211}Fr is ^{211}Ra , which has a short half-life of approximately 14.6 hrs, whereas ^{223}Fr decays to ^{223}Ra , a much longer-lived isotope with a half-life of about 11 days. While the properties of the immediate daughter are important, radiation safety considerations must account for the entire decay chain, particularly the presence and production rate of long-lived nuclides. The decay chain of ^{211}Fr

does include ^{207}Bi , which has a half-life of 32 years, implying that a small amount of long-term residual activity can persist after beam operation. However, the long-lived products display much lower activity than the initial fast decays. Consequently, their contribution to the radiation dose during typical beam delivery periods of a few days is negligible. During operation, the radiation field in the francium experimental area is instead dominated by the rapid decays of short-lived daughters, which generate high instantaneous activity while the beam is on but decay away quickly once the beam is turned off.

In contrast, isotopes such as ^{223}Fr and ^{209}Fr are less favourable from an activity-management standpoint. Their decay chains contain long-lived daughters such as ^{223}Ra (11 days), ^{207}Bi (32 years), and ^{209}Po (102 years), which can accumulate over successive beam deliveries and lead to elevated residual radiation levels. This buildup increases cool-down times, limits operational flexibility, and poses long-term challenges for radiation safety and facility operation. By selecting ^{211}Fr , we accept the presence of small amounts of long-lived residual activity while significantly reducing the dominant operational radiation hazard associated with short-term activity during beamtime. This short-lived decay structure allows higher beam intensities and longer experimental runs without exceeding facility radiation limits, in accordance with the ALARA principle¹. Although the gamma dose from ^{211}Fr accumulates over time, it remains within acceptable limits for typical experimental campaigns. Beyond radiation considerations, ^{211}Fr offers several physics and technical advantages. Its near-spherical nuclear structure leads to simpler nuclear parameters, which is advantageous for accurate theoretical modelling [34]. Furthermore, ^{211}Fr has the highest production yield among francium isotopes available at ISAC, exhibits large hyperfine splittings favourable for laser trapping, and the trapped atoms remain confined for a sufficiently long time to allow repeated measurements, enabling efficient and reliable experimental operation.

Need to trap atoms in a magneto-optical trap (MOT): The Boulder group for the cesium experiment used an atomic beam with $\approx 10^{13}$ atoms/s ([35], pages 15, 57). Even if

¹ALARA: annual dose limit for non-nuclear energy workers is 1 mSv/year.

the parity asymmetry in Fr is an order of magnitude larger, the typical ^{211}Fr delivery rate at ISAC of $\approx 1 \times 10^8$ atoms/s is clearly far from sufficient for a beam experiment. However, in a beam each atom gets used only once in the measurement. By using atomic cooling and trapping techniques, we can re-use each francium atom hundreds or even thousands of times. Wood estimates ([35], page 37) that they had $\approx 10^{10}$ 6S–7S excitations per second. This can be matched with 10^7 trapped atoms, if each is used for around 1000 excitations. Taking into account the larger parity asymmetry, even 10^6 atoms are feasible. To date, we have been able to routinely trap up to a million Fr atoms in our MOT in 1 mm^3 of volume for more than ten seconds.

2.2 Cooling and trapping of francium

A Magneto-Optical Trap (MOT) is a combined approach of optical and magnetic fields used to cool and trap neutral atoms at high densities, typically at tens to hundreds of microkelvin. A key step in this process is the deceleration of fast-moving atoms using laser-induced forces. When an atom absorbs a photon, the conservation of momentum causes a change δv in the atom’s velocity given by

$$\delta v = \frac{h\nu}{Mc}, \quad (2.1)$$

where M is the atomic mass, h is Planck’s constant, and ν is the frequency of the laser light. Although the atom re-emits the photon via spontaneous emission, the emission is isotropic on average, producing no net momentum change. This cycle of absorption and emission gives rise to the *scattering force*, which acts along the direction of the incident beam and can be used to slow the atoms [36, 37]. A high rate of momentum transfer is required to decelerate the atoms over a short distance and a short interaction time. To make this force velocity-dependent, the laser frequency is red-detuned slightly below the atomic transition frequency so that, due to the Doppler effect, atoms moving toward the beam perceive the light as closer to resonance and preferentially absorb photons from that direction, enabling effective laser cooling [38]. The MOT typically employs three pairs of counter-propagating laser

beams (six in total) aligned along orthogonal axes. When red-detuned (δ), the Doppler shift ensures that atoms moving in one direction absorb more photons from the beam opposing their motion. This arrangement forms a viscous medium known as *optical molasses*, which slows atoms and centers them within the trap. In the MOT, a position-dependent force is added by applying a spatially varying magnetic field gradient across the trapping region through the Zeeman effect. The magnetic field shifts the atomic energy levels depending on the atom's position. Using circularly polarized light, atoms preferentially absorb photons from the beam that push them back toward the trap center, creating a *restoring force*. The velocity-dependent scattering force along with the position-dependent restoring force enable simultaneous cooling and trapping of atoms. A large trapping volume and high laser power improve capture efficiency by increasing the trap depth and the number of trapped atoms. The alkali atoms are ideal candidates for MOTs due to their atomic structure, the ability to produce atomic beams or vapors, the availability of suitable lasers, and the presence of simple repeated optical transitions between the ground and excited states that allow efficient photon absorption and emission. These transitions are separated by several linewidths of the atomic transition, which ensures that only selected transitions are driven. The excited atom can decay to the ground state via spontaneous emission, resulting in a high momentum transfer. Consider a one-dimensional MOT in Fig. 2.3, where a weak magnetic field varies linearly with position. The Zeeman effect splits the magnetic sublevels, with the field proportional to the distance from the trap center [38]. When circularly polarized laser beams with σ^+ and σ^- polarization, at frequency ω_L , propagate in the $+z'$ and $-z'$ directions, respectively, they drive transitions between magnetic sub-levels. In this setup, an atom cloud moving in the $+z'$ direction sees the σ^- polarization, and the scattering force acts in the $-z'$ direction. In contrast, an atom moving in the $-z'$ direction absorbs photons from the σ^+ polarization, resulting in a scattering force that acts in the $+z'$ direction. Thus, the time-averaged force arising from the combined Doppler and Zeeman shifts pushes the atoms towards the origin, which we define as the center of the trap. The resonance is maintained by choosing the laser

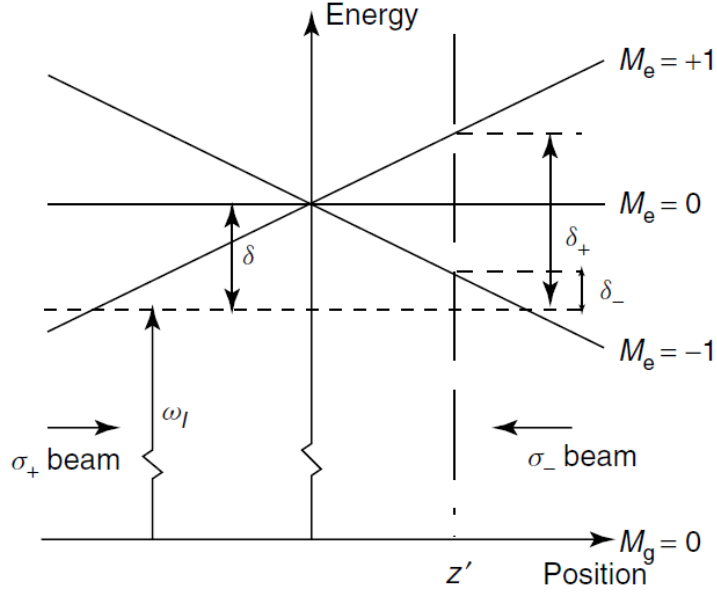


Figure 2.3: MOT in one dimension. σ^+, σ^- polarized laser beams; ω_L is the laser frequency; δ is the detuning.

detuning such that the sum of the Doppler and Zeeman shifts compensates for the laser frequency offset. At high velocities, Doppler cooling dominates and acts as a viscous drag force, while at lower velocities optical pumping into specific magnetic sub-levels becomes more significant, providing a position-dependent restoring force.

2.3 Neutralization of Fr ions with neutralizer

Because the atoms trapped in the MOT are neutral, we need to neutralize the francium ions into neutral francium atoms. Francium is delivered as singly charged ions from ISAC at energies of 20 – 30 keV and is implanted into our catching foil. Periodically, the catching foil is heated to elevated temperatures to release neutral francium atoms. Because the foil serves to neutralize implanted ions at the surface, it is hereafter referred to as the neutralizer. The embedded Fr^+ ions diffuse through the neutralizer material toward the surface, with the diffusion process itself strongly temperature dependent. Upon reaching the surface, the relative fraction of released ions and neutral atoms is determined by the equilibrium of

surface ionization, which follows the Langmuir–Saha equation [39]:

$$\frac{n_{\text{ion}}}{n_{\text{atom}}} = \frac{1}{2} \exp\left(\frac{E_{\text{WF}} - E_{\text{IP}}}{k_{\text{B}}T}\right) \quad (2.2)$$

where $\frac{n_{\text{ion}}}{n_{\text{atom}}}$ is the ratio of ions and neutrals leaving the surface, k_{B} is the Boltzmann constant, T is the absolute temperature of the neutralizer foil, and E_{WF} and E_{IP} are the work function and ionization potential of foil, respectively. From equation (2.2) for foil, the atom density becomes greater than the ion density when the energy required to remove an electron from an isolated atom $E_{\text{IP}} = 4.07$ eV is greater than the energy required to remove an electron from the metal surface $E_{\text{WF}} = 4.0$ eV. The key features that we followed while choosing our neutralizer are below:

- 1) The oxidation of the foil material could negatively impact the release of neutral atoms from the surface.
- 2) The foil material must have a low enough vapor pressure at the operating temperature to be compatible with our UHV environment and to avoid affecting the vacuum background when heated.
- 3) The material should not be too brittle, as it needs to undergo tens to hundreds of thousands of cycles of positional changes while accumulating and releasing neutral atoms.

Prior to 2015, yttrium (Yt) was used with a work function 3.6 eV, which is significantly lower than that of Zr, and the ionization potential of Fr. It was replaced by Zr due to its brittleness, which made operation unreliable. Zr is more resistant to oxidation up to higher temperature (700 – 800)°C and is mechanically stronger. The mechanical setup to collect and rotate this heated foil to release neutral Fr atoms into the capture chamber is a pneumatically operated shaft connected to a spring loaded foil holder to hold 0.001" thick Zr foil. This shaft is controlled remotely by computers and is in sync with our trapping-transfer cycle of atoms. The pneumatic actuator is expected to last for 0.7 million of rotations before requiring replacement. Fig. 2.4 represents the release fraction of ^{211}Fr taken during a test

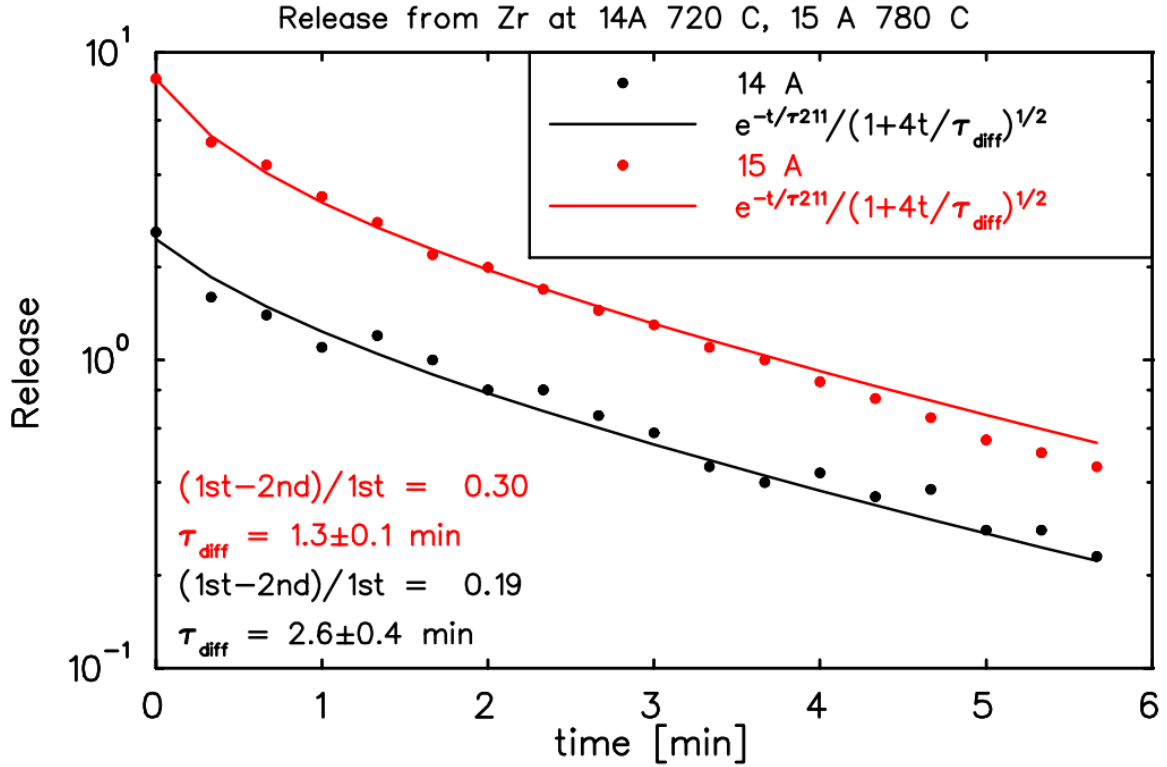


Figure 2.4: Release of neutral Fr from the Zr foil as a function of time for two different foil temperatures. The y-axis scale is arbitrary units and x-axis shows the time in minutes. Graph is provided by our collaborator Alexandre Gorelov.

run done for yield measurements in 2016 by our group. Fr was implanted in the foil, the Fr beam was turned off and then the number of atoms was measured with each heating of the foil. The diffusion lifetimes can be compared for 14 A, 720°C and 15 A, 780°C. For Fr beamtime, the Zr foil was heated for 2 s for every 28 s. The long tail of diffusion is not shown in this figure. The release of Fr from the Zr foil is not a pure exponential and is described by

$$N(t) = \frac{N_0}{\sqrt{1 + \frac{4Dt}{d^2}} e^{-t/\tau}},$$

[40] where D is the diffusion constant for Fr in the Zr foil, d is the implantation depth of Fr. The release of Fr from the foil depends on the temperature of the foil (or the current flowing through the foil), as shown in Fig. 2.5. For the work presented in this thesis, the neutralizer is pulsed at 12.0–12.3 A for 2 s every 28 s, reaching temperatures up to 830°C. The release

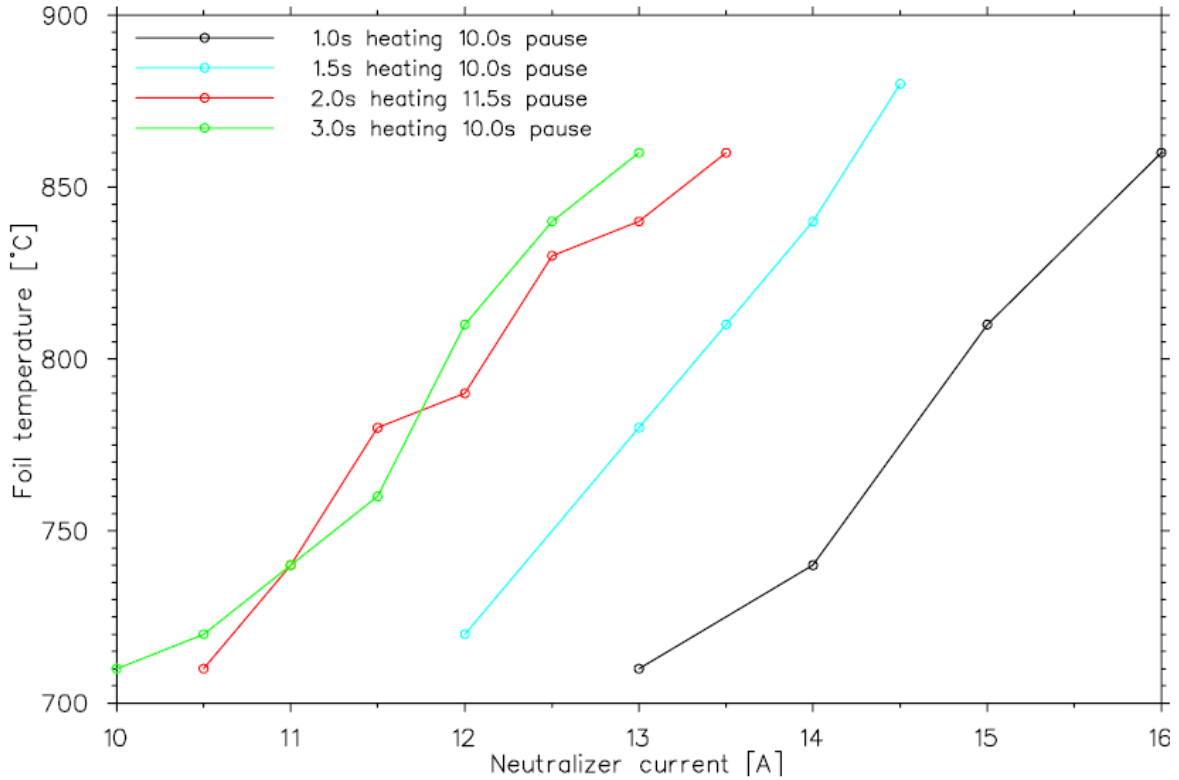


Figure 2.5: Temperature of zirconium foil vs current. Graph is provided by our collaborator Alexandre Gorelov.

efficiency of neutral ^{211}Fr from Zr increases with higher temperature and current, as shown in Fig. 2.4. However, we closely monitor the foil temperature because if it exceeds 850°C , the foil undergoes a phase transition [41] that fundamentally alters its material properties. In this regime, the foil develops a porous, cheese-cloth-like structure, which prevents efficient interaction with the surface and thus eliminates the neutralization of francium ions. A TTL signal from the computer triggers a LeCroy gate generator (LRS 222), which controls the current sent to the foil from a Kepco ATE 25-20 power supply. This dual gate-delay generator prevents erroneous re-triggering of the foil by enforcing the minimum delay set by the LeCroy. For additional protection of the foil, we have added a 4 A slow-blow fuse to the neutralizer supply circuit, whose time-current characteristic allows it to tolerate currents of approximately 12 A for several seconds before blowing; this permits ~ 2 s heating pulses at 12 A while ensuring that sustained currents of about 12.5 A lasting longer than approximately

5 s will open the circuit and prevent damage to the foil. Additionally, the heating time of the neutralizer must be carefully managed relative to the total cycle time. If the cycle time is too short (less than 10 s), the repeated heating pulses prevent sufficient thermal relaxation between cycles, causing the baseline temperature of the neutralizer assembly to rise and potentially exceed 850°C. Conversely, if the cycle time is too long (greater than 30 s), the diffusion efficiency becomes suboptimal, as the foil does not reach or maintain the temperature–time conditions (approximately 830°C for 2 s) required for efficient release.

2.4 Trapping apparatus

To achieve efficient atom production for precision measurements of APV, the trapping apparatus is divided into two sequential chambers. The first chamber, located adjacent to the neutralizer assembly, is optimized for trapping francium atoms with large magnetic field gradients and high laser intensities, and has a compact geometry. These features collectively maximize the loading rates but make the chamber unsuitable for high-precision measurements requiring a low-background environment and well-controlled field geometries.

To overcome these limitations, the trapped atoms are transferred from the first trap to a second trap located approximately 1 m away by optically pushing the atoms. The second trap provides increased optical access, a low-background environment, and houses a Fabry–Pérot cavity inside the trapping chamber. Once the atoms are transferred to the lower chamber, laser spectroscopy is performed for precision measurements. The details of these chambers are described in the following section.

2.4.1 Capture trap

The capture trap (CT) is mounted directly onto the radioactive beamline delivering Fr^+ ions. A critical component for capturing atoms in the MOT within the CT is the neutralizer foil, which is positioned just in front of the Faraday cup, as shown in Fig. 2.6. The Zr foil receives the Fr^+ ions, is subsequently heated to release neutral atoms into the glass cell. The

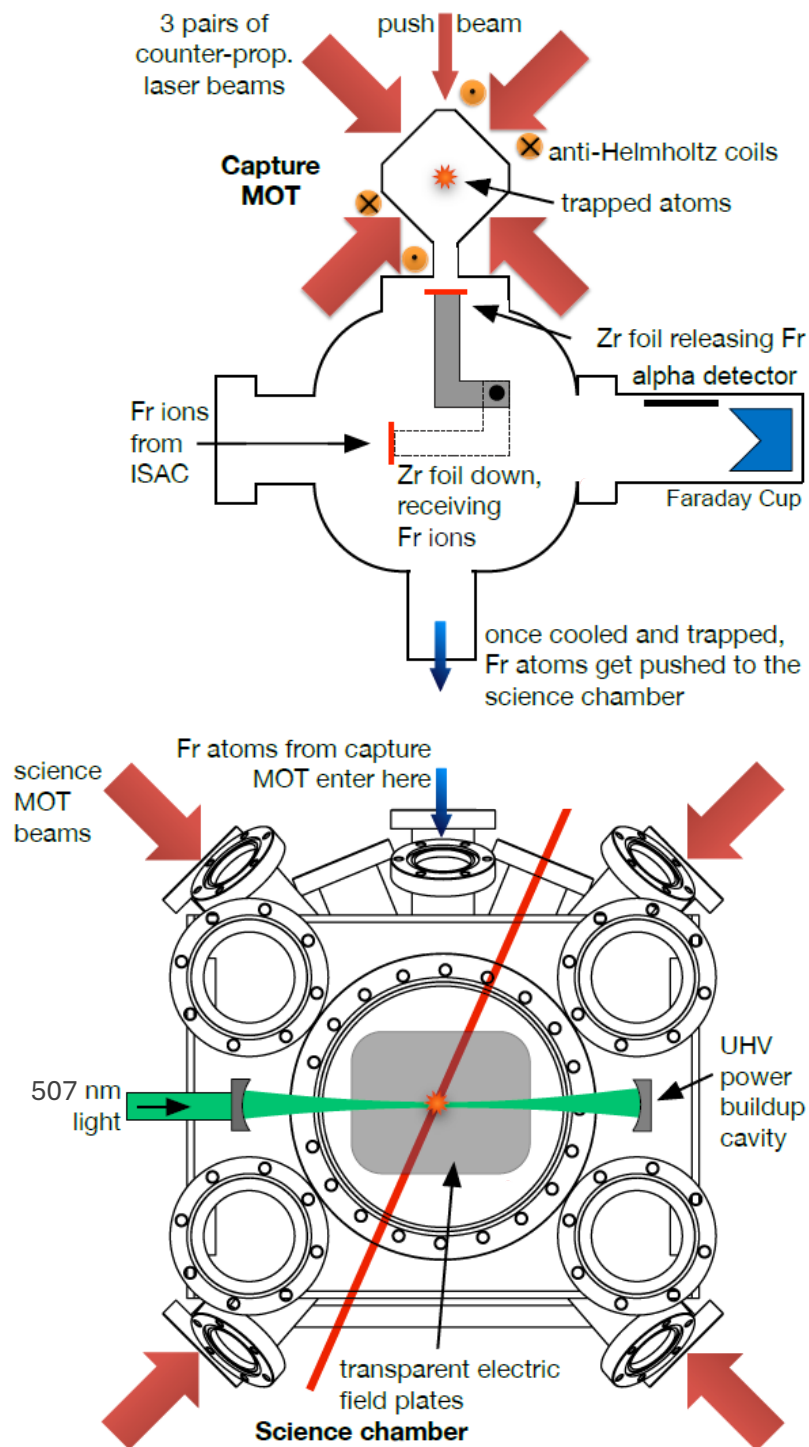


Figure 2.6: Trapping apparatus for atoms in a magneto-optical trap (MOT). The upper half of the figure shows the capture trap, while the lower half shows the science chamber. The figure includes the MOT beams, the neutralizer, magnetic coils, the push beam, the power build-up cavity, and the electric field plates. The trapped atoms are pushed from the capture trap to the science chamber, where 7S–8S spectroscopy is performed.

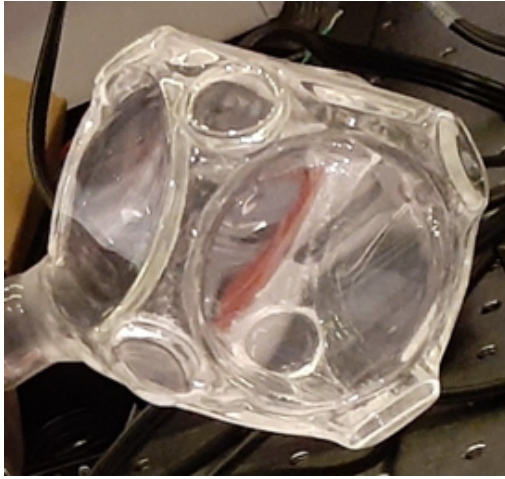
Fr atoms can collide with the walls of the glass cell; such collisions may depolarize the atoms or lead to their adhesion to the surface, resulting in losses from the MOT. To minimize these effects, a well-established approach is to coat the Pyrex walls with an organic polymer. This coating promotes elastic-type collisions and helps maintain the atom population in the vapour phase without significant loss [42]. The capture trap consists of a glass cell whose inner walls are coated with a silane based dry film coating to avoid sticking of Fr atoms to walls. This coating also assists in re-thermalizing the atoms to a Maxwell–Boltzmann velocity distribution. For about 92% of the total trapping-transfer cycle, the Fr beam is directed onto the neutralizer foil; during the remaining time, it is dumped onto the Faraday cup.

The glass cell (in Fig. 2.7(a)) is protected with an exo-skeleton cage to which all MOT beam optics are mounted. The quadrupole field for the MOT is generated with water cooled coils to provide an axial magnetic field gradient of ~ 10 G/cm. There are bias coils to compensate ambient magnetic field such as the Earth’s magnetic field, or from nearby sources such as ion pumps, and two super-conducting magnets operated by the TITAN facility approximately 3 meters above the Fr traps. The MOT beams are 2 inches in diameter ($1/e^2$) to fill the whole glass cell with slowing/trapping light for maximum trapping efficiency. The linearly polarized incoming beams pass through quarter-wave plates to become circularly polarized before going into the glass cell, then are retro-reflected by mirrors to form the six-beam configuration required for 3D trapping of Fr. The retro-reflected MOT beams, whose handedness is flipped upon reflection, pass through the same quarter-wave plates so that the MOT light acquires the opposite helicity and thus maintains the required circular polarization. The trap and repump beams are delivered from the optics benches to the CT via optical fibers. They are first combined (mixing the trap and repump light) and then split into three MOT beams using an Evanescence Optics Inc. 2×3 fiber splitter/combiner. The power is distributed equally as $\sim 33\%$ (trap) and as $\sim 67\% + 16\% + 16\%$ (repump) into each MOT arm at 718 nm using polarization-maintaining fibers. Two CCD cameras capture MOT images of the atom cloud

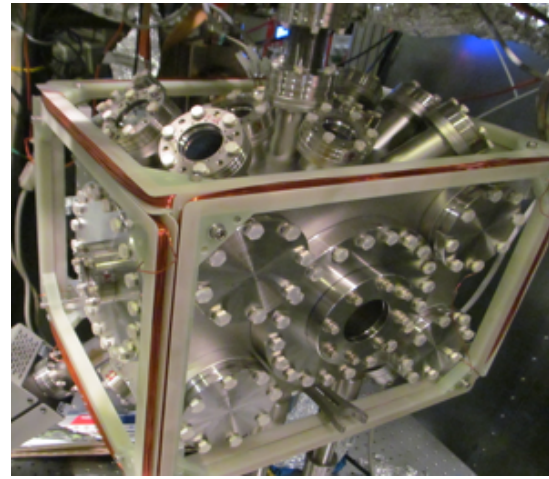
in the CT to estimate the atom number and the decay lifetime of the MOT. The bias coils also allow us to move the MOT within the CT to better optimize the transfer of atoms from the capture trap to a second trap in a different chamber, which will be described below. The chambers are connected by welded bellows that suppress vibrations originating from the radioactive beamline, which is mechanically coupled to the ISAC beamline and therefore subject to its associated vibrations. A pneumatic valve located between the capture trap (CT) and the lower chamber acts as a mechanical shutter and is used to isolate the vacuum between the two chambers whenever either chamber needs to be opened for repairs or for the installation of new components. The available open space around the CT is very limited, leaving insufficient room to perform laser spectroscopy. In addition, the background signal is elevated due to radiation from the nearby radioactive beamline. Because of the tight spatial constraints inside the CT, it is not possible to install electric field plates or a power build-up cavity, since the CT is dedicated primarily to achieving high trapping efficiency. Once the atoms get trapped in the CT, the trapped atoms get pushed vertically about 70 cm with push beam to the lower chamber, referred to as ‘Science Chamber’, SC. The push beam is tuned to resonance with the $7S_{1/2}(F = 5) \rightarrow 7P_{3/2}(F' = 6)$ atomic transition.

2.4.2 Science Chamber

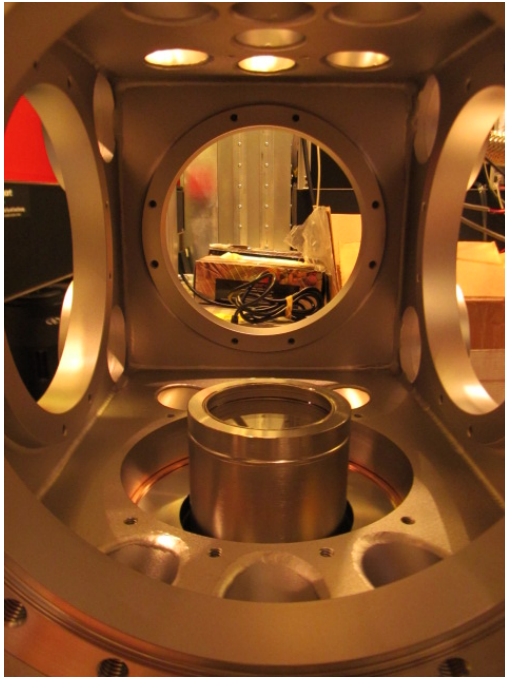
The atoms are first pushed from the Capture Chamber to the Science Chamber (SC), where they are immediately recaptured in a second MOT. The MOT beams in the SC have a 1 inch ($1/e^2$) diameter and are delivered in the same configuration as in the CT. The 718 nm trap and repump light is routed from the laser benches to the SC via optical fibers and distributed using an Evanescent Optics Inc. 2×3 fiber splitter/combiner, with the same power distribution as in the CT. The SC is a non-magnetic, stainless steel, chamber built by Pfeiffer Vacuum [34]. This chamber is equipped with multiple viewports for MOT beam delivery and retro-reflection, as well as for auxiliary beams, the detection system, ion pumps, and alkali-metal dispensers (described later). It also features recessed flanges to hold quadrupole



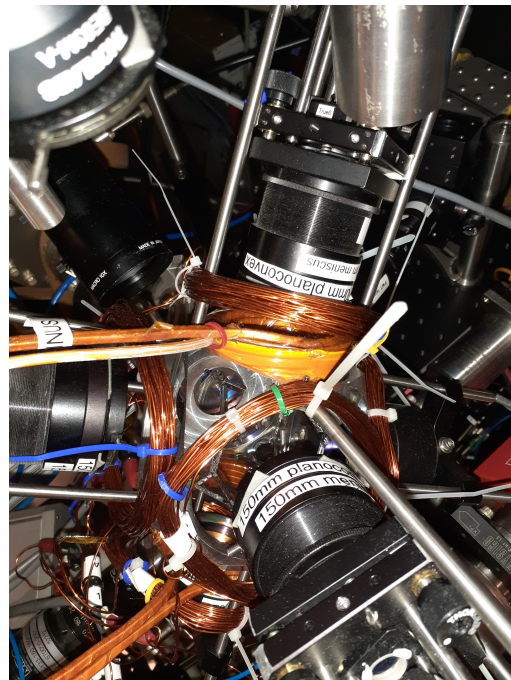
(a) Capture trap: Glass cell to trap atoms from neutralizer. In the real cell, the small circular side faces present in the mock-up cell (as shown in the image) are replaced with triangular-shaped faces.



(b) Science Chamber to do 7S–8S spectroscopy. In the current image, the front (visible to the reader) and back-facing viewports have been replaced with a flange near which the quadrupole coil is positioned.



(c) Inside view of the science chamber. The recessed viewport can be seen clearly in the center of science chamber.



(d) Capture trap with exo-skeleton on glass cell with all the optics mounted.

Figure 2.7: Images of the trapping apparatus: the Capture trap and the Science chamber.

coils in an anti-Helmholtz configuration, generating the magnetic field gradient required for the MOT. These water-cooled coils typically operate at 60 A, and the recessed design helps reduce eddy currents in the copper gaskets. An additional recessed flange at the bottom of the SC (as shown in 2.7(b and c)) supports the detection setup, including fluorescence collection optics, an aperture, and other components, positioning the collection optics closer to the MOT to enhance light collection. The SC also contains three pairs of bias/shim coils to compensate for ambient magnetic fields. Since we perform $7S - 8S$ spectroscopy in the SC, one of the significant contributions to the total transition rate is the Stark-induced transition, which arises from the mixing of hyperfine levels when an external electric field is applied. The SC accommodates two transparent electric field plates made of BK7 glass and coated with indium-tin oxide (ITO). The plates measure $4.75'' \times 4.75''$ and are separated by spacers with a distance of $1.125'' \pm 0.0004''$ [43]. The field plates are transparent and are spaced out from each other to facilitate the MOT beams. The trapped atoms in the SC are observed using two CCD cameras, which view the MOT from different planes. One of the cameras shares the line of sight with the detection system, facilitating easy alignment of the detector (discussed later in this chapter). The camera data is used to estimate the number of atoms trapped in the MOT, which is then treated as one of the observables for evaluating systematic uncertainties. Furthermore, the SC also houses a Fabry-Pérot cavity, referred to as the power build-up cavity, which is discussed later in the thesis.

2.4.3 Vacuum

The pressure in the CT and SC is maintained at approximately 10^{-10} mbar. The vacuum in the CT is maintained using a hybrid vacuum pump, SAES NEX Torr D100-5, which combines a non-evaporable getter (NEG) and an ion pump to remove active gases such as O_2 , H_2 , and CO. The SC is equipped with a Pfeiffer turbo pump, a Gamma TITAN ion pump (Ion Pump - 45S-DIX-4D-SC-110-N), and a SAES NEG pump (CapaciTorr D 400-2). To evacuate the SC, the pressure is first reduced to $10^{-6} - 10^{-7}$ mbar using the turbo pump. After the

all-metal valve is closed, the ion pump is activated to further reduce the pressure; the ion pump effectively pumps the residual gases such as H_2 and H_2O . Once the pressure reaches approximately 10^{-8} mbar, typically after a day of turning on the ion-pump operation, the NEG pump is turned on. The NEG pump is first pre-conditioned for 30 minutes followed by activation of the getter material for one hour. The SC is also equipped with an SRS RGA100 residual gas analyzer, which monitors the gas composition during pump-down after venting the chamber. This vacuum system collectively enables atom lifetimes of 10-20 s in the MOT. Given the limited annual availability of Fr (5–6 days per year), rubidium dispensers are installed in both the CT and the SC to facilitate apparatus construction and testing.

2.4.4 Trap and repump laser

To work with Rb and Fr, we work with 718 nm for Fr and with 780 nm for Rb. The suitable wavelength is achieved with continuous-wave, tunable Ti:Sapphire lasers. For ^{211}Fr , the ground state $7S$ has two hyperfine ground states $F_g = 4, 5$, while the excited state $7P_{3/2}$ has hyperfine levels ($F'_e = 3, 4, 5, 6$), as shown in Fig. 2.8. The atoms are trapped by using a cycling transition on $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F'_e = 6)$, consistent with the selection rule, $\Delta F = \pm 1, 0$. This laser is referred to as the ‘trap laser’. Due to the finite linewidth of the trap laser and the number of atoms, there is a finite probability that the atoms will be driven to the nearby transition $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F'_e = 5)$. Atoms excited to $7P_{3/2}(F'_e = 5)$ can decay to either $7S_{1/2}(F_g = 4)$ or $7S_{1/2}(F_g = 5)$. Once atoms decay to the lower hyperfine ground level $F_g = 4$, the trap laser is no longer resonant. To prevent atoms from falling out of the cycling transition, another laser called ‘repump laser’, drives the $7S_{1/2}(F_g = 4) \rightarrow 7P_{3/2}(F'_e = 4)$ transition. This return atoms to the upper hyperfine ground level, ensuring that they remain in a closed cycling transition.

The trap laser is derived from a Ti:Sapphire Solstis laser from M Squared, pumped with 8.5 W of 532 nm light from a Coherent Verdi G pump. This configuration produces approximately 850 mW of output power at 718 nm. The laser is frequency stabilized by locking it to a

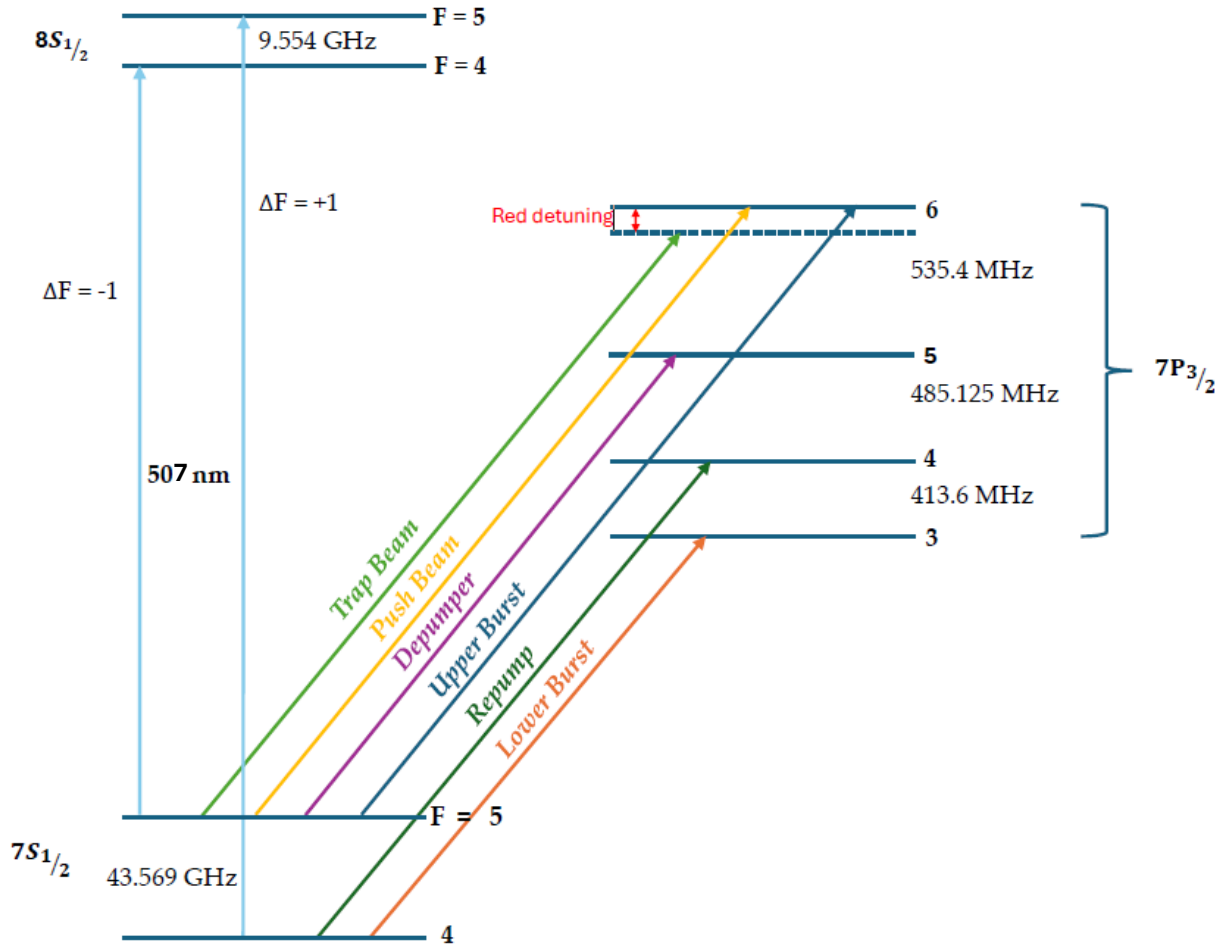
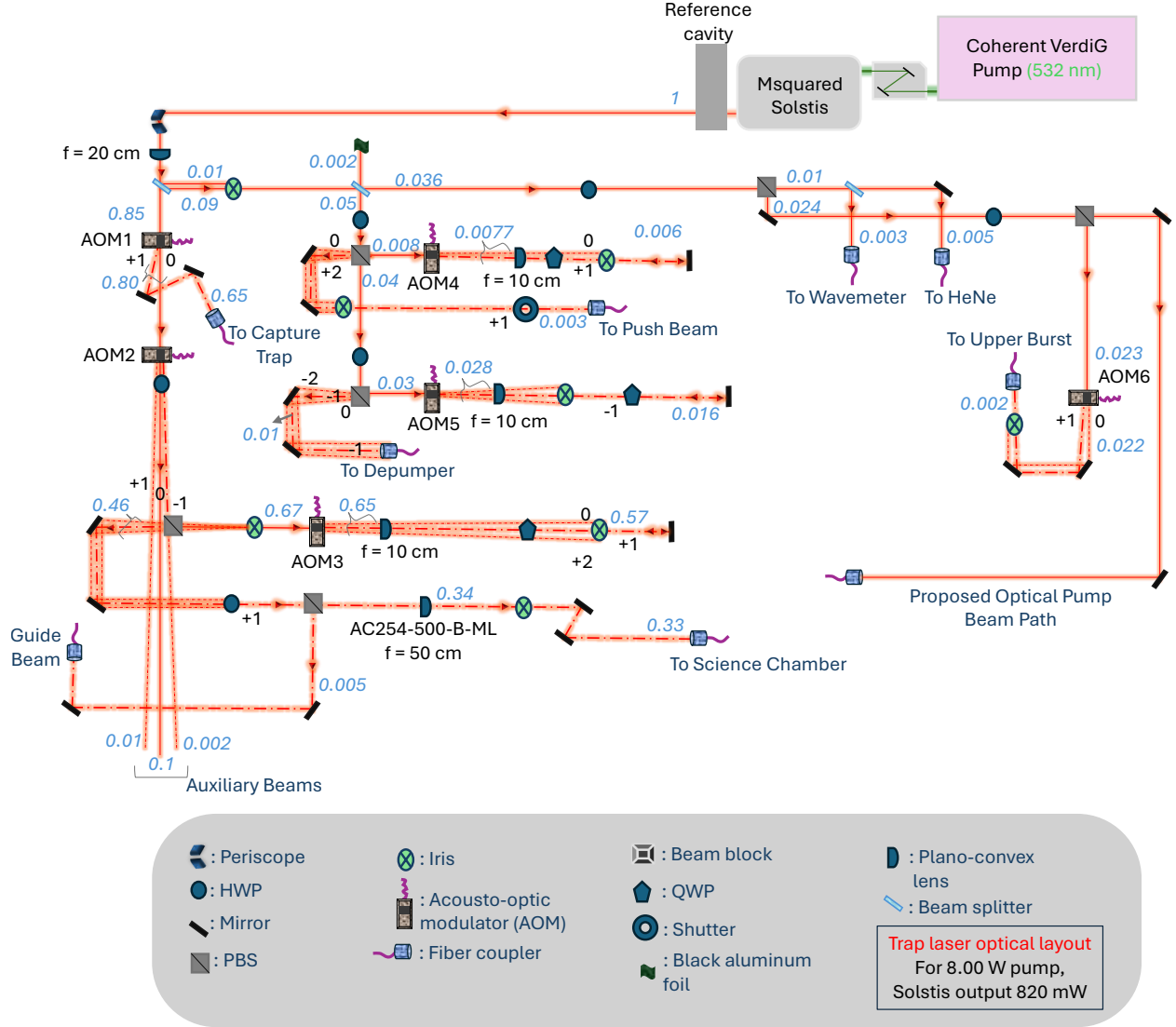


Figure 2.8: ^{211}Fr structure showing hyperfine splitting and all driven transitions .

reference cavity from M Squared. The repump laser is also a Solstis from M Squared. It is pumped with ~ 6.5 W at 532 nm light from a Sprout G laser, producing about 650 mW output at 718 nm. We derive several additional probe beams from the trap and repump laser, as shown in Figs. 2.9 and 2.10. Another key laser in the FTF is the ‘spectroscopy laser/blue laser’, which is used to drive the $7S \rightarrow 8S$ transition. This transition requires a powerful laser source at 507 nm, but sufficient laser intensity is not directly available at this wavelength. The solution is to frequency-double a Ti:Sapphire laser. For this purpose, a Ti:Sapphire Solstis laser is pumped with 16 W at 532 nm light from a Coherent Verdi V-18, producing approximately 900 mW of output power at 1013 nm. The spectroscopy laser at 1013 nm is then frequency-doubled using an external cavity doubler (Solstis ECD-X) to obtain 507 nm. We measure the wavelengths of these beams with an Ångström WS U10 wavemeter with an absolute accuracy of 10 MHz, and calibrate our wavemeter periodically to the $5S_{1/2} \rightarrow 5P_{3/2}$ transition of rubidium saturation spectroscopy.

2.4.5 Laser beam directory

The $7S \rightarrow 8S$ transition is driven between both hyperfine components, following $\Delta F = \pm 1, 0$, which means that the trapped atoms should be pushed into the lower hyperfine level $7S_{1/2}(F_g = 4)$. Atoms are initially trapped using the cycling transition $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F_e = 6)$. To enable excitation of the $7S \rightarrow 8S$ from the lower hyperfine ground level, the trap operates on the upper hyperfine level while a depumping resonant with $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F_e = 5)$ transfers population to $F_g = 4$, as shown in Fig. 2.8. Most of the beams addressing the $7S_{1/2}$ and $7P_{3/2}$ transition are derived from the trap laser. We shine a resonant laser beam to $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F_e = 6)$ to push the atoms from CT to SC. The trapping laser beam is switched between the CT and SC using an AOM, maintaining the same terminal frequency. There is a mechanical shutter in front of push fiber coupler on optical bench with trap laser to protect any unintentional leak of push beam. The trap laser beam also goes to wavemeter which is used as frequency readout of trap laser, to HeNe laser to



frequency stabilize the trap laser, and to upper burst beam on $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F'_e = 6)$ (details in section 2.6.2). The Fig. 2.9 shows the optical bench layout of trap laser beam with details of pick up of the all beams. The other important beam is the guide/clear beam which is on the same red-detuned frequency as the trap frequency in SC. This beam is used to push the trapped atoms out of the interaction region at the end of the data taking cycle so that we can estimate the background in the data taking window. The frequency shifts for the trap and related beams are implemented using acousto-optic modulators (AOMs). Tables 2.1 and 2.2 summarize the AOM configurations and are essential to understand how the required beams are turned on and off at different times. Most AOMs are switched off by either detuning the RF frequency from the center or reducing the drive amplitude. A few use a double-pass setup to achieve the large frequency shifts (in MHz) needed for transitions involving the $7P_{3/2}$ state, as shown in Fig. 2.9.

From repump laser, a laser beam goes to HeNe for frequency readout, to HeNe to stabilize the frequency of repump laser shown in Fig. 2.10. We also drive a lower burst beam on $7S_{1/2}(F_g = 4) \rightarrow 7P_{3/2}(F_e = 3)$ (details in section 2.6.2). The trap and repump lasers are frequency stabilized using a digital feedback system that locks them to a HeNe reference via a Fabry–Pérot cavity, providing long-term stability. The relative peak positions of the target lasers and the HeNe are monitored and corrected in real time, achieving a stability of about ± 2 MHz/h. Laser frequencies are monitored with a wavemeter through a fiber-optic switch, calibrated using rubidium spectroscopy, and periodically verified to correct for drift during beamtimes.

For blue laser, 1013 nm from M2 Solstis Ti:Sapp laser is branched off with a beam-splitter to wavemeter and the other branch goes to a fiber EOM1 (NIR-MPX-LN-20-P-P-FA-FA) before going to an ultra-low expansion (ULE) cavity. This ULE is an external reference Fabry–Pérot cavity that provides the frequency stability upto 120 kHz. This ULE is temperature controlled with a Thorlabs TEC controller. The details of this setup are in [43]. As shown in Fig. 2.11, the ULE beam path has a resonant EOM2 (New focus 4002) that produces

Table 2.1: Detailed description of AOMs used in the trap laser setup, including their types, frequency shifts, functions, and operation.

AOM	AOM Type	Frequency Shift	Function	Operation
AOM1	Switching AOM (capture trap and science chamber)	+1 (90 MHz)	Switches the trap beam between the capture trap and the science chamber.	On-state RF frequency: 90 MHz, -8 dBm, set on the controller. Off-state RF frequency: terminated on the RF switch.
AOM2	Chopping AOM (science chamber)	-1 (90 MHz)	Chops the trap beam in the science chamber.	Switches RF frequency between 90 MHz (-3.95 dBm) and 160 MHz (-45 dBm). Controlled by the controller.
AOM3	Double-pass AOM on science chamber	(+1,+1) (90 MHz)	Brings trap light into resonance with the trap frequency.	Constant amplitude and frequency provided by the controller. Always on.
AOM4	Double-pass AOM to push atoms (Rb and Fr)	(+1,+1) (55 MHz)	Brings light on-resonance to push atoms from the capture trap to the science chamber.	Constant amplitude and frequency from the controller. Push beam always on. TTL-controlled mechanical shutter on the push beam.
AOM5	Double-pass AOM to depumper (science chamber)	(-1,-1) (213 MHz)	Sends light to depump atoms to the lower hyperfine ground state.	Constant frequency from the controller. Amplitude pulse provided by a delay generator when depump is on.
AOM6	Single-pass AOM for upper burst	+1 (95 MHz)	Brings light on-resonance for the upper burst beam.	Switches RF frequency between two values. RF generator sends signal to the RF switch. Constant amplitude.

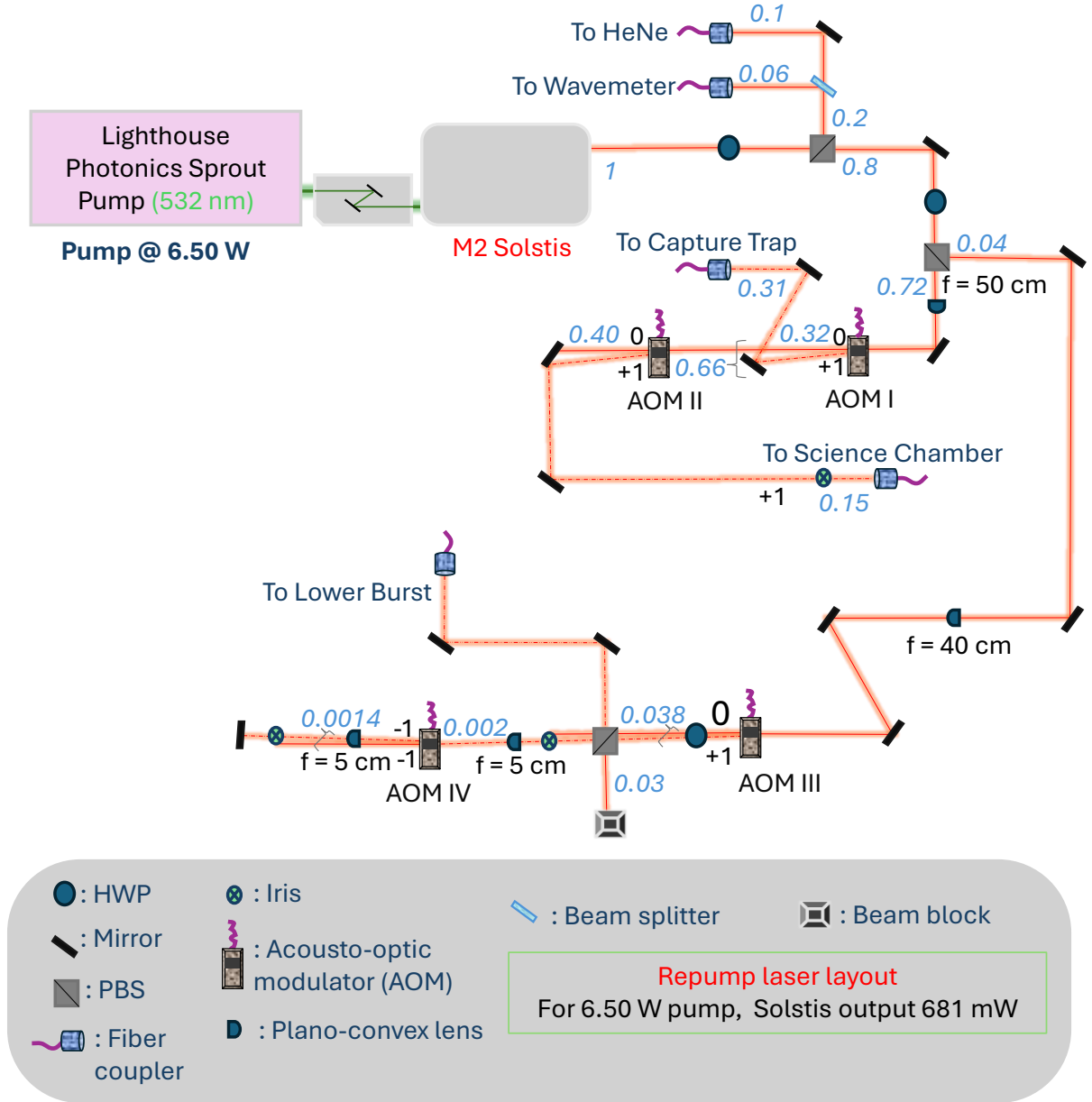


Figure 2.10: The repump laser beam at 718 nm or 780 nm (the output power mentioned in the figure) is branched out and sent to various destinations. The laser output for 718 nm for 2024 beamtime was 330 mW for 9.0 W pump. Along the way, the beams pass through several AOMs (see Table 2.2 for details), beam splitters, irises, and other optics. These elements help isolate and couple the desired sideband with right frequency into fiber couplers, enabling delivery of the beams from the optics table to the CT and the SC.

Table 2.2: Description of AOMs on the repump laser path, including frequency shifts, functions, and operation.

AOM	AOM Type	Frequency Shift	Function	Operation Description
AOM1	Switching AOM (capture trap and science chamber)	+1 (90 MHz)	Switches the beam between the capture trap and the science chamber.	On-state RF frequency: 90 MHz, -8 dBm, set on the controller. Off-state RF frequency: terminated on the RF switch.
AOM2	Chopping AOM (science chamber)	+1 (90 MHz)	Sends light to the science chamber.	Switches RF frequency between 90 MHz (-3.95 dBm) and 160 MHz (-45 dBm). Controlled by the controller.
AOM3	Single pass to burst beam	(+1) (90 MHz)	Brings burst light into resonance with the repump frequency.	Constant amplitude. Switches RF frequency between two values via delay generator.
AOM4	Single pass for Rb and double pass for Fr	($-1, -1$) (153 MHz)	Brings light on-resonance to the lower burst cycling transition.	Constant amplitude and constant frequency, both provided by the controller.

sidebands at 15.507MHz to produce ± 1 sidebands to lock the laser to ULE cavity. This light that gets reflected back from first mirror of ULE cavity goes as the reflected signal into photodiode (Thorlabs, PDA10A2). The signal from photodiode goes to mixer that mixes the signal into two signals where one comes from reference oscillator and other signal comes from the photodiode. This ULE cavity is locked to the TEM₀₀ mode created by a fiber EOM1 (iXblue). The output of the PID lock goes to the fast piezo of M2 Solstis laser as an active feedback. The RF frequency for the fiber EOM1 is controlled by a frequency synthesizer (FSW-0200). We set the frequency steps remotely as a list. The synthesizer connects to the computer via USB and syncs with data taking through a trigger from the SRS DS345 frequency generator. The phase matrix goes through 81 steps in both ‘up’ and ‘down’ modes, totalling 160 steps. Each frequency step lasts 50 ms, with a step size of 500 kHz and a sideband power of about 15 dBm. This 1013 nm is doubled with ECD to 507 nm and it goes through an AOM1 shown in Fig. 2.11 producing a first order sideband at 110 MHz which gets coupled into a fiber coupler to deliver the blue light from blue light optical bench to the SC.

Need of power build up cavity (PBC): The transition rate can be increased by increasing the number of trapped atoms, enhancing the laser intensity, or improving the fluorescence detection efficiency. The limited Fr⁺ yield at ISAC (TRIUMF) makes an increase in the number of atoms unfeasible. A practical alternative is to use a Fabry–Pérot cavity, where the 507 nm beam undergoes multiple reflections, effectively increasing the intensity seen by the atoms. A Fabry–Pérot resonator was built and installed in the vacuum chamber for this purpose that would increase the intensity of blue light within the cavity to tens of kW/cm². The details of the PBC are in the next chapter.

2.4.6 Controller

We used the NI PCI-e 6713 card to control the AOMs, the neutralizer’s motion, and the transfer of trapped atoms between the chambers, along with other timed actions. These

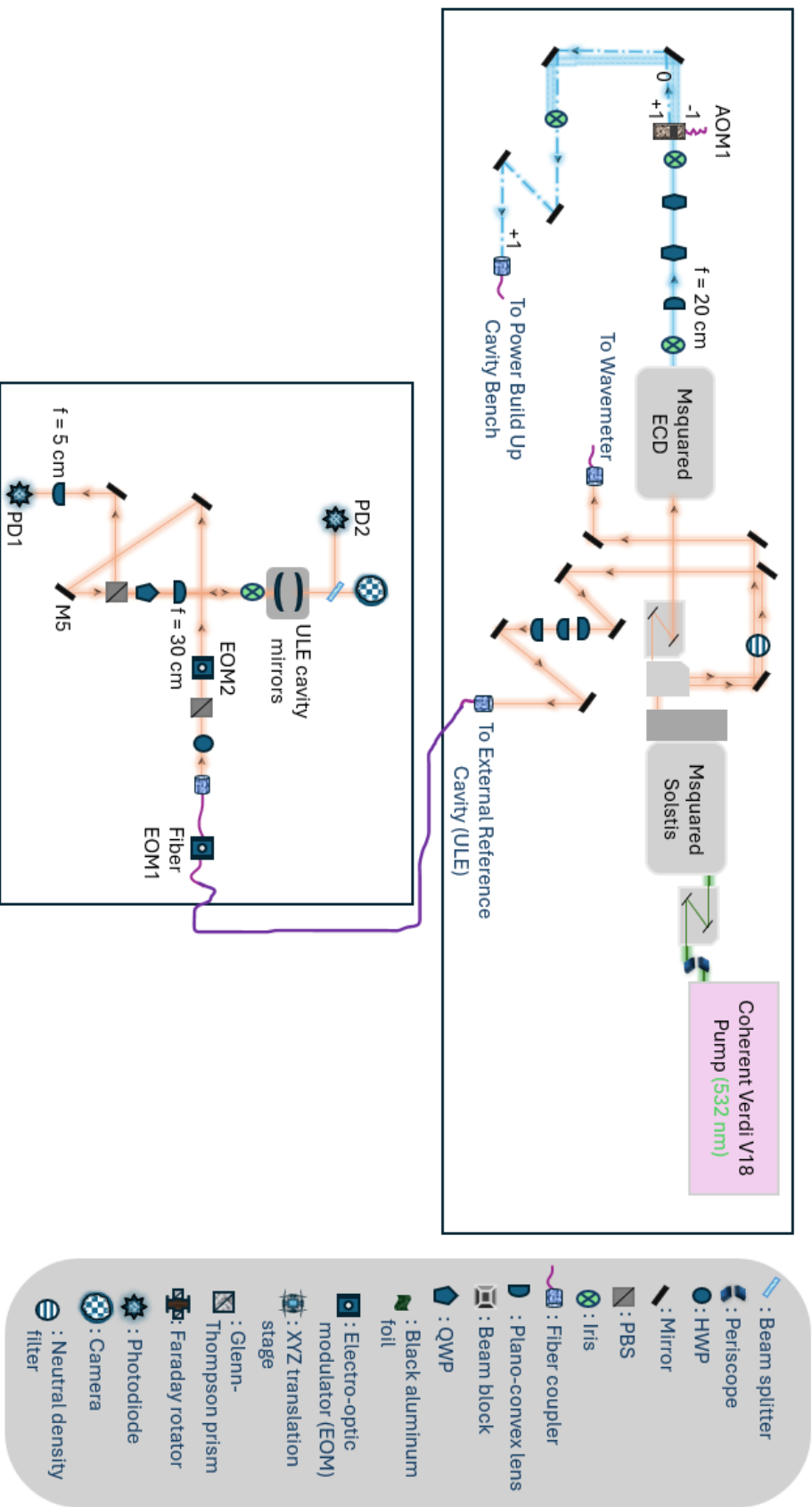


Figure 2.11: Coherent 18 W pump generates 1013 nm, goes to the wavemeter and to ULE which serves as an external reference cavity for the laser. 1013 nm is doubled with ECD to 507 nm which goes to SC with high power optical fiber.

cards send voltage signals to the AOMs and shutters through analog and digital buffers. The Pulse Blaster card generates TTL signals to trigger the shutters and AOMs. It has 24 programmable TTL pulse channels, so let us set the timing/pulse width for each pulse. Our entire process, including loading, trapping, transferring, and data collection, is based on the precise timing of the Pulse Blaster card. The transfer AOM for moving atoms and the chopping AOMs on the trap and repump beams are controlled by direct digital synthesized signal generators (Model 409B), which produce synchronized radio frequency sine waves for the AOMs.

The data taking part has a few internal cycled loops that include repetition of many similar sub-cycles in data taking period. A ‘primary trigger’ signal from the controller/pulse blaster is sent to trigger the SRS 365 function generator that outputs an arbitrary waveform (a waveform as in rectangular pulse train) with 160 pulses. This SRS DS345 function generator triggers another SRS DS345 function generator which also outputs an arbitrary waveform with 39 counts. The trigger from this SRS DS345 function generator triggers two DG535 delay generators. These delay generators control the trap, bin trigger, repump, and blue beam on/off with pulse amplitudes and their widths. One of the delay generator further triggers the other DG535 delay generator that triggers the burst and depump beams on/off.

The total timing of a single cycle that includes loading atoms on a neutralizer foil, trapping, transferring, taking data is ~ 28 s.

2.5 Detection of fluorescent photons

The detection system is placed under the bottom in the recessed viewport of the SC. It includes light-collection optics consisting of two 2-inch achromatic doublet lenses, each with a focal length of 10 cm. The lenses are positioned to form an f – $2f$ – f imaging system, providing a straightforward relay between the MOT and the detection plane. The fluorescent light is gently focused at the iris such that the aperture can be closed down to reduce the background from scattered light. The collimated light passes first through a fixed polarizer to suppress

the push-beam light, followed by a half-wave plate and a polarizing beam splitter, which together allow selection of the light directed either to the camera for observing the MOT position in the SC or to the detector for photon counting. The splitting of light with PBS avoids the overloading of detector from the linearly polarized push beam. A stack of filters consisting of a color glass filter to cut out the blue light at 497 nm and 507 nm, and a FBH720-10↓ bandpass filter at 718 nm. The collimated light is then reflected by a mirror oriented at 45° and subsequently passes through a set of 1" diameter convex lenses with a focal length of 3 cm, which focus the fluorescence light from the MOT onto the silicon photomultiplier (SiPM, Hamamatsu C14456-3050GD). It is a photon counting unit with a comparatively low dark count rate, 60,000 counts per second, with an integrated in-built temperature control unit. A finite fraction (20–30% of the total) of the light from the PBS is directed to the camera branch to visually locate the MOT, which is later used to estimate the number of atoms in the MOT. The same branch has another PBS with half waveplate so that a second photon counting device can be used. The detector assembly is mounted on a three-dimensional micrometer stage to image the MOT cloud onto the detector. The detection efficiency is approximately 1/4000 per photon, limited by the quantum efficiency of the SiPM, the PBS, filters, the solid angle of the collection optics, and the quantum efficiency of the detector (30%). The signal from SiPM goes into SCB-68A, a shielded I/O connector block for easy signal connection to a counting channel on NI PCI-7811R with FPGA module.

2.6 Bursting of atoms in detection

The detection efficiency of the current system is approximately 1/4000, resulting in an insufficient number of detected counts per signal. In principle, the efficiency could be improved by using a detector with higher quantum efficiency, and by increasing the collected solid angle. However, these options are largely impractical. Commercially available detectors offer a maximum quantum efficiency of only about 30%. Increasing the solid angle would require placing the fluorescent light collection optics closer to the MOT, which is not feasible

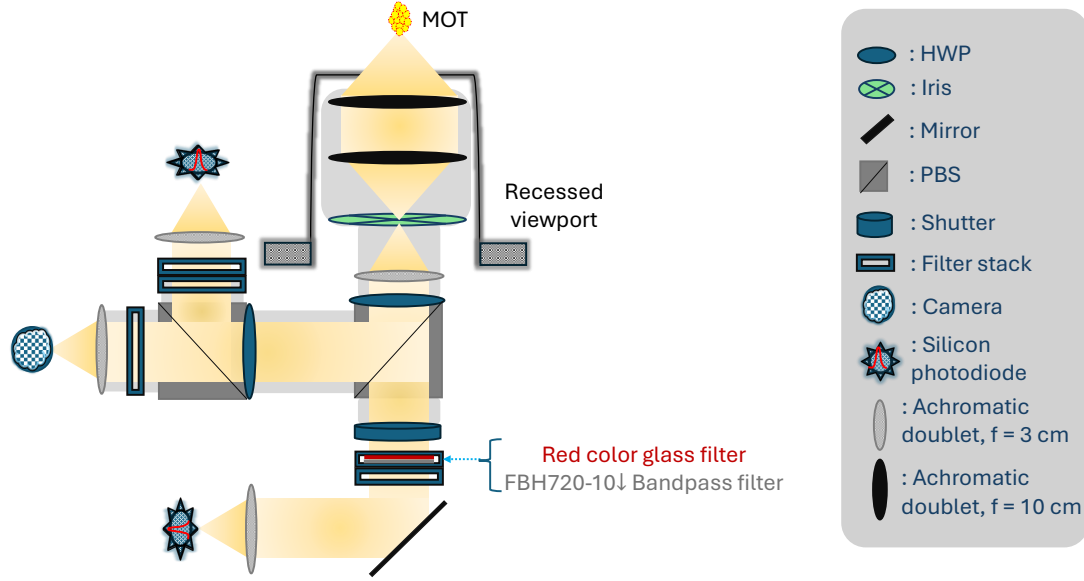


Figure 2.12: Current fluorescence detection system with two SiPMs and a CCD camera.

because the lenses would obstruct the MOT beams. However, the Boulder group [35] used cycling transitions to optically amplify the photons carrying the signature of atoms undergoing the $S \rightarrow S$ transition. The resulting burst of photons compensated for the low detection efficiency. We have implemented a similar scheme in francium.

2.6.1 Steps for the Burst technique

1. Operate the MOT to trap atoms in the science chamber. Trapping beams are applied for 1 ms.
2. Turn off the MOT light and use a repumper or depumper to optically pump all atoms into a single hyperfine ground level, leaving the other ground level empty.
3. Excite the atoms from the populated ground state via the $7S \rightarrow 8S$ transition using an intense 507 nm blue beam applied for 400 μs .
4. Allow the atoms to decay from the $8S$ state; a finite fraction decays through the $7P$ manifold into the initially depleted hyperfine ground state.

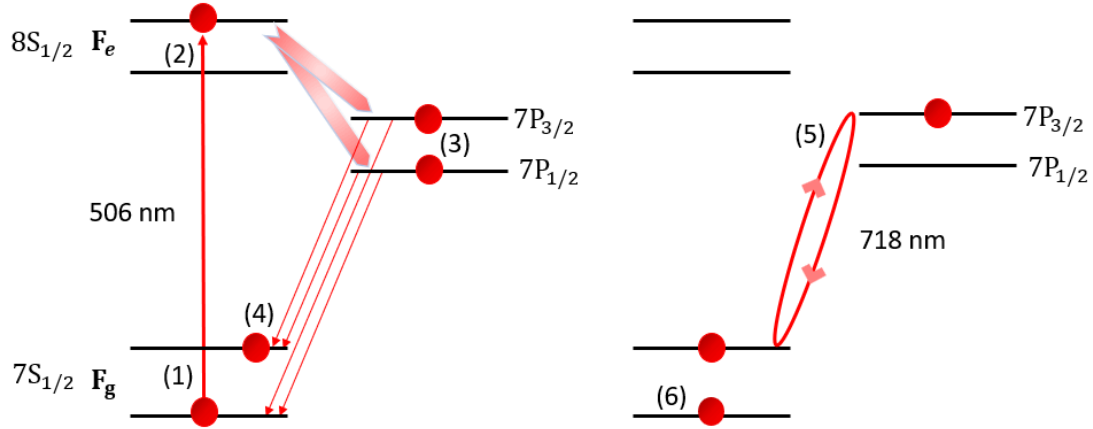


Figure 2.13: Experimental steps for 7S - 8S transition with burst cycle. Assume all atoms are pumped into the lower hyperfine ground state, and steps are in following order: (1) excite atoms to hyperfine level of 8S (2) the excited atoms decay to 7S via 7P (3) and a significant fraction of atoms ~ 50 percent from 7P state decay to the upper hyperfine level (4) and the probe laser cycles these decayed atoms on upper cycling transition (5). All atoms that were not excited to 8S state remains ‘dark’ in 7S lower hyperfine level (6). The hyperfine levels in this figure are not to scale.

5. Introduce a burst window of $400 \mu\text{s}$, overlapped with the data-taking timing window. During this period, all other trapping lasers and the blue beam are turned off.
6. Apply a cycling transition to repeatedly excite these atoms and generate a burst of fluorescence photons. For clarity, we explicitly define two cycling transitions whose names will be used in later sections:

- **Upper burst:** the cycling transition $7S_{1/2}(F_g = 5) \rightarrow 7P_{3/2}(F_e = 6)$.
- **Lower burst:** the cycling transition $7S_{1/2}(F_g = 4) \rightarrow 7P_{3/2}(F_e = 3)$.

Applying either the *upper burst* or *lower burst* transition drives the atoms repeatedly, producing a corresponding burst of fluorescence photons.

7. Detect the resulting fluorescence during the measurement window.
8. Turn the MOT beams back on to retrap the atoms and repeat the sequence.

2.6.2 How do we derive the burst beam?

Since the burst beam is on the D2 line, we derive the upper burst beam from the trapping laser. The burst beam on the lower cycling transition is derived from the repump laser on $7S_{1/2}(F_g = 4) \rightarrow 7P_{3/2}(F_e = 3)$. The burst beam is turned on/off with an acousto-optic modulator which acts as a shutter. The beam diameter of the burst beam is 5 mm with beam power at saturation intensity. We minimized the momentum transfer to the atoms during the bursting by using a retro-reflected beam. For the red-detuned burst transition, the beam power and duration are optimized to approach saturation without heating the trapped MOT atoms, balancing efficient excitation with minimal perturbation.

2.6.3 Measured decay time of cycling transition

The burst beam drives a cycling transition that repeatedly excites the atoms; however, as mentioned earlier, the transition is not fully closed because there is a finite probability that atoms leak into the other hyperfine ground level after being excited to $7P_{3/2}(F_e = 5, 4)$ for the upper and lower burst transitions. We developed this technique for Fr by measuring these decay times for the upper and lower cycling transition with SRS SR430 multichannel scalar (MCS) for the upper and the lower burst in ^{87}Rb . Since Fr has larger hyperfine splittings than Rb, the cycling transition decay time is expected to be longer than in the case of Rb. The details can be found in Chapter 6 of thesis [33]. The decay time $1/e$ of the upper and lower bursts for Rb is measured as 1.60 ms and 330 μs [33]. An estimate from a simulation says that $\sim 10^4$ photons could be emitted before the atoms get lost in the dark hyperfine level for Fr. This number of photons gives almost 100% efficiency for the detection system. Thus the burst technique improves the current detection system's efficiency by $\sim 10^3$ fold.

2.7 Summary

To summarize, the francium trapping facility receives Fr^+ ions from ISAC, which are accumulated on a thin Zr foil and heated to release neutral ^{211}Fr atoms into the capture trap. Zirconium is used instead of yttrium due to its mechanical stability, and ^{211}Fr is chosen for its low activity, providing a safer radiation environment. The capture trap is a glass cell with anti-stick coatings, where neutral atoms are cooled and trapped in a MOT using red-detuned light and a magnetic field gradient.

The MOT is then pushed into the Science Chamber and retrapped, enabling 7S–8S spectroscopy with 507 nm light under controlled electric and magnetic fields. Atoms excited to 8S decay through 7P, producing 718 nm photons for detection. The previous detection system had an efficiency of about 1/4000, which was overcome using the burst technique: atoms driven on a cycling transition on the D2 line repeatedly absorb and emit photons, amplifying the signal by ~ 4000 and achieving unit detection efficiency.

Both the Capture Trap and Science Chamber operate under ultra-high vacuum. The trap and repump lasers are frequency-stabilized using a HeNe-based digital feedback system, while the 507 nm light is stabilized with an external ultra-low expansion reference cavity. A key component of the setup is the power-build-up cavity, which enhances the 507 nm intensity and provides the sensitivity required to measure a transition thirteen orders of magnitude weaker than the allowed transition.

CHAPTER 3

POWER BUILD UP CAVITY

The power build-up cavity (PBC) plays a crucial role in driving the highly forbidden 7S–8S transition. In this chapter, I discuss the performance of the PBC. The cavity was first implemented in 2021, and a detailed description of its mechanical design and construction can be found in [33]. However, to provide context for the discussion that follows, I will briefly outline the cavity mirror specifications, the experimental implementation, and the key formulas used to characterize the performance of the PBC.

3.1 PBC setup

The PBC shown in Fig. 3.1 is a two-mirrored cavity where 507 nm light bounces back and forth between two highly reflective mirrors. The radii of curvature r_1 , r_2 of the two mirrors are both 100 cm, and the spacing L is 16 cm. The transmissivities of the input and output mirrors are $T_1 = 900$ ppm and $T_2 = 50$ ppm, respectively. The high reflectivity comes from the multilayer coating of Ta_2O_5 and SiO_2 with ion beam sputtering. Mirror M_1 has a multilayer coating with 22 layers with a thickness of $1.70\text{ }\mu\text{m}$, and mirror M_2 has coatings consisting of 30 layers with a total thickness of $2.27\text{ }\mu\text{m}$. They were manufactured by Five Nine Optics and were received by us in early 2021. The backs of the mirrors are coated with standard anti-reflection coatings. A standing wave is established inside the cavity as a result of the reflective properties of the mirrors. The majority of the injected light is coupled into the fundamental TEM_{00} Gaussian mode rather than higher-order modes. The cavity is operated under high vacuum around 10^{-10} mbar in the science chamber.

To operate the cavity, its length is locked to a transmission peak of the ULE cavity-stabilized laser light using a PID feedback loop acting on the mirror piezos. The locking setup is illustrated in Fig. 3.1. The cavity lock is highly sensitive to mechanical and acoustic vibrations,

in particular since it is not mounted on an optics table but on an ISAC beamline. The

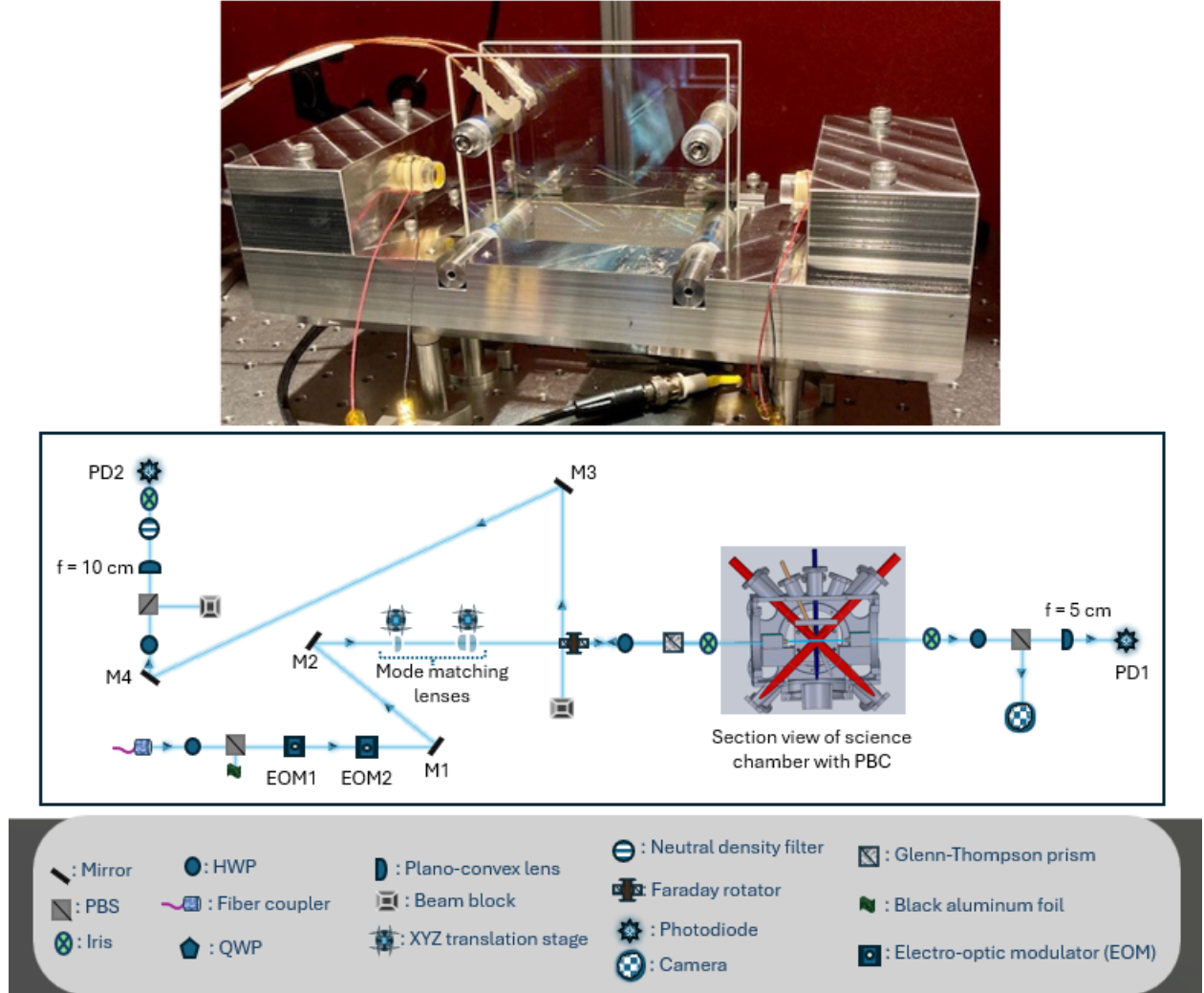


Figure 3.1: Our Fabry-Perot cavity setup with mirrors glued to piezos, electric field plates mounted on the cavity base on a test bench (top). The layout (middle) shows the launch of the blue 507 nm beam (or 496 nm for Rb) to the science chamber and pick up beams for reflected and transmission signals. Electro-optic modulator EOM1 creates FM sidebands at ± 4.0207 MHz to generate a Pound-Drever-Hall (PDH) error signal from the reflected light, which is isolated by a Faraday rotator and steered with mirrors M3 and M4. The signal is detected by photodiode PD2 (Thorlabs, PDA36A2). The Glan-Thompson prism is the final optical component, and it delivers highly linear polarized light into the Fabry-Perot cavity. The transmitted light is sent to the camera for mode observation and also to photodiode PD1 (Thorlabs, PDA10A2) which has a bandwidth of 150 MHz to enable cavity decay time measurements.

structural rigidity needed to suppress mechanical vibrations is achieved by mounting both mirrors on a common cavity base made of 316L stainless steel. The base is shaped to ac-

commodate electric field plates and to allow for clearance of the trapping beams. To isolate acoustic vibrations from nearby equipment and the ISAC beamline, the structure is mounted on PEEK spacers on top of four legs that are bolted to the floor of the science chamber. The purpose of the PEEK spacers is to decouple the cavity from the environment vibrationally. The mirrors are attached to ring piezos using UHV-compatible epoxy.

3.1.1 Quantification of the PBC

The mirror surfaces are characterized by their reflection R , their transmission T , and their absorption/scattering losses B , with energy conservation implying that $R + T + B = 1$. Below we follow the exposition of [33], which is based on the books by Siegman [44] and Bachor [45]. The performance of the cavity is characterized by the amount of light that is circulating inside it, transmitted, and reflected, relative to the input. This can be expressed in terms of gain coefficients when the cavity is on resonance:

$$\begin{aligned} G_{\text{cav}} &= \frac{4T_1}{(B + T_1 + T_2)^2}, \\ G_{\text{trans}} &= \frac{4T_1T_2}{(B + T_1 + T_2)^2}, \\ G_{\text{ref}} &= \frac{(B + T_2 - T_1)^2}{(B + T_1 + T_2)^2}. \end{aligned} \tag{3.1}$$

Ultimately we are interested in the cavity gain G_{cav} , which is an indicator of how much power circulates inside the cavity. Our spectroscopy signal will be proportional to it. Generally, the spectroscopy signal is not available during cavity setup and tuning, and in that case, the cavity gain is not easy to assess as we have no other probe inside the cavity. Here, the transmission gain G_{trans} and the reflection gain G_{ref} become important, as they correspond to the amount of (observable) light transmitted through, and reflected from, the cavity, respectively. Their observation, and that of another quantity, the *finesse* F of the cavity, permit to estimate the gain.

Finesse is defined as the ratio between the *free spectral range* ν_{FSR} of the cavity (the frequency separation between neighbouring resonances) and the width (FWHM) γ of a cavity resonance.

$$F = \frac{\nu_{\text{FSR}}}{\gamma} = \frac{2\pi}{B + T_1 + T_2}. \quad (3.2)$$

The free spectral range is $\nu_{\text{FSR}} = c/2L$ and amounts to 0.9375 GHz for our PBC. The finesse is straightforwardly related to the lifetime τ_{cav} of the light field inside the cavity:

$$\tau_{\text{cav}} = \frac{1}{2\pi\gamma} = \frac{F}{2\pi\nu_{\text{FSR}}}. \quad (3.3)$$

If our PBC is performant, this lifetime should be on the order of one microsecond, and thus can be readily measured as the decay of the transmitted light with a fast photodiode, after the input laser light is quickly turned off. Mirror degradation manifests itself as a reduced lifetime and finesse.

The light is coupled into the TEM_{00} standing wave mode. We lock the cavity when the laser frequency matches the resonant frequency of the cavity, increasing the intensity of light inside the cavity. The initial characterization of the cavity, measured via the decay time of the intra-cavity light field (1.1 μs) was compatible with losses being entirely due to the finite mirror transmissions (intentional, “good”, losses so to speak) and negligible amounts of “bad” losses from absorption and scattering. In a PBC setup as shown in Fig. 3.1, the reflected photodiode measures the light that is not coupled into the cavity and is instead reflected back from the input mirror. When the cavity is off resonance, most of the incident light is reflected, and this appears as a flat line (a high constant level) on the reflected signal as shown in Fig. 3.2. From the third panel in Fig. 3.2, it is evident from a single FSR scan that the input beam is well mode-matched to the cavity, with most of the light in the TEM_{00} mode, ensuring efficient coupling when the cavity is on resonance. At that point, we see a dip in the reflected signal (the cavity is admitting the light), and a corresponding peak in the transmitted signal (light passes through the cavity) as seen in respective panels in Fig.

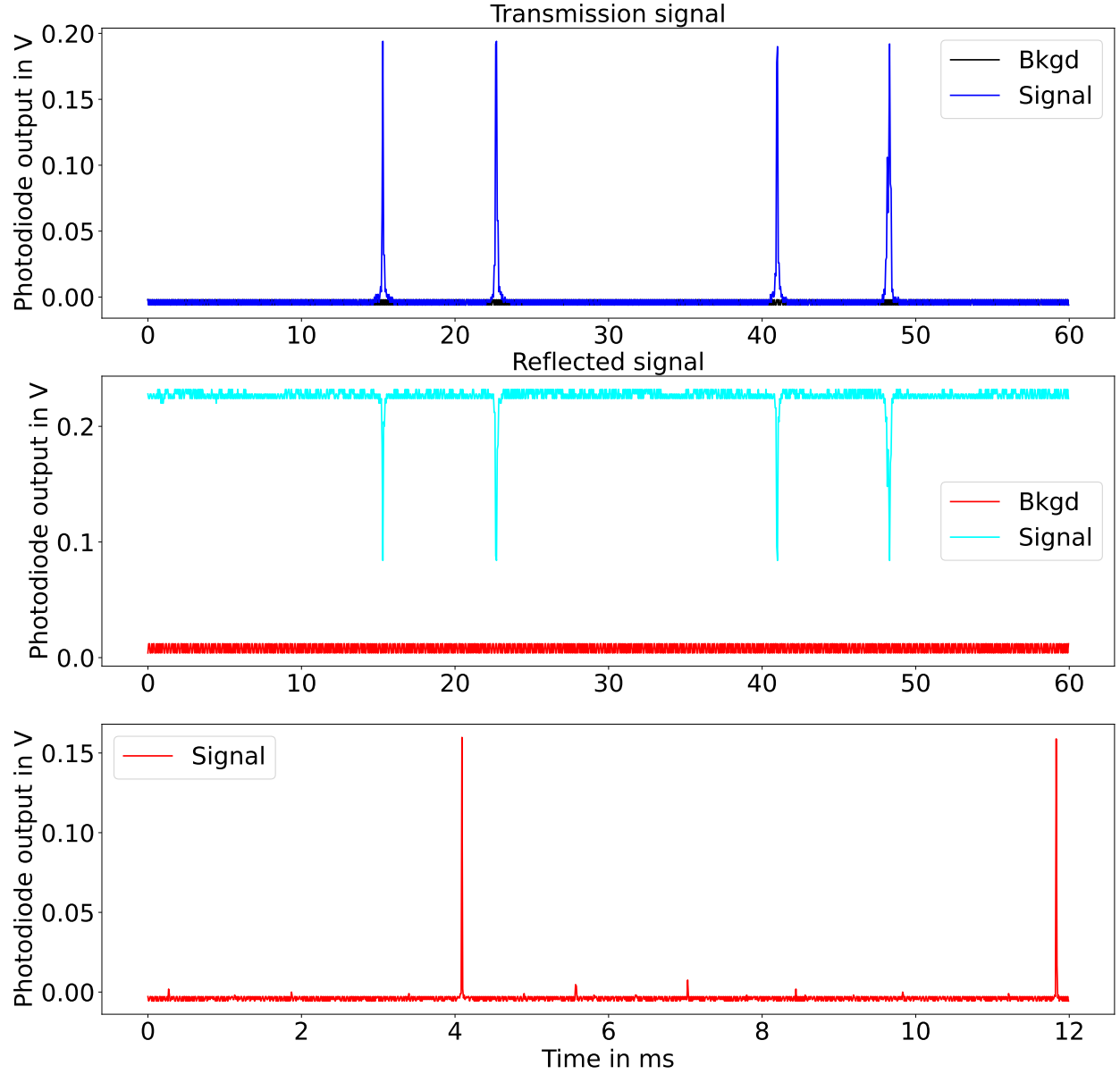


Figure 3.2: shows the transmission signal, the reflected signal measured with photodiodes, and the signal observed for one free spectral range (FSR) scan; the legend ‘bkgd’ refers to the background when the photodiodes were not seeing any light. All three signals give information about the transmission and reflection gains using Eq. 3.1. We scan the cavity over one free spectral range by applying ~ 64 V to the piezo. The x-axis is expressed in ms and corresponds to the voltage ramp that drives the piezos that adjust the cavity length. The coupling efficiency in TEM_{00} mode is estimated by taking the ratio of the TEM_{00} peak height to the sum of all modes seen in one FSR; in this figure, the coupling efficiency is 78%.

3.2.

3.1.2 Locking the PBC

The PBC is locked to the TEM_{00} mode using the Pound-Drever-Hall (PDH) method. The PDH error signal, after processing through an active feedback loop, is sent to a piezo actuator on the cavity mirror to maintain the resonance. In the following section, I briefly introduce the electrical components used to keep the PBC locked. Model names are included to give the reader a sense of how each component is selected to ensure compatibility in bandwidth and frequency range, enabling fast and low-noise control of the piezo.

The reflected signal for PDH is picked up from PD2 shown in Fig. 3.1 and passed through a high-pass filter (Mini-Circuits, ZFHP-1R2-S+) to remove high-frequency components. The filtered signal is then amplified by a low-noise amplifier (Mini Circuits, ZFL-500LN) before being sent to the phase detector (Mini Circuits, ZRPD-1+). The output of the RF power amplifier (Intracore model PA) is split by a splitter (Mini-Circuits, ZSC-2-1), with one branch driving EOM1 in Fig.3.1 and the other serving as a reference signal for PD2. The DC output of the phase detector is fed into an analog PID controller (Stanford Research Systems, SIM 960), which adjusts the proportional, integral, and differential gains. For our setup, the differential gain is crucial for quickly responding to rapid changes in the error signal. The typical PID values used are:

$$P = 2 \times 10^0, \quad I = 1 \times 10^1, \quad \text{and} \quad D = 15 \times 10^{-5}.$$

The output from the PID controller is filtered through a low-pass filter (Stanford Research Systems, SIM 965) to remove DC terms, and the resulting signal is used to actuate the slow piezo on the cavity mirror.

The second piezo on the cavity mirror is driven by an ultra-low noise, high-voltage driver [46] delivered from Joint Quantum Institute (JQI) at University of Maryland. The signal from the JQI box is filtered with the low pass filter (Stanford Research Systems, SIM 965) to remove high-frequency components before being fed to the fast piezo. The fast piezo is

modulated by the voltage from the JQI box to scan the cavity. These piezos (PAHH+0049) from PI Ceramic support modulation from DC to 1 kHz (positive voltage only).

3.2 Performance of PBC

The PBC mirrors and their performance met our expectations as per manufacturer’s specifications during the testing period of the cavity at atmospheric pressure. Once the cavity was placed under vacuum, the decay time was measured to be optimal $1.1 \mu s$. Over the years, we observed that the decay time of the PBC dropped by more than half compared to when the cavity was first introduced into the vacuum. In March of 2023, we opened the science chamber to replace the Rb dispensers and discovered that one of the ion pumps had burned out. We replaced it with a new one. After completion of the maintenance and sealing the chamber, we measured the cavity decay time again and found that it had improved.

Recovery of loss: Over several months, the cavity performance degraded again. Breaking the vacuum and opening the science chamber to expose the cavity optics is not ideal. Although the simplest way to vent the chamber is to completely open it, this approach carries several risks. The air contains water vapor, which is slow to remove under vacuum. In addition, opening the chamber increases the risk of contaminating the interior surfaces and optical components. To mitigate these issues, we explored an alternative recovery method by filling the chamber with air without fully opening it.

Ventilate the vacuum chamber with atmospheric air: We went for a way to vent the chamber with air that is free of water vapor, without opening the science chamber. To achieve this, we installed a Supelco molecular sieve 5A water filter (503118) in the line to leak air into the chamber. This filter with a $5 \text{ \AA}$ pore size removes hydrocarbons and moisture to make air anhydrous. It is a porous layer open tubular gas chromatography column made up of specialty carbon adsorbents. This filter has zeolites ($<80\%$), magnesium aluminosilicate ($<30\%$) and quartz ($<5\%$) [47]. The filter has inner SS Swagelok valves at both ends on

top and bottom. An additional SS-4H Swagelok valve connects the filter to the air supply on one side and to the science chamber on the other. We added an in-line particulate filter with a 2-micron pore size after the molecular sieve to remove any dust to further protect the system.

Procedure to vent the chamber: The first step is to open the air valve and the inner valve of the molecular sieve. Then, bake the molecular sieve filter at up to 80°C overnight to prepare it for absorbing water to vent air through the sieve. Once done, close the bottom valve and wait for the filter to return to room temperature before proceeding with venting. It is important to follow the correct order when opening the valves in order to have controlled air flow to the science chamber. One can choose their course of action that best suits their application. We followed the series of steps discussed below:

The pressure gauge between the top swagelok valve and UHV valve to science chamber is monitored such that the pressure later inside the science chamber can reach the pressure we aimed for. The ion pump is turned off to keep it safe under such high pressure and the valve to the turbo pump is opened to make sure that the turbo pump is ready for the pumping down. The volume up to the final UHV valve to the science chamber is filled with air. The top and bottom swagelok valves are closed down. Finally, we slowly open the UHV valve to the science chamber while taking extra caution. This procedure was used to bring the pressure inside the chamber up to ~ 10 mbar. To bring the pressure inside the science chamber up to 1000 mbar, the bottom Swagelok valve was opened, and the UHV valve to the chamber was also opened. The top Swagelok valve after the filter remained closed initially and was then carefully opened to control the gas flow into the chamber. Once the desired pressure was reached, both the top and bottom Swagelok valves were closed.

3.3 Improvement in PBC performance

We locked the PBC to the peak of the TEM_{00} mode during venting to both photo-induce recovery of the cavity and to monitor the real-time increase in the transmission signal. As

reported in [48]), blue light enhances the recovery process. However, due to limited time and the need to stay on schedule for beamtime preparations, we were unable to study the recovery behavior in the absence of blue light to directly assess its necessity. We vented the cavity up to 10 mbar. In the top graph of Fig. 3.3, a real-time increase in the cavity transmission signal is visible. It was challenging to maintain the lock of the cavity due

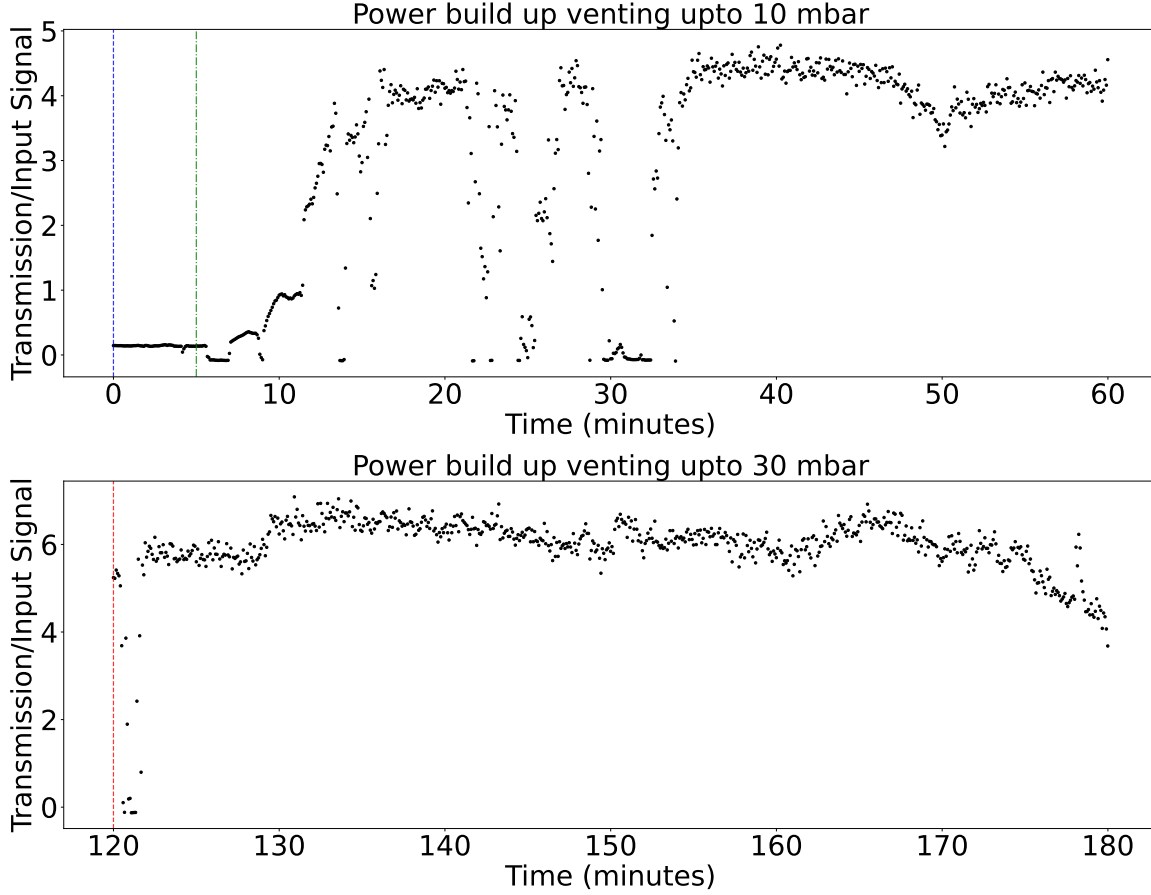


Figure 3.3: The PBC was locked to the peak of the transmission signal to observe real-time recovery. The top figure shows the transmission recovery as the cavity was vented up to 10 mbar. The blue dotted line marks the signal level before venting, and the green line indicates when venting began. Within approximately 10 minutes, the cavity appears to have recovered most of its transmission. The bottom figure shows the recovery after venting to 30 mbar, with the red dotted line marking the point at which the pressure was increased. The drops in the signal correspond to moments when the cavity lost lock.

to pressure variations within the chamber, which caused drops (top panel of Fig. 3.3) in the transmission signal when the cavity lost its lock. Later, the air pressure was raised to 30 mbar, and a slight increase in the transmission signal was observed (bottom panel of

Fig. 3.3). The cavity appeared to recover most of its losses in approximately the first 10 minutes of venting. This recovery of the cavity appeared temporary and gradually began to degrade. One and a half months after the initial venting up to 30 mbar, we vented the cavity up to 1000 mbar and kept the chamber in air overnight. During all these venting procedures, the vacuum remained free of water vapor, which helped us recover the vacuum to 10^{-10} mbar in just over 90 hours, despite having a vacuum chamber of volume $\sim 0.04 \text{ m}^3$. To measure

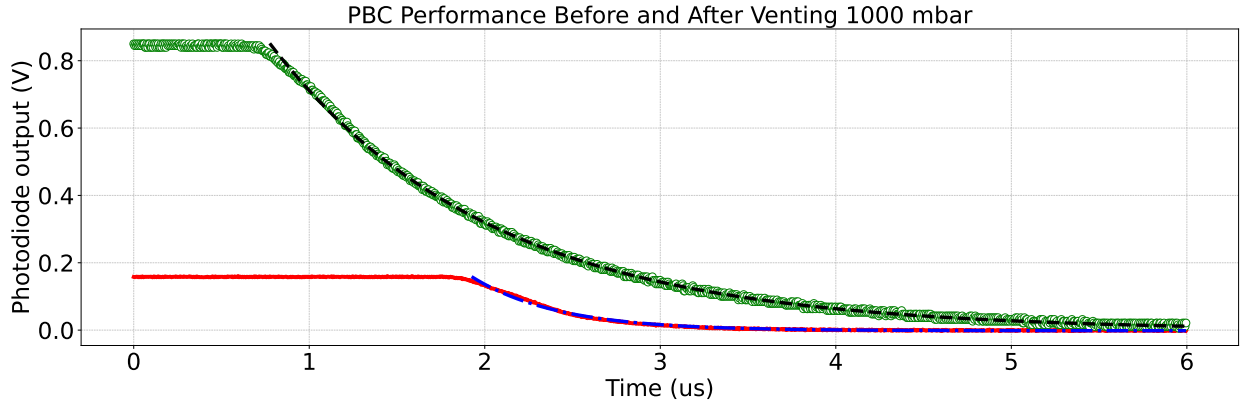


Figure 3.4: PBC transmission signal before and after venting.

Figure 3.5: shows the improvement in cavity performance before and after venting the cavity to 1000 mbar. The cavity was locked with the PDH method, and the decay time was measured via transmission with PD1 in Fig. 3.1. The red curve shows the decay prior to venting, and the green one right after venting. Exponential fits yield decay times of $\sim 384 \text{ ns}$ and $1.1 \mu\text{s}$, respectively. These numbers have been corrected for the finite switch-off time ($\sim 108 \text{ ns}$) of the AOM controlling the incoming light.

the decay constant, the rf signal to the AOM is turned off to interrupt the light going to the cavity, which exponentially decays the built-up intensity inside the cavity. The decay time of the cavity was measured to be 384 ns before venting the cavity, which recovered to $1.1 \mu\text{s}$. Once the cavity was recovered, we measured the decay time of the blue beam inside the cavity to assess the performance degradation over time. Fig. 3.6 illustrates the behavior of the cavity over several days following venting to 1000 mbar. According to equation (3.1), the transmission gain scales with the product $T_1 T_2$ multiplied by the square of the finesse. Thus, any reduction in finesse leads to a quadratic drop in transmission gain as observed by the blue curve in Fig. 3.6. As the loss term B increases over time, the denominator in equation (3.1)

increases, leading to a continuous decline in both finesse and transmission gain. Under the

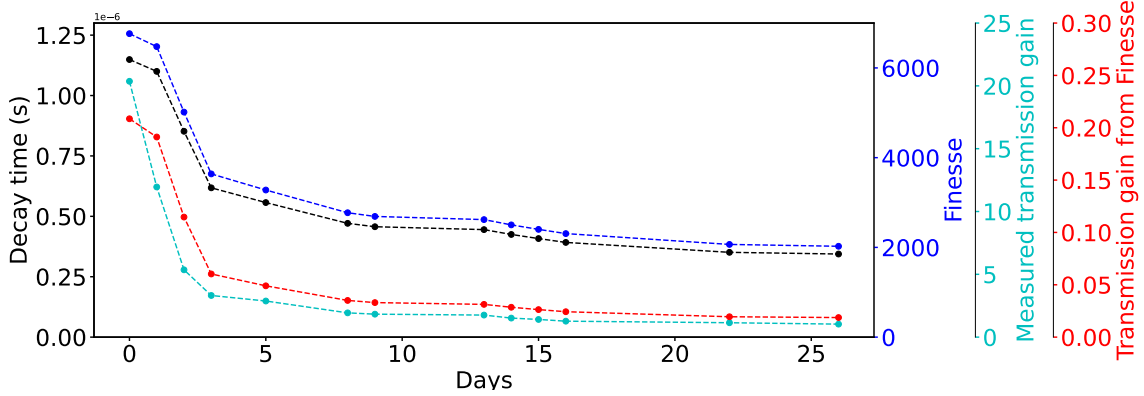


Figure 3.6: It shows the measured transmission gain and the finesse estimated from the decay time. The black line represents the directly measured decay time, the blue line shows the corresponding finesse, the cyan line indicates the transmission gain measured from photodiode data, and the red line shows the transmission gain estimated using the finesse. The numbers for decay time have been corrected for the finite switch-off time (~ 108 ns) of the AOM controlling the incoming light.

assumption that T_1 and T_2 remain constant, the observed degradation is attributed to the increasing B , which quadratically reduces the circulating power and, consequently, G_{cav} . The coupling into the cavity is efficient, and mostly in the TEM_{00} mode, the off-resonance level of the reflected signal (the flat part) represents almost the entire input power shown in Fig. 3.2. We measured the transmission signal and the reflected signal as shown in Fig. 3.2. It can be used as a proxy for the input power when normalizing the transmitted signal. The cyan-colored graph in Fig. 3.6 serves as an additional observable, using G_{trans} to estimate performance degradation. It follows the same trend as the red-colored graph, which is derived from finesse (or decay time).

We found that venting the cavity to 10, 30, or 1000 mbar leads only to temporary recovery, with performance gradually decreasing over time with ~ 50 ppm per day. The decay time of the cavity provides a direct way to estimate losses, and using that, we have tracked how total losses ($B+T_1+T_2$) evolve over time (see Fig. 3.7). After recovering the cavity, the losses start with “good” transmission losses at around 950 ppm and gradually rise to about 2500 ppm over the span of 55 days as absorption increases. To get a rough sense of this progression,

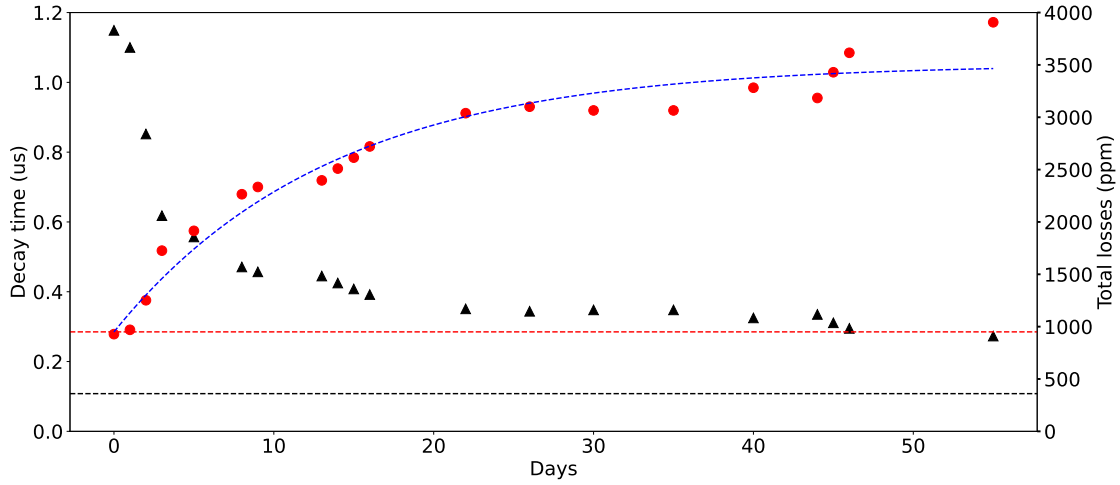


Figure 3.7: shows the degradation in cavity performance after venting the cavity up to 1000 mbar. The triangular data points show the decay time of the cavity, and the ‘red’ circles show the increment in total losses of the cavity. The decay time numbers have been corrected for the finite switch-off time (~ 108 ns) of the AOM controlling the incoming light. The black dotted line shows the time taken by the AOM to turn off the 496 nm light when the rf signal to the AOM is cut. The red dotted line shows the “good” transmission losses as the sum of the transmissivities of M1, M2 ($T1 + T2 = 950$ ppm). The blue dotted line is the fitted exponential curve with function, $y = C + A(1 - e^{Bt})$, where C is offset, fixed at 950 ppm, A is 2558 ppm, rate constant, B as 0.07 per day.

we applied a simple exponential fit. While this captures the overall trend reasonably well, the actual behavior is not perfectly smooth; there are periods of sharp increase, followed by plateaus or slower changes. This could suggest that the loss dynamics might be more complex than a single exponential process. At this point, it is difficult to say whether these features reflect real physical processes, such as surface effects on the mirrors or changes inside the cavity, or if they are simply due to measurement variability. Exploring that in detail would require a deeper and more focused investigation, which is somewhat beyond the scope of the present study. However, the observed rise in losses appears to be consistent with increased absorption based on what we have seen from other reports in the literature.

3.4 Degradation in cavity performance in vacuum

We found that the degradation in cavity performance under vacuum has been reported in [49] at different wavelengths, for different mirrors and coatings but no literature is present to represent the cavity behavior around 500 nm. The literature provides scattered evidence on the root cause of this behavior for different cases [50–53].

Possible cause of degradation: The degradation of mirrors in vacuum could be due to the following possible reasons:

- 1) There could be physical permanent damage to mirror or its surface or a laser induced damage but the repetitive recovery of degradation of mirror performance does not support this possible reason of decline in performance.
- 2) Outgassing of materials from an optical adhesive could contaminate the vacuum. The glue used to attach the piezo and the mirror to the mount is EPO-TEK 302-3M A/B which meets the NASA outgassing requirements. The total mass loss is 0.70, and collected volatile condensable material is less than 0.01 [54] at 25°C.
- 3) There could be contaminants on the mirror surfaces. The mirror surfaces were cleaned during installation in the vacuum chamber, ensuring optimal cavity performance. However, during this period, as mentioned earlier in this chapter that one of our ion pumps began to slowly burn out, which may have introduced contaminants into the vacuum. The hydrocarbon molecular contaminants might have undergone adsorption-desorption competition on the mirror surfaces [55].

Possible explanation for the behavior: There are two possible reasons for observing the degraded cavity behavior.

Case A: Contamination on mirrors 1) As shown in Fig. 3.7, the cavity performance degrades over time. One possible explanation is contamination of the mirrors. A photochemical reaction may occur between surface contaminants and the incident light. Based on insights from the literature, light could photo-dissociate adsorbed contaminant molecules

via a two-photon absorption process, forming chemical bonds on the mirror surface [49].

2) With time, the contaminant grows into a multi-layer deposit on mirror surface affecting the optical properties of mirror, mainly the effective refractive index of the mirror.

Case B: Depletion of oxygen from the dielectric stack of mirror coatings:

A dominant cause of vacuum-induced cavity degradation is the creation of oxygen depletion points in the surface layers of Ta_2O_5 and SiO_2 , as evidenced by increasing absorption losses over time [48]. This depletion leads to non-stoichiometric compositions, such as Ta_2O_x and SiO_y . Reference [48] mentions that the cavity loss was recovered by filling it with pure oxygen, and the degradation was prevented by applying a 110 nm thick SiO_2 layer atop Ta_2O_5 for wavelengths of 370 nm and 420 nm. They also reported that the degradation reduced by a factor of 10 by adding a 1 nm thick passive SiO_2 layer on Ta_2O_5 .

A similar configuration is present in our cavity mirrors, with SiO_2 as the top layer. The exact thickness of this layer is unknown, but is likely greater than 1 nm, judging by the total thickness of the coating and the number of layers. Oxygen depletion in the dielectric stack could be the cause of the observed decrease in the performance of our cavity.

3.5 Summary

To drive the M1 transition, a power-build-up cavity (PBC) was used to intensify the 507 nm light by a factor of ~ 4000 . The cavity mirrors are coated with Ta_2O_5 and SiO_2 , with SiO_2 as the outermost layer. The PBC is locked to the peak of the transmission signal and remains locked while the 507 nm light is frequency-scanned across resonance and chopped to avoid photoionization from the $7\text{P}_{3/2}$ state.

After placing the system under vacuum, we observed degradation in the optical properties of the cavity mirrors at 496 nm. We quantified cavity losses by measuring the cavity decay time, which provides an independent characterization of the cavity. Because the finesse (or G_{cav}) scales quadratically with the decay time, a reduction of the decay time by half results in a quadratic decrease in intracavity power, which directly impacts signal sensitivity. The losses

were recoverable by flushing the cavity with anhydrous air—produced by passing atmospheric air through a porous open-tubular gas-chromatography water filter—at pressures of 10 mbar, 30 mbar, and 1000 mbar. In all cases, the cavity performance improved and the mirror losses were temporarily recovered. However, the recovery did not persist, and the mirrors degraded again over time.

The 496 nm light remained present in the cavity during each flushing cycle. The exact cause of the degradation is unclear. One possibility is oxygen depletion in the mirror coatings; alternatively, failure of an ion pump may have introduced hydrocarbons into the chamber, leading to surface contamination, as discussed in [56]. In practice, venting the cavity to air approximately 20 days before beamtime allows the system to reach a steady loss state and gives the vacuum sufficient time to recover, providing a reliable approach for maintaining intense blue light inside the cavity.

CHAPTER 4

M1/ β MEASUREMENT

4.1 Motivation to measure the ratio of M1 to β

As discussed in Section 1.5.2, when driving the 7S-8S transition, the total transition rate(1.48) receives a contribution from the Stark-induced amplitude $E1_{ST}$, the magnetic dipole amplitude $M1$, and the parity-violating amplitude $E1_{PV}$. The Stark induced amplitude contains two contributions, each characterized by its own polarizability: the scalar transition polarizability, α and the vector transition polarizability, β . Any absolute measurement of the transition amplitude requires precise information of several parameters: the laser power in the PBC, the number of atoms trapped in the MOT, the MOT position between the electric field plates, and the fluorescence detection efficiency in the 8S cascade decay. Determining all of these with the required level of precision makes absolute measurements of the amplitudes extremely challenging. In addition, the parity-violating amplitude is unobservable directly on its own because it has an extremely small oscillator strength 10^{-21} in Fr. The practical solution is to measure one component of the transition amplitude relative to another, so that the unknown experimental factors cancel out in the ratio. If one of the components can be reliably determined from theory, then such ratios constrain the others, which are difficult to calculate and therefore are of particular interest. It is also advantageous to increase the effective signal size of the parity-violation component by interfering it with a much larger parity-conserving amplitude. Thus, to determine $E1_{PV}$, we will measure the ratio of the parity violating component to a well known component of the R_{7S-8S} transition amplitude. The Stark-induced amplitude is a natural choice because it is generated by an applied electric field and can therefore be tuned. In particular, we use the relation $E1_{ST} \propto \beta E$, where E is the known applied electric field. Once β is characterized, the absolute value of $E1_{PV}$

can be extracted from the measured ratio.

1) One way to determine β is by measuring the ratio α/β , which has a predicted value of 5 for francium [57] because α is more reliable to calculate. This ratio is experimentally measurable and provides a meaningful test of atomic-parity violation theory. The quantity α is larger than β [57] and can be calculated more reliably. This is because α has contributions from the $nP_{3/2}$ and $nP_{1/2}$ intermediate states in Eq. 1.44, which contribute with the same sign, as discussed in [57, 58] for cesium. In contrast for β , these contributions enter with opposite signs, thus, making β more sensitive to theoretical uncertainties in the matrix elements. The quantity α is independent of the magnetic sub-level m_F , whereas β depends on it. Consequently, determining β requires knowledge of the population distribution of atoms along the ground state ($7S_{1/2}$) magnetic sublevels of ^{211}Fr . Our current setup does not allow us to control or measure this distribution. In other words, with our current apparatus, we are not prepared to populate atoms in a known m_F state, or we cannot be sure that the m_F distribution in the MOT is sufficiently equal to carry out a meaningful α/β measurement. Therefore, optically pumping the atoms into a well-defined m_F state is the next major objective in the FrPNC lab. The steps taken toward enabling this capability are discussed in the final chapter of this thesis.

2) Another way to determine β is by measuring the ratio $M1/\beta$, since both the $M1$ transition and β have the same dependence on m_F as seen in expressions (1.42) and (1.46). This method also serves as a good way to better understand the magnetic dipole ($M1$) transition amplitude for the $7S \rightarrow 8S$ transition. $M1_{HF}/\beta$ is measured experimentally and due to the fact that $M1_{HF}$ can be determined accurately semi-empirically due to well-known hyperfine splittings, β can be determined.

4.2 Magnetic dipole amplitude

The Hamiltonian that describes the magnetic dipole (M1) transition is

$$H_{M1} = \vec{\mu} \cdot \hat{B}_{\text{laser}}, \quad (4.1)$$

characterizing the interaction between the magnetic dipole moment of the electron, $\vec{\mu} = \frac{e\hbar}{2mc}(\vec{L} + 2\vec{S})$, and the magnetic field component of the laser, $\hat{B}_{\text{laser}}$. Magnetic dipole transitions couple states of the same parity and are generally not restricted to $\Delta N = 0$. However, in a hydrogenlike system, non-relativistically, the matrix element between states with different principal quantum numbers vanishes. The M1 amplitude between the 7S and 8S states in francium is zero in this limit. The relativistic corrections and hyperfine mixing admixes the other states leading to a very weak magnetic dipole transition.

Details of hyperfine-induced mixing and relativistic effects: The hyperfine mixing of 7S and 8S is given by [59]

$$\begin{aligned} |\overline{7SF}\rangle &= |7SF\rangle + \frac{\langle 8SF|H_{hf}|7SF\rangle}{E_{7S} - E_{8S}} |8SF\rangle, \\ |\overline{8SF'}\rangle &= |8SF'\rangle + \frac{\langle 7SF'|H_{hf}|8SF'\rangle}{E_{8S} - E_{7S}} |7SF'\rangle, \end{aligned} \quad (4.2)$$

where $F(F')$ is the total angular momentum quantum number and H_{hf} is the hyperfine-interaction Hamiltonian. To drive the M1 transition between 7S and 8S states, the M1 matrix element between these states is

$$\langle \overline{8SF'}|\mu|\overline{7SF}\rangle = \langle 8SF'|\mu|7SF\rangle + [\langle 8SF|H_{hf}|7SF\rangle - \langle 7SF'|H_{hf}|8SF'\rangle] \frac{\langle F'|\mu|F\rangle}{E_{7S} - E_{8S}}. \quad (4.3)$$

The first term of expression(4.3) is arising from the relativistic effects and the second term of Equation(4.3) refers to the hyperfine interaction mixing. The total magnetic dipole am-

plitude is given by

$$A_{M1} = M1_{\text{rel}} + (F - F') M1_{hf}, \quad (4.4)$$

where $M1_{\text{rel}}$ is the relativistic contribution and $(F - F') M1_{hf}$ is the hyperfine correction. The terms containing parity-violating terms in the total transition rate (1.48) can be neglected because of their small contribution due to their small oscillator strength. The total transition rate can be written as

$$R_{7S \rightarrow 8S} \propto (\beta^2 E^2 + M1^2). \quad (4.5)$$

The above equation shows the quadratic dependence of the transition rate on the electric field. We apply different electric field values across the field plates in SC and measure the total transition rate. There is a standing wave in the PBC where the direction of $\hat{k}$ changes in the interference term, eventually cancelling out the interference term. The M1 / E1 ratio will cancel out the light intensity, the atom number, and the detection efficiency. At $E = 0$ V/cm, we have a pure M1 transition. At values of non-zero electric fields, the transition amplitude is given by the sum of M1 and β , and the quadratic dependence on the electric field has a coefficient equal to β^2 . Due to the hyperfine dependence of the M1 components as shown in (4.4), we measured the amplitude of E1 and M1 in the $\Delta F = \pm 1$ transitions. Using Equation (4.4), the transition rate, $R^\pm$ for $M1/\beta$ for $\Delta F = \pm 1$ is

$$R_{7s \rightarrow 8S}^\pm = (\beta E)^2 + (M1_{rel} \pm M1_{hf})^2. \quad (4.6)$$

The value $M1_{hf}$ can be semi-empirically determined from the well-known 7S and 8S hyperfine splittings. In turn, β could then be extracted from $M1_{hf}/\beta$. The single-particle approximation for non relativistic expression for $M1_{hf}$ is [60]

$$M1^{hf}(\text{single particle}) = - \left| \frac{\mu_B}{c} \right| \cdot \frac{\sqrt{\Delta W_{7S}^{hf} \cdot \Delta W_{8S}^{hf}}}{E_{8S} - E_{7S}} \cdot \frac{1}{2} (g_S - g_I), \quad (4.7)$$

where $\Delta W_{hf}^{7s} = 43.5695 \times 10^9$ Hz[61], $\Delta W_{hf}^{8s} = 9.5625 \times 10^9$ Hz[57], and $E_{8s} - E_{7s} =$

2.4467081 eV^1 $g_I = 0.0004743$ [62], $g_S = 2.002319$. This yields

$$M1_{hf} = -3.450 \times 10^{-5} \frac{\mu_B}{c}. \quad (4.8)$$

$M1_{hf}$ in cesium has been calculated with very high confidence. The theoretical uncertainty is up to one quarter of a percent [63], which means that the semi-empirical calculation of $M1_{hf}$ is extremely precise and reliable.

4.3 Details of the measurement

The $7S \rightarrow 8S$ transition rate (for $\Delta F = \pm 1$), as shown in Fig. 4.1, is measured using a 507 nm beam in a power build-up cavity (PBC). A positive and negative voltages are applied to the electric field plates inside the science chamber (SC) using SRS P300 series power supplies (P375 for positive, P370 for negative voltage) to drive Stark-induced transitions. The 1013 nm near-infrared (NIR) light from the blue laser is sent to the ultralow expansion (ULE) [43] cavity via a fiber electro-optic modulator, as shown in Fig. 2.11, which has a Phase Matrix Synthesizer to scan the frequency over the resonance. This NIR light is locked to a sideband in the TEM₀₀ mode of an ULE [43] cavity. The sideband, generated using a Phase Matrix synthesizer (see Chapter 2), was scanned across resonance from 6050 to 6090 MHz in 80 upsteps and 80 downsteps (500 kHz/step, 50 ms dwell time, 15 dBm sideband power) on the $\Delta F = -1$ transition, with the wavemeter reading from 9865.82448 to 9865.82323 cm⁻¹. For the $\Delta F = +1$ transition, the Phase Matrix was scanned from 6454 to 6494 MHz under similar settings, with the wavemeter reading from 9866.71004 to 9866.70872 cm⁻¹. The trap and repump laser readings on the wavemeter were 13923.38370 cm⁻¹ and 13924.80373 cm⁻¹

¹For the transition $\lambda_{7S \rightarrow 8S}$, the wavelength is given by $\lambda_{7S \rightarrow 8S} = \frac{hc}{E_{8S} - E_{7S}}$, where the energy difference is in eV. To convert from eV to MHz: use Planck's constant $h = 6.626068 \times 10^{-34} \text{ J} \cdot \text{s}$, and the conversion factor $1 \text{ eV} = 1.602176634 \times 10^{-19} \text{ J}$. The frequency in MHz is given by:

$$\text{MHz} = \frac{E_{7S \rightarrow 8S} (\text{eV}) \times 1.602176634 \times 10^{-19}}{6.626068 \times 10^{-34}} \times 10^{-6}$$

$$E_{7S \rightarrow 8S} = 2.4467 \text{ eV} = 591.7 \text{ THz} \Rightarrow \lambda_{7S \rightarrow 8S} \approx 507 \times 10^{-9} \text{ m}.$$

respectively. These values may assist the local team in tracking resonance drift in the wavemeter reading over time in the future. Before entering the PBC, the beam enters the doubling cavity as NIR light and exits as blue light. The frequency of the blue light is scanned over a range of 80 MHz across the 7S-8S resonance in a symmetric up-down pattern in 12.8 s. The transition signal is detected using the burst detection technique (as introduced in chapter 2), targeting upper and lower burst transitions. To derive the β transition, the

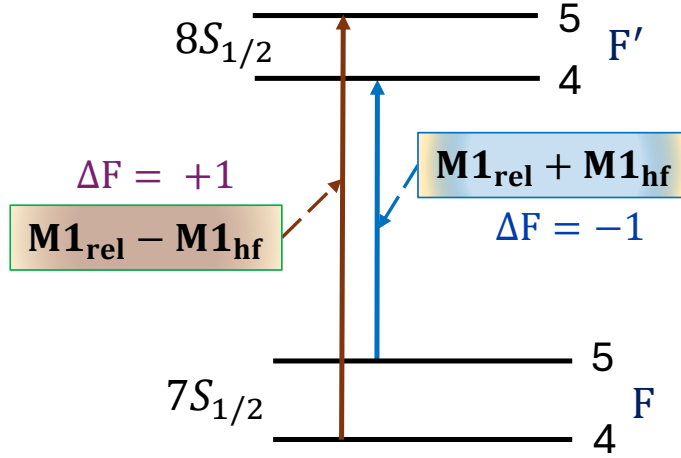


Figure 4.1: $\Delta F = \pm 1$ hyperfine transitions ^{211}Fr .

polarization of the blue laser is perpendicular to the electric field and to the propagation direction. This condition of light alignment to the electric field is achieved by using a half-waveplate and a Glan-Thompson prism as the last element before light enters the PBC. The blue beam (507 nm) is passed through an AOM with +110 MHz. The blue beam is sent to the PBC which stays locked to the TEM_{00} mode using the Pound Drever Hall (PDH) method, and is cut off by turning off an AOM. The timing sequence for the $M1/\beta$ measurement shown in 4.2 is structured such that each (re)trap-measurement cycle within a scan has a duration of 2 ms. In each cycle, atoms are (re)trapped for 1 ms. After that, the repumper laser stays on for an additional 90 μs to ensure that the atoms end all up in the upper hyperfine ground state. If we want to probe the transition starting from the lower hyperfine ground state (the $\Delta F = +1$ transition), the atoms are first trapped on the upper hyperfine ground state and then get pumped into the lower hyperfine state using a depumper beam.

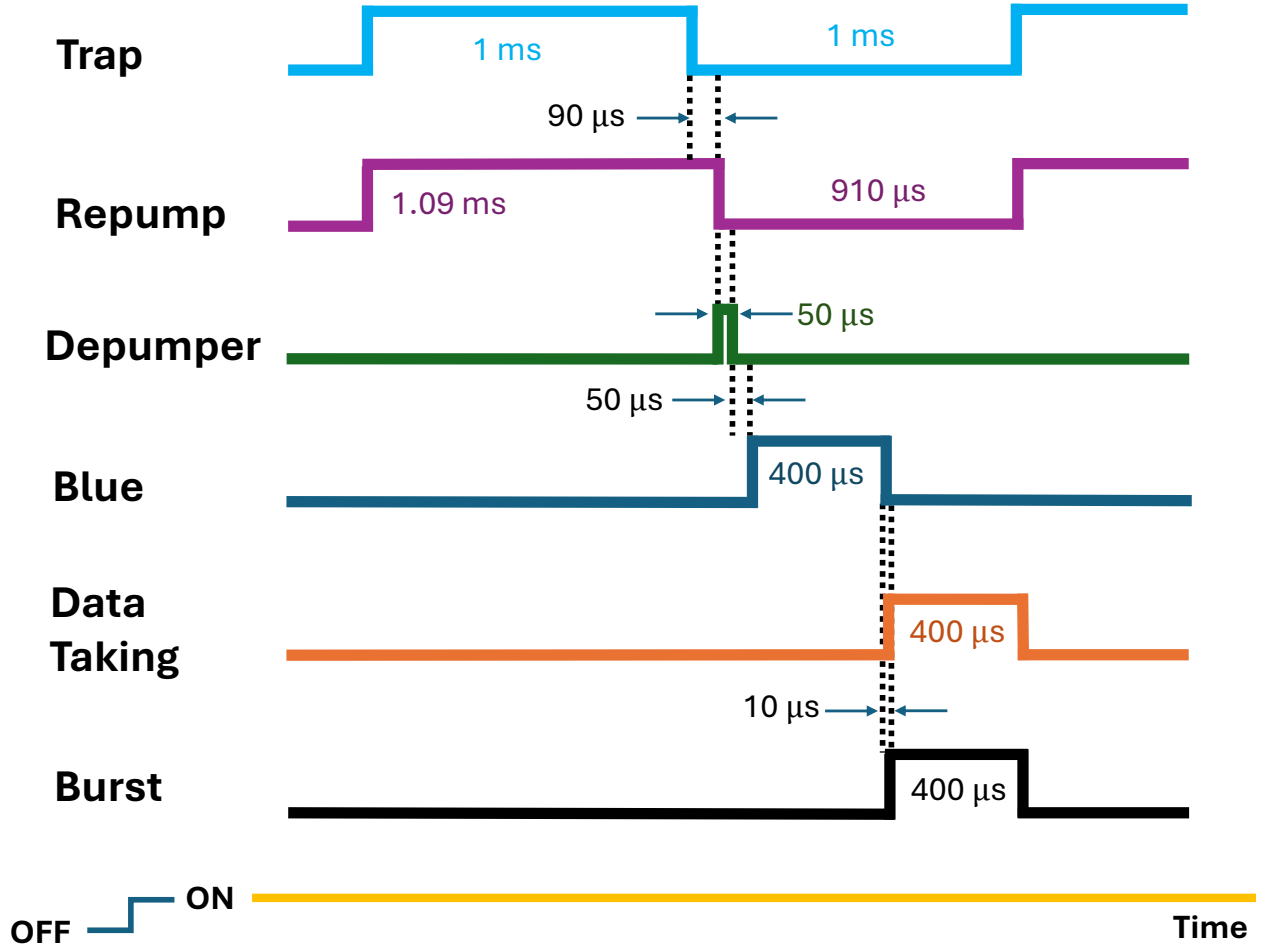


Figure 4.2: Timing sequence for the $M1/\beta$ measurement.

Once the atoms are prepared in the desired hyperfine ground state, a 400 μ s pulse of blue light is applied to drive the forbidden $7S_{1/2} \rightarrow 8S_{1/2}$ transition. Atoms that complete the subsequent $8S \rightarrow 7P \rightarrow 7S$ cascade decay are then probed using a burst beam for 400 μ s. The detection window overlaps with the burst pulse, allowing data collection during this time.

4.3.1 Atom number and its calculation

The number of atoms in the MOT is estimated from the CCD camera data by fitting an exponential decay function:

$$y = Ae^{-t/\tau} + \text{offset}. \quad (4.9)$$

Here, A (amplitude) and τ (decay time) are free parameters, and offset is a constant. The atom number is calculated as the difference between the fitted amplitude and the mean offset after the trapped atoms in SC are pushed away. The amplitude and decay time are floating parameters, while the offset is kept as a fixed parameter. Data are collected for 12.8 s. The background (offset) is determined by pushing the atoms out of the trapping region using a “reset beam” at around 12 s. This reset beam is the light resonant to the trap frequency and is picked off with a PBS and a half-wave plate from the trap beam going to the SC. Each experimental (re)trap-measurement cycle records a `.txt` file containing three columns: time, buffer, and signal. The data are filtered by removing duplicate entries from the third column, corresponding to the CCD signal. The raw signal³ is converted to an atom number using:

$$\begin{aligned} \text{Number of atoms} = & \frac{\text{Photons per count} \times \text{ROI}}{\text{Exposure time} \times \text{Fraction of solid angle}} \\ & \times \text{Scattering rate} \times \text{Camera fraction}. \end{aligned} \quad (4.10)$$

Photons per count are derived from electrons per count and quantum efficiency. The camera, Point Grey Flea3 (FL3-GE-03S1M-C), is expected to have quantum efficiency, η , of $\sim 35\%$ ⁴. Although the exact quantum efficiency η of the FL3-GE-03S1M-C camera at 718 nm is not known, a similar model, the Blackfly BFLY-PGE-05S2M-CS, should have a comparable sensitivity⁵. For 718 nm with a 12-bit ADC and a saturation capacity of 14508, the photons per

³Internal records for the elog: elog.triumf.ca/Francium/General/696 and elog.triumf.ca/Francium/Experiments/395

⁴Near IR η is mentioned to be $\sim 21\%$ at 850 nm; see this PDF (page 10).

⁵See this datasheet for camera at upper chamber (page 3) for QE values at various wavelengths.

count comes out to be $\frac{\text{electrons per count}}{\eta} = \frac{\text{Saturation Capacity}}{2^{\text{ADC}} \times \eta} \approx 10$. The ROI (region of interest) encompasses the area of maximum MOT fluorescence plus background. The exposure time was fixed at 15 ms. The solid angle fraction based on the lens diameter and its distance to the atoms is 0.015. The scattering rate is estimated from MOT beam diameters, beam intensity, detuning⁶, and light coupling efficiency input to the science chamber. The camera fraction denotes the portion of light that reaches the camera after division by a polarizing beam splitter and a half-wave plate ($\lambda/2$) positioned between the SiPM and camera, and is 0.22. The raw data were summed over all cycles, then averaged, and the decay model was

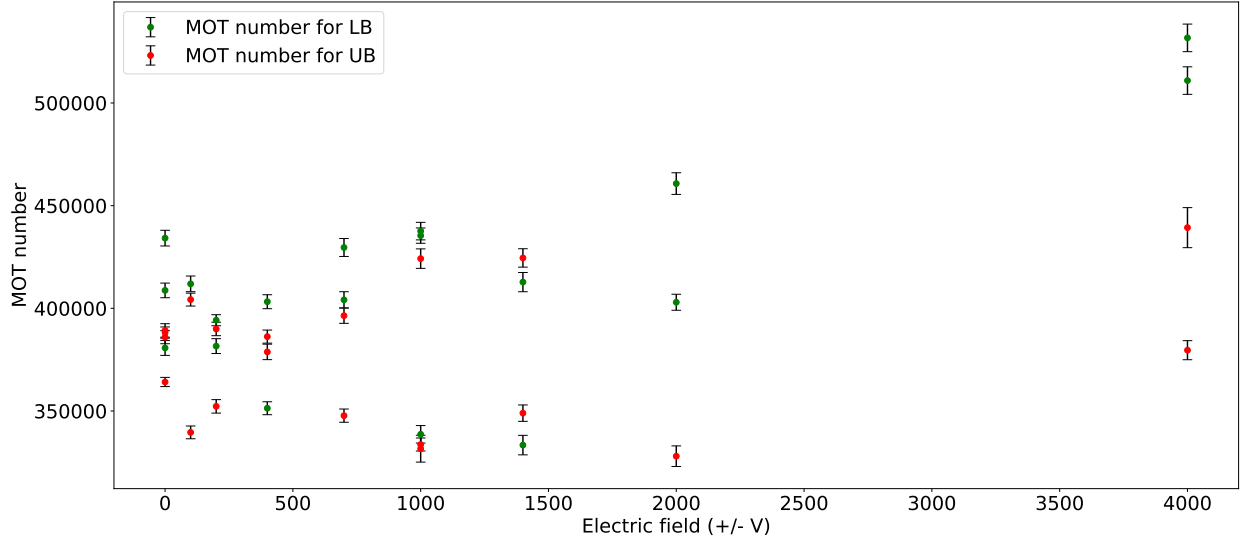


Figure 4.3: Atom number for UB and LB.

fitted to this averaged curve to estimate the atom number. The uncertainty in the number of atoms in each data point is $\sqrt{N}$. Assuming that each cycle is independent, the errors from each file are added in quadrature, and the mean variance is calculated as:

$$\sigma_{\text{avg}}(t) = \sqrt{\frac{1}{N} \sum \sigma_i^2(t)}. \quad (4.11)$$

⁶The detuning is considered here to be -3Γ . The exact resonant frequency for $7S_{1/2}(F=5) \rightarrow 7P_{3/2}(F'=6)$ are not well established like rubidium, so our estimates for the trapping frequencies for Fr MOT (translated from Rb MOT because the trapping parameters like trapping power in MOT beams, MOT beam diameter, B field gradient and so on remain similar for both Fr and Rb) are red-detuned by $3-4\Gamma$.

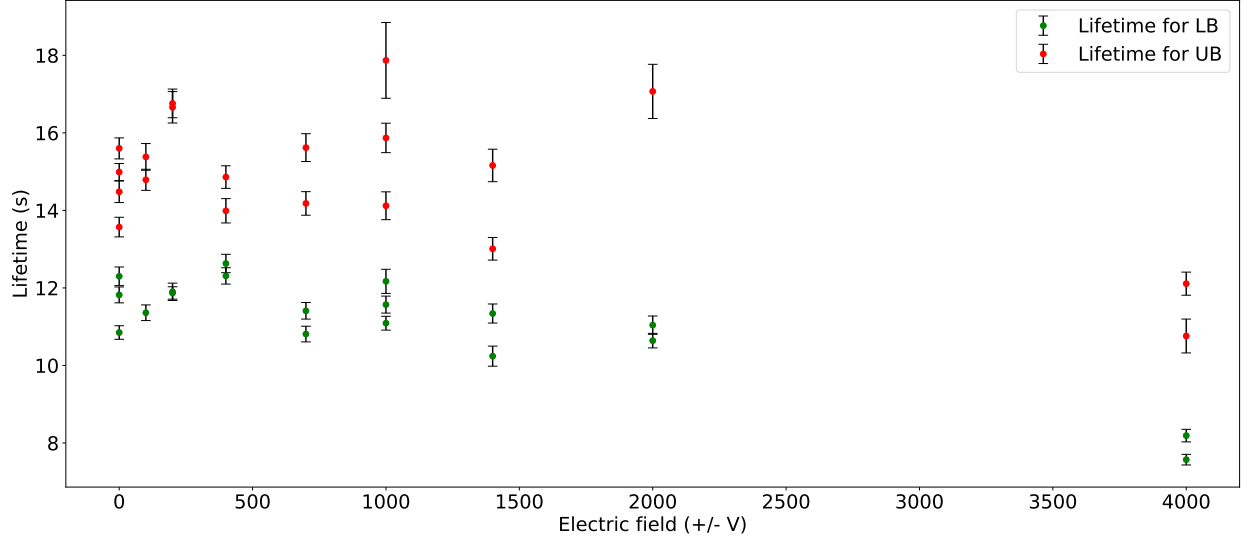


Figure 4.4: Lifetime of atoms in MOT for UB and LB.

We fit the exponential decay of atom number with *iminuit* and the exponential fit is only fitted to the data before the reset beam pushes the trapped atoms out of MOT in data taking window.

$$\chi^2 = \sum \left(\frac{\text{data} - \text{model}}{\text{error}} \right)^2. \quad (4.12)$$

The χ^2_{reduced} is when χ^2 is divided by the number of data points minus the number of free parameters (degrees of freedom) and is then scaled by multiplying by $\sqrt{\chi^2_{\text{reduced}}}$ to ensure that the error estimates are not off or $\chi^2_{\text{reduced}} \sim 1$. The estimated number of trapped atoms in the MOT is about 400000 as shown in Fig. 4.3. The average lifetime for lower burst (LB) and upper burst (UB) is ~ 13 s as shown in Fig. 4.4.

Blue light: The blue light plays an important role in normalization because variations in its intensity within the interaction region can affect the transition rate. Therefore, we monitored the blue light as one of the systematic factors used to normalize the transition rate. It was monitored using an oscilloscope (in milli-volt scale). During scans of the Phase Matrix Synthesizer over the resonance, the PBC lock exhibited a spread that fluctuated, or “breathed.” This behavior appears to be linked to the ULE cavity lock. The ULE is locked to a sideband generated by the phase matrix, which means that the distance between the laser

frequency and the sideband is kept fixed. As the phase matrix moves through the frequency steps, the sideband shifts. To maintain lock, the laser beam follows the sideband, ensuring that the ULE remains locked to the sideband. As the phase matrix scans, any feedback from the ULE lock system—particularly from its PID loop—is passed to the master blue laser. As a result, any fluctuations, lags, or corrections made by the ULE PID appear in the laser behavior, and this influences the PBC lock. Although this effect is present, it behaves consistently across all scans, affecting them in the same way. Additionally, during data collection, the blue light in the PBC is chopped, and the cavity lock takes some time to reach the peak of the transmission signal. However, since the fluctuation is consistent, the effect averages out over the scan. The minimum and maximum values (in mV) of the locked blue light on the PBC were estimated with limited precision across all scans. If we look at the variation in the maximum bound on the blue light, for UB the measured values have a mean of 35.89 mV with a standard deviation of 0.94 mV, indicating a moderate spread around the average. The uncertainty in the mean, calculated as the standard error, is approximately ± 0.24 mV. For LB, the mean is 73.31 ± 0.25 mV with standard deviation of 1.10 mV. A relatively generous error estimate of 5% was assigned to account for the uncertainty inherent in this rough measurement approach.

4.4 Analysis

We started the measurement with $\Delta F = -1$ hyperfine transition where we excited the $7S(F = 5) \rightarrow 8S(F' = 4)$ transition. We use Python 3.9.12 for data analysis. The process involves preparing the data, fitting a profile to extract the peak area under the curve, and plotting the quadratic relationship between the electric field and the area to determine the value of $M1/\beta$. A total of 39 measurements are taken at each of the 160 frequency steps, resulting in 6240 data points. The data is reshaped into appropriately sized arrays, and the counts at each frequency step are obtained by summing the measurements. The counts (binned data) are divided by the number of cycles in each scan, bin width (data-taking

timing window, 400 μs) and summed points per cycle for each frequency step (39 points). The standard deviation (uncertainty) of the counts is equal to the square root of the counts, and the uncertainty is also divided by the number of cycles, bin width, and points per cycle. Since the Phase Matrix goes twice over the resonance, say going up to down and down to up in frequency. We observe two resonances in our full scan, where technically either of them could be used for analysis. However, we prefer mostly the first one, where the signal size is relatively larger than the second peak because the atoms in MOT have not largely decayed for the first peak. A Lorentzian and Voigt profile multiplied with exponential trap decay is fitted to the forward resonance at each electric field value. The background under the Lorentzian and Voigt profile² is a constant offset plus a ramp. The initial parameter guesses were chosen as follows: the background offset was set to the mean signal after atoms left the interaction region; the peak height to the maximum signal value; the decay rate from an exponential fit to the MOT lifetime using CCD data; the center frequency from the frequency corresponding to the maximum signal; the linewidth as 3 MHz; and an initial guess was made for the background ramp. The exponential decay of atoms in MOT is ~ 13 s. The fitted parameters are extracted from the minimized χ^2 using *iminuit*

$$L_{mod}(\nu) = b_0 + e^{-\nu/b_2} \frac{2b_1}{\pi\gamma} \frac{(\gamma/2)^2}{(\nu - \nu_0)^2 + (\gamma/2)^2}, \quad (4.13)$$

where b_0 is the background (constant offset plus ramp), b_1 is the scaling factor times the area under the normalized profile (which is 1), b_2 is the lifetime of trapped atoms in MOT, γ is the linewidth, ν_0 is the peak position. The Voigt profile is modified the same way as the Lorentzian profile and is as follows:

$$V_{mod}(\nu) = b_0 + b_1 \text{Voigt Profile}(\nu, \sigma, \gamma) e^{-\nu/b_2}. \quad (4.14)$$

²The Voigt Profile is a convolution of a Lorentzian and a Gaussian function. For more details, see the SciPy documentation and the Wikipedia article.

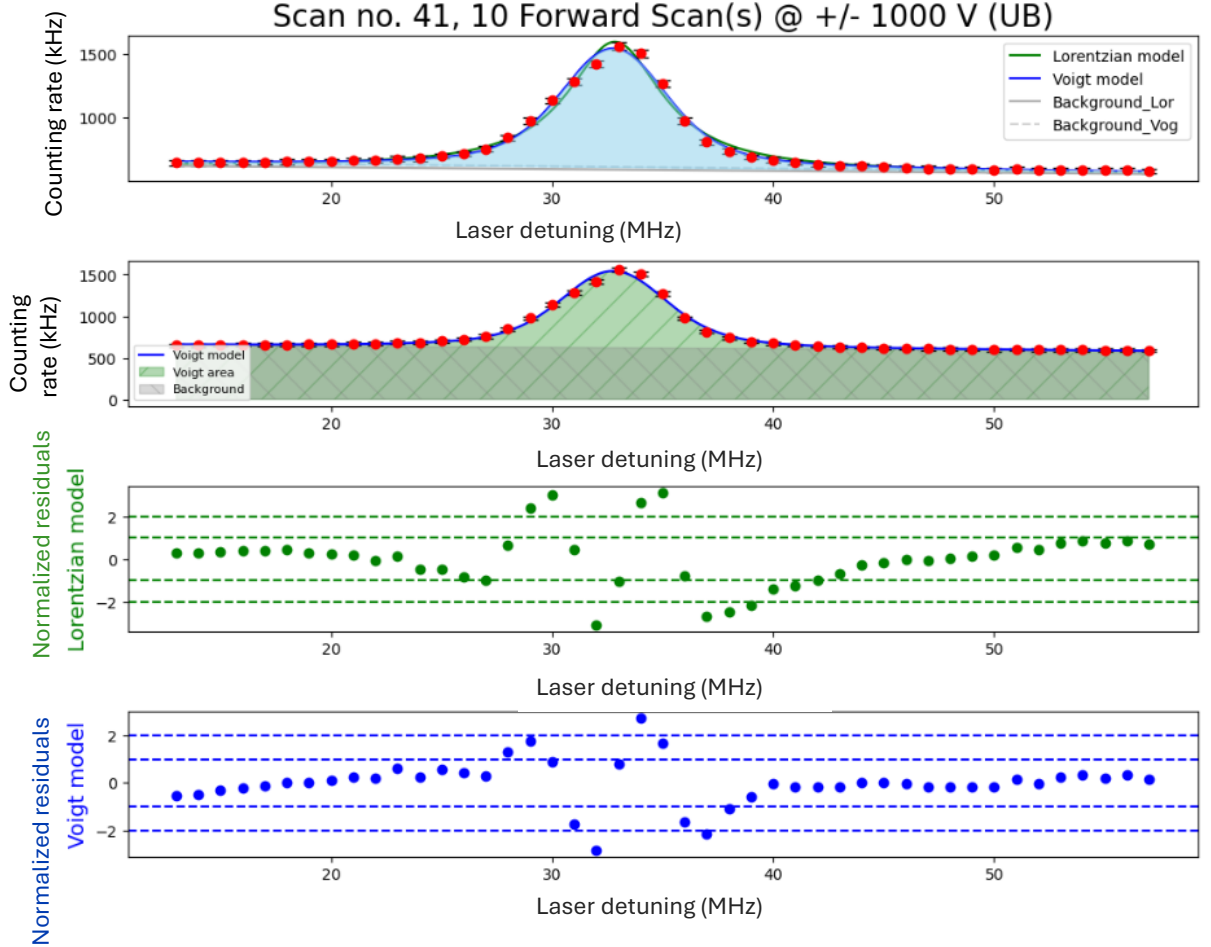


Figure 4.5: One of the scans from the FrPNC beamtime in December 2024. This figure shows a *forward scan*, meaning the phase matrix crosses the resonance first in the up-down and down-up scans, with the electric field set to 1000 V on each field plate for the $\Delta F = +1$ transition. The first subset shows the fitted Lorentzian and Voigt profiles to y-axis with suppressed background, while the second subset shows the same plot with the y-axis starting from zero. This helps to visualize the role of background in fits because for instance for the signal at large electric field values, the background played minimal role and the fit function was mostly ruled by the resonance. But at the low electric field values, the signal observed on top of large background shows that the background plays a vital role for fitting and small variation in background could have affected the signal significantly. The third and fourth subsets show the normalized residuals for the Lorentzian and Voigt fits, respectively.

The fitted profile is normalized to the unit area, so the area under the peak is determined by the scaling factor applied during the fitting. The χ^2 is minimized by varying the background, linewidth, and peak position fit parameters, while keeping the MOT trap lifetime fixed in both fit models; the MOT lifetime is determined from the CCD data. For the Voigt model, the reduced χ^2 ranges from 0.82 to 76.33 (LB) and 1.14 to 98.26 (UB), increasing with electric field strength. For the Lorentzian model, the reduced χ^2 ranges from 1.07 to 249.45 for LB, and 1.17 to 154.74 (693.60 for one scan at ± 4 kV, Scan 55) for UB. The high χ^2 value is attributed to the pattern seen in the residuals (fourth panel of Fig. 4.5). To ensure a consistent evaluation of model behavior and fit quality, the error bars are rescaled by $\sqrt{\chi^2}$ to set the reduced χ^2 to 1. Although the Voigt model yields a lower χ^2 value than the Lorentzian model, this may be due to its additional degree of freedom. However, even when the degrees of freedom are matched, the Lorentzian model does not outperform the Voigt model. This suggests that the Voigt model provides a better fit, particularly at higher electric-field values. The transition rate is obtained when the counts are normalized to the atom number from the CCD and has an additional uncertainty added from blue light. The numerical parameters obtained from peak extraction across all scans at varying electric fields are presented in Tables C.7 and C.8.

Linewidth of transition: The expected linewidth for the 7S–8S transition is 2.99 MHz. However, during this 2024 beamtime, we observed a broad linewidth of 4.8–6.73 MHz, nearly twice the expected value at all electric field strengths. If the broadening were due to saturation effects from hyperfine pumping, as seen in the 2021 measurement [33] at high electric fields, it would not be present at low fields. In contrast, our data show significant broadening even at 0 V for both the UB and LB, indicating that the effect is not solely field dependent. In addition to the broadening, we observe a small but consistent asymmetric distortion (a shouldered feature as seen in Fig. 4.5), near resonance, with peak centers separated by approximately 10 kHz.

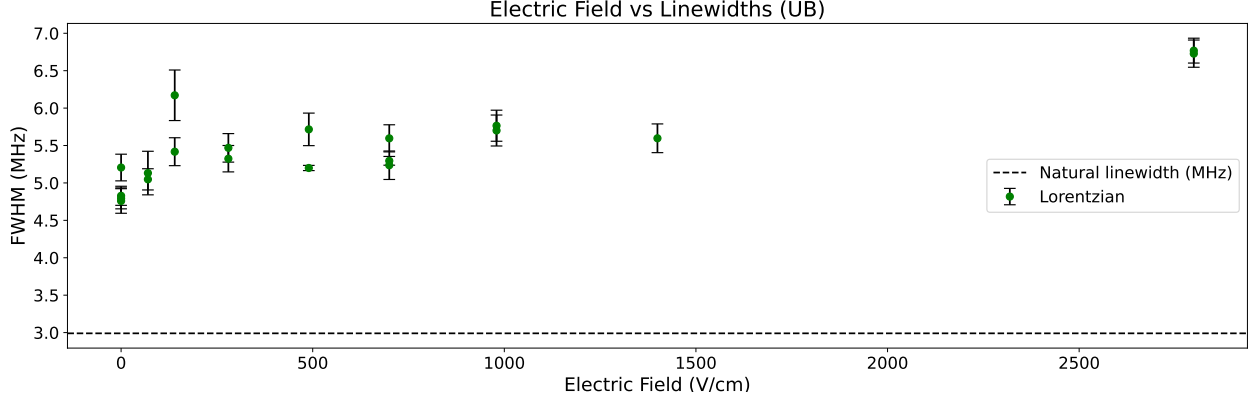


Figure 4.6: Linewidth for UB.

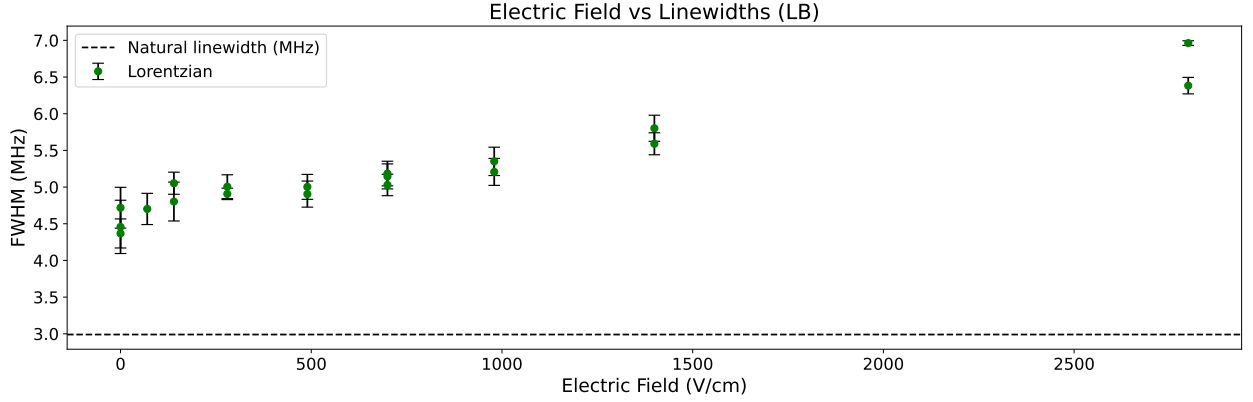


Figure 4.7: Linewidth for LB.

Carl Wieman et al. reported similar line-shaped asymmetries in their M1 measurement in Cs, where they attributed the distortion to spatially varying ac Stark shifts in a strong standing-wave field [64]. In our data, the distortion persists even at 0 V and shows a systematic directionality: it appears more frequently on the right side of the resonance for LB and on the left side for UB. The origin of this effect remains unresolved.

4.4.1 Transition rate as a function of electric field $\Delta F = \pm 1$

The transition rate is plotted as a function of electric field strength (V/cm) for the upper and lower bursts on the $\Delta F = \pm 1$ transitions, as shown in Figs. 4.8 and 4.9. The data show a clear quadratic dependence of the resonance strength on the electric field, with a non-zero

offset at 0 V caused by the M1 contribution. The curvature of the parabola gives the value of β^2 . Although the two quadratic curves appear different, the quadratic part of the transition rate should be identical. However, the two quadratic curves cannot be directly compared because $M1/\beta$ differs due to the M1 offset, $(M1_{\text{rel}} \pm M1_{\text{hf}})^2$. At higher electric fields, both figures reveal a saturation effect, shown by the solid blue lines in Figs. 4.8 and 4.9, which arise from *hyperfine pumping*. For example, in transition $7S(F=4) \rightarrow 8S(F'=5)$, the atoms can decay to $7S(F=4)$ or $F=5$. Due to the large ground-state hyperfine splitting (~ 45 GHz), atoms that decay to $F=5$ no longer interact with 507 nm light and are effectively lost from the excitation cycle. At high electric fields, the excitation rate of the forbidden $7S_{1/2} \rightarrow 8S_{1/2}$ transition becomes large enough so that the atoms are rapidly transferred to the dark state. This leads to population depletion of the initial state, which causes observable lineshape broadening, as shown in Fig. 4.5. To account for this, we introduce a saturation factor c into the fit function and perform a modified quadratic fit $(1 - e^{\frac{-(ax^2+b)}{c}})c$, where a , x and, b signify β^2 , E , $M1^2$, respectively. This saturated quadratic curve fits the data, with the lowest reduced χ^2 : 5.18 for the UB and 2.28 for the LB. For comparison, the dotted curve uses fixed parameters a and b from the saturated fit, but omits the saturation factor. The saturation effect is stronger at high electric fields because the transition rate varies quadratically for the electric field E . It is less for the UB than for the LB because at the time of the measurement, the amount of blue light for the UB was roughly half that of the LB. The error bars are also inflated to make the reduced χ^2 one so that the error consideration be realistic. The second panels in Figs. 4.8 and 4.9 show the normalized residuals of the saturated quadratic fits. The third panels zoom into the low electric field region, where both the standard and saturated quadratic fits align closely, indicating that saturation and hyperfine pumping are negligible at low electric fields. To avoid monotonic or slow drifts in parameters like laser power or PID locks, data were collected in two sets with electric field values jumbled over time for each transition. A quadratic fit to the data yields the parameters β^2 and $M1^2$. The square root of the ratio of the offset (M1) to the quadratic coefficient (β) gives $M1/\beta$.

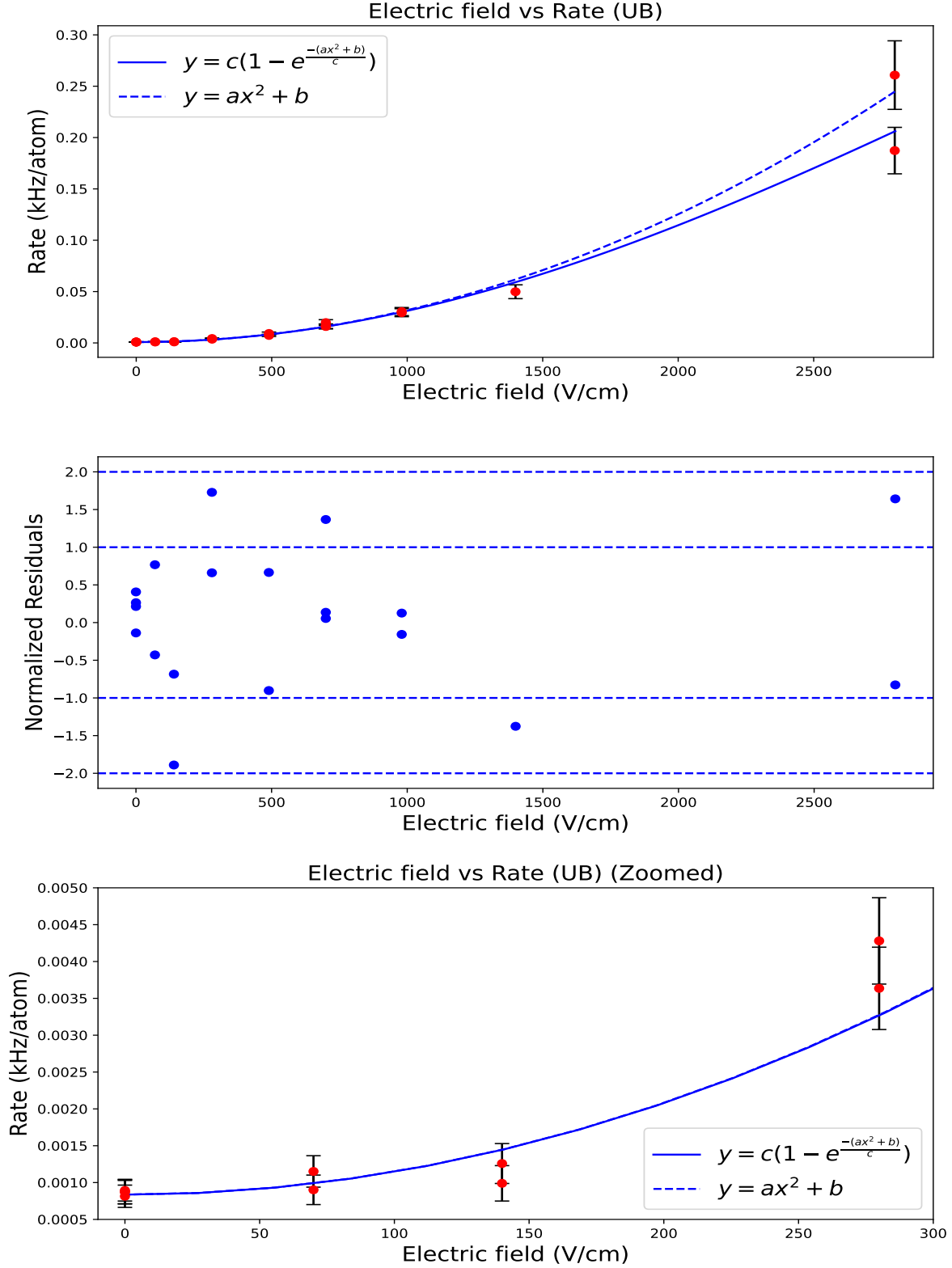


Figure 4.8: Quadratic dependence of transition rate as a function of electric field on $\Delta F = +1$, $7S_{1/2}(F = 4) \rightarrow 8S_{1/2}(F' = 5)$ with Voigt model. The values of fitted parameters from solid line are $a = (3.1 \pm 0.2) \times 10^{-8}$, $b = (8.4 \pm 0.6) \times 10^{-4}$, $c = (6.9 \pm 4.8) \times 10^{-1}$. $M1/\beta$ is 163.8 ± 8.1 V/cm .

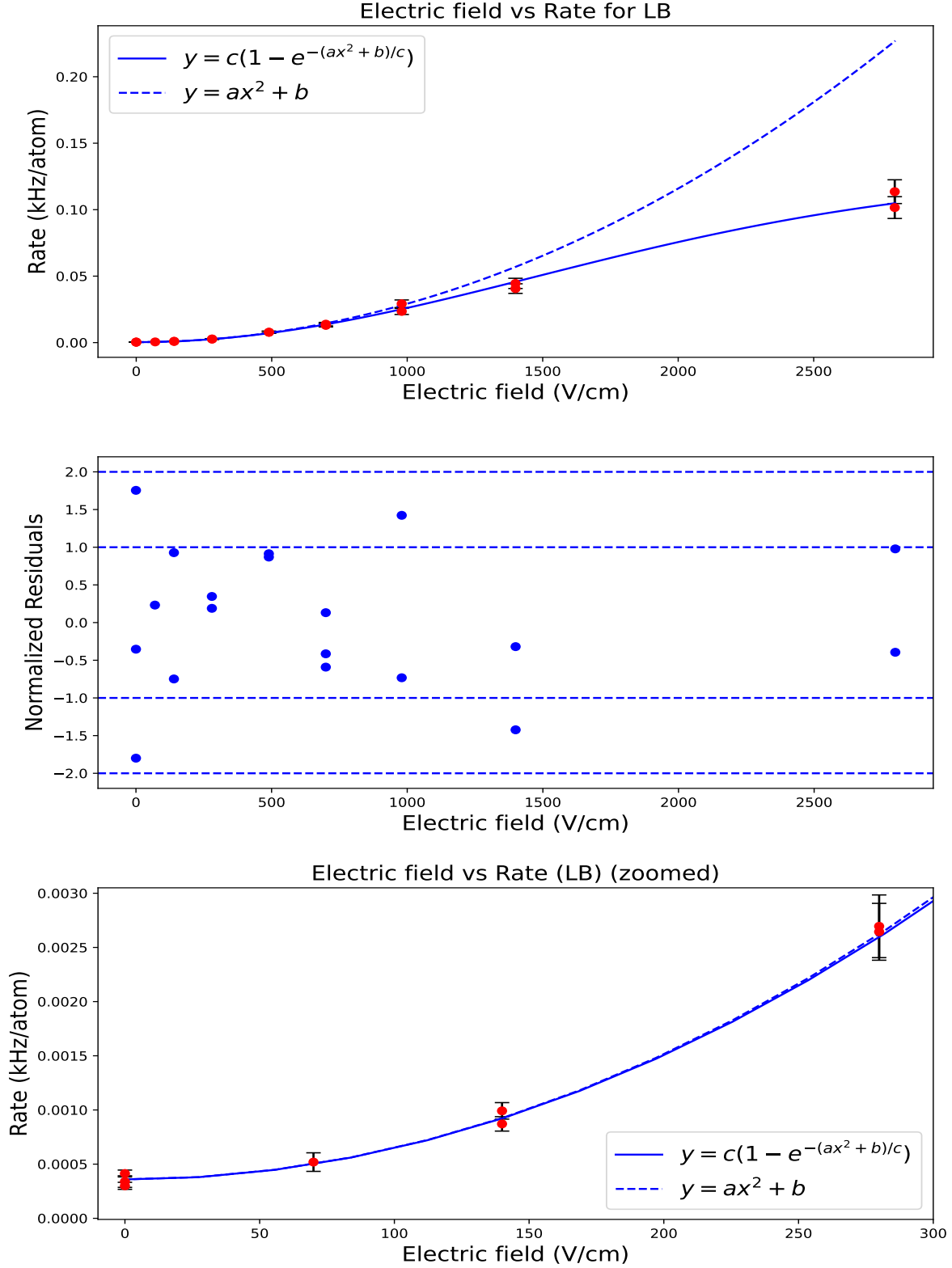


Figure 4.9: Quadratic dependence of transition rate as a function of electric field on $\Delta F = -1$, $7S_{1/2}(F = 5) \rightarrow 8S_{1/2}(F' = 4)$ with Voigt model. The values of fitted parameters from solid line are $a = (2.9 \pm 0.1) \times 10^{-8}$, $b = (3.6 \pm 0.2) \times 10^{-4}$, $c = (1.3 \pm 0.1) \times 10^{-1}$. $M1/\beta$ is 111.4 ± 3.6 V/cm.

Since the saturation effect is mostly present at high electric field values, a regular quadratic fit (shown with the green line in Fig. 4.10 for UB and in Fig. 4.11 for LB) was fitted to the data. The values for $M1/\beta$ are extracted from the regular quadratic fit up to 500 V/cm for both UB and LB.

4.4.2 To get the right unit for $M1/\beta$

The normalized transition amplitudes for both hyperfine branches, extracted from the quadratic fits for the upper and lower bursts, are as follows:

$$\left(\frac{M1}{\beta}\right)_{\text{UB}} = \frac{M1_{\text{rel}} - M1_{hf}}{\beta}, \quad \left(\frac{M1}{\beta}\right)_{\text{LB}} = \frac{M1_{\text{rel}} + M1_{hf}}{\beta}. \quad (4.15)$$

Summing and subtracting these expressions gives the following expressions:

$$\left(\frac{M1}{\beta}\right)_{\text{UB}} + \left(\frac{M1}{\beta}\right)_{\text{LB}} = \frac{2M1_{\text{rel}}}{\beta}, \quad \left(\frac{M1}{\beta}\right)_{\text{LB}} - \left(\frac{M1}{\beta}\right)_{\text{UB}} = \frac{2M1_{hf}}{\beta}. \quad (4.16)$$

Using known $M1_{hf}$, one solves the difference equation to determine β (in units of a_0^3), then uses the sum equation to obtain $M1_{\text{rel}}$ (in units of μ_B/c). Theoretical numbers are typically quoted in the literature in atomic units. Since β arises from the induced Stark effect, the unit of electric field (V/cm) is carried over to M/β . To obtain $M1_{\text{hf}}$ or $M1_{\text{rel}}$ in units of μ_B/c , all quantities must have consistent units:

$$M1 \left(\frac{\mu_B}{c}\right) = \left(\frac{M1}{\beta}\right)_{\text{meas}} \left(\frac{\text{V}}{\text{cm}}\right) \times \beta (a_0^3).$$

The electric field unit V/cm is converted to V/m by: $1 \frac{\text{V}}{\text{cm}} = 100 \frac{\text{V}}{\text{m}}$. The electric field must then be converted from SI units to atomic units (AU) by multiplying by the conversion factor:

$$E_{\text{SI} \rightarrow \text{AU}} \left(\frac{\text{V}}{\text{m}}\right) = 1.9446906 \times 10^{-12} \cdot \frac{\hbar^2}{e a_0^3 m_e},$$

where $\hbar$ is the reduced Planck constant, e the elementary charge, a_0 the Bohr radius, and m_e the electron mass. Finally, the Bohr magneton μ_B is defined as:

$$\mu_B = \frac{e\hbar}{2m_e} = \frac{\alpha}{2}, [65]$$

where α is the fine-structure constant. To convert the SI magnetic moment obtained from $(M1/\beta)_{\text{meas}}, E_{\text{SI}}, \beta$ into atomic units, we note that this product has the dimensions of a magnetic dipole moment in SI units. We then divide by the SI value of the reduced M1 matrix element, $M1_{\text{redME2SIamp}}$, which converts SI magnetic moments into atomic-unit magnetic moments: $M1_{\text{AU}} = \frac{M1_{\text{SI}}}{M1_{\text{redME2SIamp}}}$. In atomic units, $e = \hbar = m_e = 1$, so the atomic unit of magnetic moment is $\mu_{\text{AU}} = \mu_B$, giving $M1_{\text{AU}} \rightarrow M1(\mu_B)$. Finally, because the reduced M1 operator is defined relativistically as $M1 = \mu/c$, the natural unit of the final result is μ_B/c , matching the convention used in the numerical evaluation. We experimentally obtain $M1/\beta$ for both hyperfine components, the average of both will give us $M1_{\text{rel}}$ and the difference will give $M1_{\text{hf}}$. Once $M1_{\text{hf}}$ is known, we can extract β . The peak area estimated from both models differed moderately, with the absolute difference between their values at specific voltages relative to the average peak area of $\sim 15\%$ for UB and $\sim 13\%$ for LB.

This prompts an examination of the $M1/\beta$ value for both models, as presented in the Table 4.1. The values of $M1/\beta$ for UB in Table 4.1 exhibit strong internal consistency relative to their mean, with an average absolute deviation of approximately 2.2%. In contrast, the LB dataset shows a higher degree of variability, with an average absolute deviation of about 3.5% when assessing relative agreement between the two models. The discrepancy arises because β is the parameter that defines the curvature of the parabola, which governs how the parabola behaves near its minimum. As mentioned above, the saturation comes into affect at higher electric fields, so we considered obtaining the values for β and M1 by fitting the regular quadratic function up to the electric field values where the saturation is negligible, that is 500 V/cm. For a sanity check, we estimated β by fixing the M1 offset as the weighted

average of all M1 values at 0 V for the respective $\Delta F = \pm 1$ measurements. Let us see how these numbers work towards the final results for $M1_{rel}$, $M1_{hf}$, β in Table 4.1.

Table 4.1: $M1/\beta$ values for UB and LB for 2024.

#	Fitted model	$\frac{M1}{\beta}_{UB}$ (V/cm)	$\frac{M1}{\beta}_{LB}$ (V/cm)	Remarks
1	Voigt	163.8 ± 8.1	111.4 ± 3.6	Fit with quadratic saturation up to 2500 V/cm
2	Voigt	163.2 ± 10.3	108.6 ± 4.4	Fit with regular quadratic up to 500 V/cm
3	Voigt	166.1 ± 10.9	108.8 ± 4.3	Anchored M1 and got β from regular quadratic fit ³
4	Lorentzian	155.9 ± 6.7	120.0 ± 3.4	Fit with quadratic saturation, up to 2500 V/cm
5	Lorentzian	155.2 ± 8.9	116.6 ± 3.0	Fit with regular quadratic up to 500 V/cm
6	Lorentzian	158.3 ± 9.3	115.7 ± 3.8	Anchored M1 and got β from regular quadratic fit ⁴

The values for $M1_{hf}$, β , $M1_{rel}$ for each model agree internally with each other (#1 – #3 and #4 – #6, well within their error bars $\leq 0.3\sigma$. On comparing the output from two models with each other, show a discrepancy of 1.0σ for $M1_{hf}$ and of 1.6σ for $M1_{rel}$ which is due to how well these two models fit the data (as mentioned above). The reduced χ^2 performs well for the Voigt model, so we will further discuss the comparison of our result #2 in Table 4.2 with theory.

³Anchored the M1 offset as 0.029 ± 0.001 for UB and 0.019 ± 0.001 for LB by taking weighted average of M1 from rate and got value of β^2 as in parameter ‘e’ from regular quadratic fit up to 500 V/cm in Fig.4.10 for UB and 4.11 for LB

⁴Anchored the M1 offset as 0.031 ± 0.001 for UB and 0.021 ± 0.001 for LB by taking weighted average of M1 from rate and got value of β^2 as in parameter ‘e’ from regular quadratic fit up to 500 V/cm in Fig.C.2 for UB and C.4 for LB

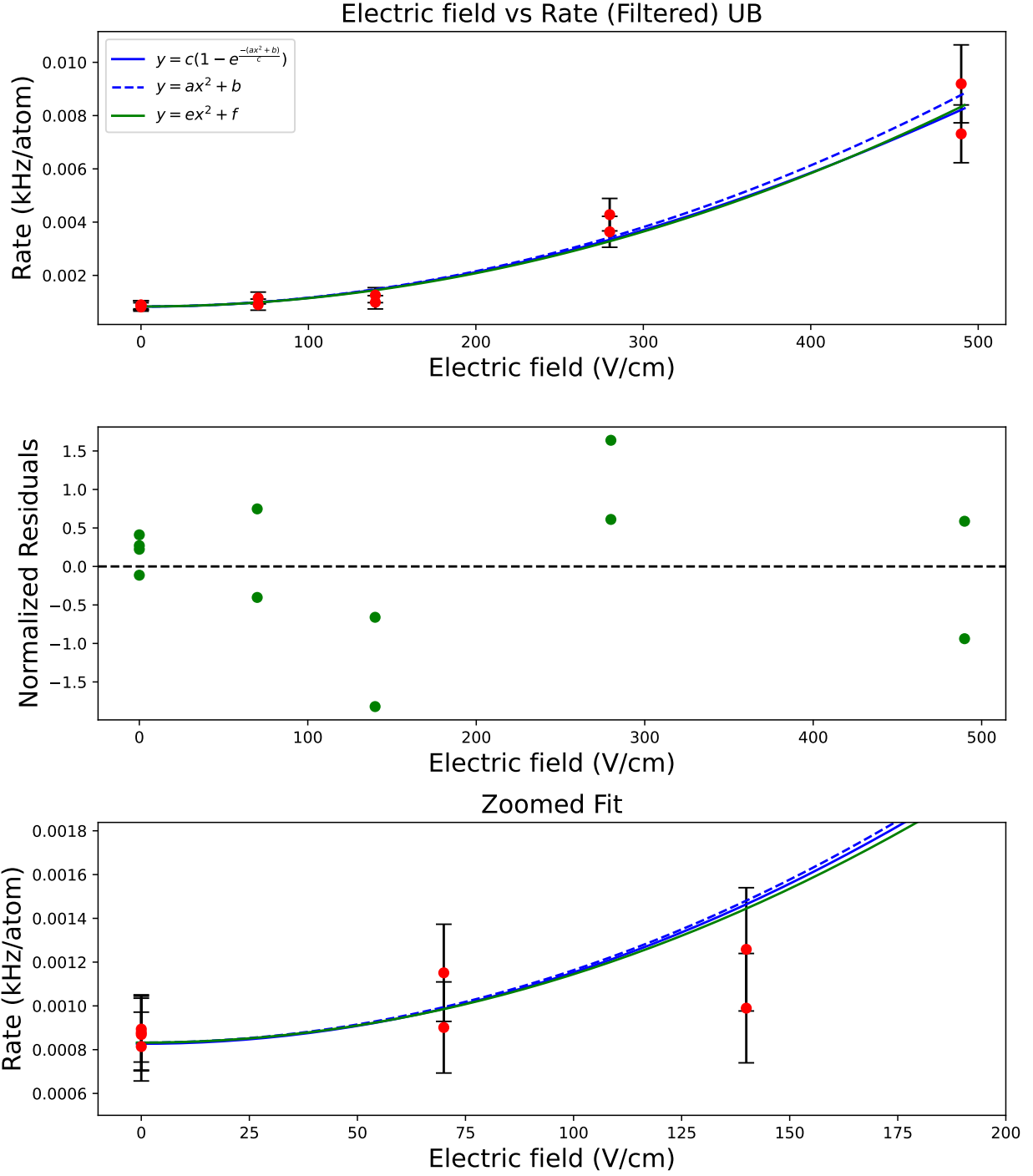


Figure 4.10: Quadratic dependence of transition rate as a function of electric field on $\Delta F = -1$, $7S_{1/2}(F = 5) \rightarrow 8S_{1/2}(F' = 4)$ with Voigt model. The values of fitted parameters from solid line is $a = (3.3 \pm 0.8) \times 10^{-9}$, $b = (8.3 \pm 0.7) \times 10^{-4}$, $c = (6.6 \pm 21.4) \times 10^{-2}$, $e = (3.1 \pm 0.3) \times 10^{-8}$, $f = (8.3 \pm 0.7) \times 10^{-4}$. $M1/\beta$ is 163.2 ± 10.3 V/cm.

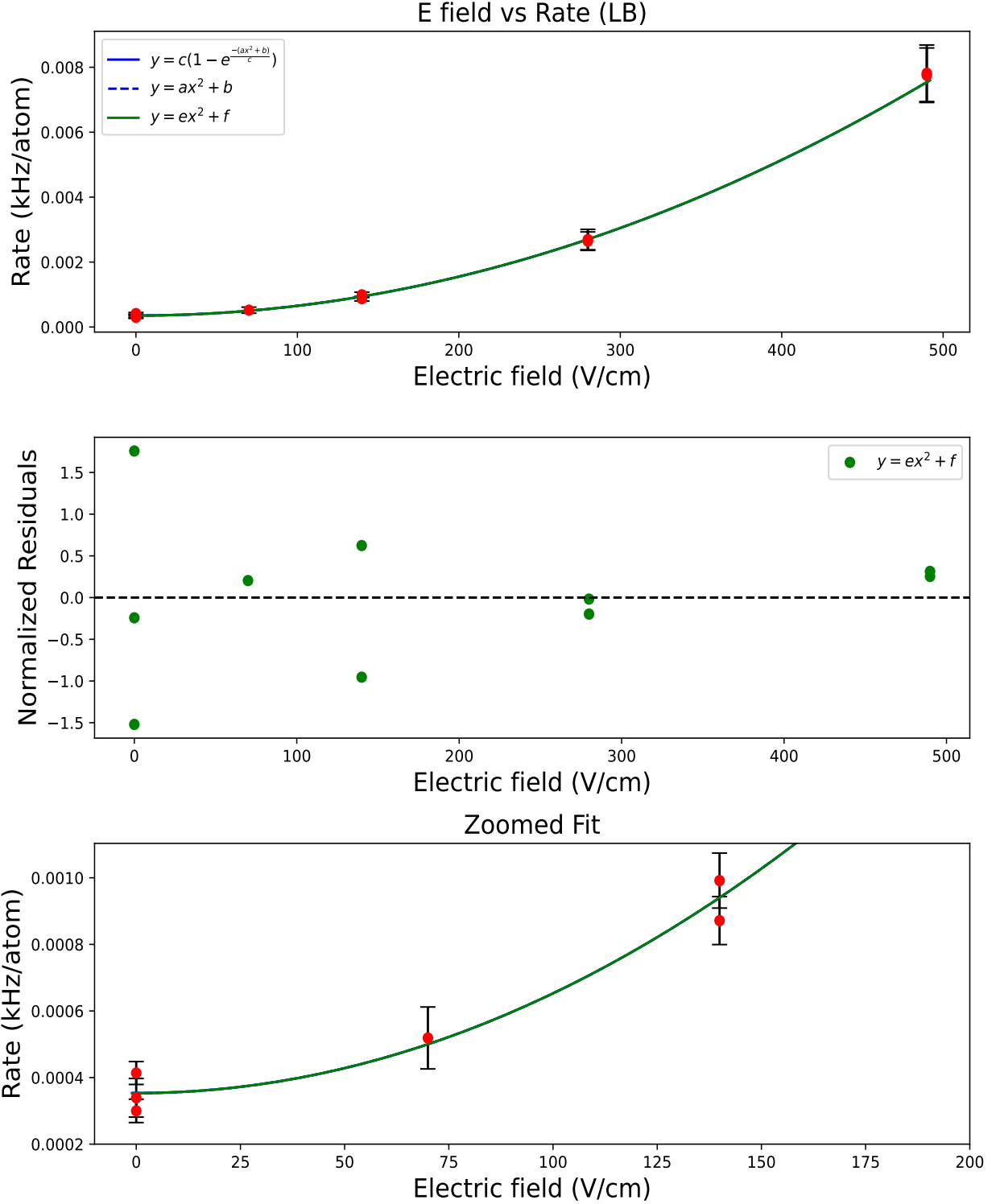


Figure 4.11: Quadratic dependence of transition rate as a function of electric field on $\Delta F = -1$, $7S_{1/2}(F = 5) \rightarrow 8S_{1/2}(F' = 4)$ with Voigt model. The values of fitted parameters from solid line is $a = (3.0 \pm 0.2) \times 10^{-9}$, $b = (3.5 \pm 0.2) \times 10^{-4}$, $c = (5.6 \pm (2.5 \times 10^2)) \times 10^1$, $e = (3.0 \pm 0.2) \times 10^{-8}$, $f = (3.4 \pm 0.2) \times 10^{-4}$. $M1/\beta$ is 108.6 ± 4.4 V/cm.

Table 4.2: Outcome of $M1/\beta$ measurement for 2024.

#	$M1_{hf} \times 10^{-5}(\mu_B/c)$	$\beta(a_0^3)$	$M1_{rel} \times 10^{-5}(\mu_B/c)$
Remarks	Used theo. value of β	Used theo. value of $M1_{hf}$	Used theo. value of β
1	-10.4 ± -1.7	24.7 ± 4.1	133.5 ± 4.4
2	-10.8 ± -2.2	23.7 ± 4.8	131.8 ± 5.5
3	-9.8 ± -1.7	26.2 ± 4.7	129.5 ± 4.4
4	-7.1 ± -1.5	36.0 ± 7.6	133.8 ± 3.9
5	-7.6 ± -1.9	33.6 ± 8.2	131.8 ± 4.7
6	-8.4 ± -2.0	30.4 ± 7.2	132.9 ± 5.0
Theoretical values	-3.450	74.3 ± 0.7	137.4

Using the theoretical value of $M1_{hf}$ and the measured data, we determined the following. $\beta = 23.7(4.8) a_0^3$. The theoretical prediction for β is $74.3(7) a_0^3$, which makes our result for β differ from the theoretical value by a factor of 3. As already mentioned, in Cs the theoretical and measured values agree very well with each other [58]. It is therefore very unlikely that, in our case for Fr, the theory is off by such a large factor. We then calculated $M1_{rel}$ using Eq. 4.16, taking the theoretical value of β . As shown in the fourth column of Table 4.2, the resulting $M1_{rel}$ is quite close to the theoretical prediction. On the other hand, if we use the theoretical value of β from theory and use it to estimate $M1_{hf}$, the result is in poor agreement with the theoretically calculated value of $M1_{hf}$. We may need to question the validity of the measurement scheme.

Comparison of $M1/\beta$ for 2021, 2023 and 2024: Figure 4.12 shows a comparison of the measured values of $M1/\beta$ for three different beamtimes. The saturated quadratic curve was fitted to the data to estimate the value of $M1/\beta$.

In 2021, we were only able to measure the $7S_{1/2}(F=5) \rightarrow 8S_{1/2}(F'=4)$ transition [33] because the cyclotron broke down before we could measure the other transition. The transition rate was measured using a photomultiplier tube (PMT). This is the same transition for which we used the lower-burst (LB) configuration in both 2023 and 2024. Because only this transition was accessible in 2021, we used the theoretical values of $M1_{hf}$ and β to predict

$M1/\beta$ for the transition on which we would later use the UB configuration.

In **2023**, we measured $M1/\beta$ for both UB and LB using the full PBC power, as well as using a reduced PBC power (by a factor of ~ 8), the latter shown as “2023*” in Table 4.3. During this beamtime, the PBC did not have sidebands, and the dominant source of uncertainty arose from the atom number.

In **2024**, we measured $M1/\beta$ on both $\Delta F = \pm 1$ transitions, and the PBC included sidebands at one FSR. The UB measurements from all years agree well with one another, with all values lying within $< 0.3\sigma$ of the mean value of 165.4. For LB, the measurements are mutually consistent only at the $\sim 2\sigma$ level, giving a weighted mean of 119.3. Although the LB values show poorer internal agreement compared to UB, they still overlap within their quoted uncertainties.

Table 4.3: Outcome of $M1/\beta$ measurement for 2021, 2023, 2024.

#	$M1_{hf} \times 10^{-5}(\mu_B/c)$	$\beta(a_0^3)$	$M1_{rel} \times 10^{-5}(\mu_B/c)$
Remarks	Used theo. value of β	Used theo. value of $M1_{hf}$	Used theo. value of β
2021 ⁵	-3.450 (theo.)	74.3 ± 0.7 (theo.)	152.0 ± 8.4
2023	-5.2 ± 2.2	49.6 ± 20.8	140.9 ± 5.5
2023*	-6.0 ± 3.0	42.8 ± 21.0	142.0 ± 7.3
2024	-9.9 ± 1.4	26.0 ± 3.6	132.2 ± 3.6
Theoretical values	-3.450	74.3 ± 0.7	137.4

⁵Used theoretical value of β and $M1_{hf}$ to evaluate the value of $M1_{rel}$.

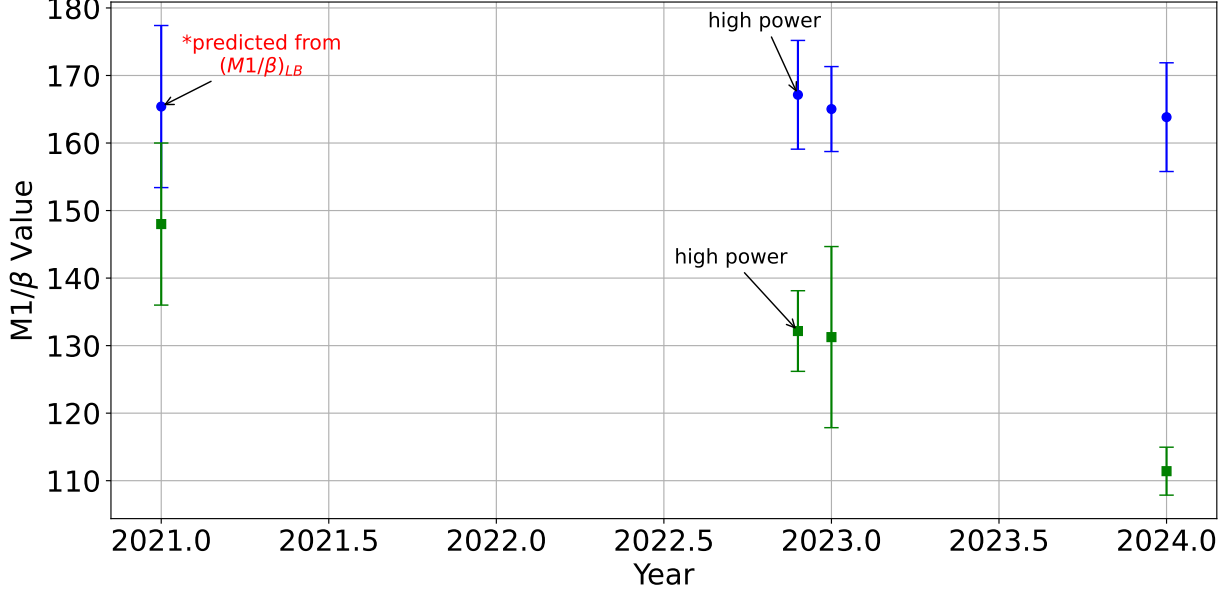


Figure 4.12: Comparison of $M1/\beta$ for 2021, 2023, 2024. In 2021, the predicted value shown with * is for UB. The ‘green’ dots show the measurement for the LB and the ‘blue’ dots show the measurement of $M1/\beta$ with UB. For 2023, we measured $M1/\beta$ for high and low blue power.

4.5 Discussion

The 2024 results for $M1/\beta$ deviate from the theoretical predictions in a way that warrants further examination. There are several possible avenues to investigate this discrepancy: 1) to examine whether the theoretical predictions for these quantities require revision, 2) to evaluate the possible contribution of the quadrupole transition amplitude to $M1_{hf}$, and 3) to determine a potential systematic issue of the measurement that calls for closer investigation and could lead to an improved measurement scheme.

The theoretical value of β is generally considered reliable [57], over 99% of its contribution comes from valence excited states calculated using high-accuracy single-double (SD) wavefunctions, but its evaluation remains sensitive to small uncertainties because the $np_{1/2}$ and $np_{3/2}$ terms partially cancel. Overall, the theoretical uncertainty in β is estimated to be $\sim 1\%$, driven mainly by correlation effects and limitations in scaled SD data for the key contributing transitions. The Boulder group carried out an experiment that evaluated the

violation of atomic parity in cesium, resulting in a value of $\beta = 27.2(4) a_0^3$ [35], which agrees very well with the theoretical value of β predicted by Safronova, $27.11(22) a_0^3$ [57]. More recently, a team at Purdue University, pursuing an independent reevaluation of parity violation in cesium, reported a new measurement of $\beta = 27.139(42) a_0^3$ [66].

As mentioned above, β can be measured in two ways: one via α/β measurement (which uses a theoretically calculated value of α) and the other via $M1/\beta$. For cesium, the discrepancy between the values of β obtained using these two methods is 0.68% [66], corresponding to 2.8σ . In addition, the theoretical estimate of $M1_{\text{hf}}$, arising from the off-diagonal hyperfine interaction, is expected to be highly reliable. Given the well-known hyperfine splittings, the amplitude $M1_{\text{hf}}$ should be determined with good accuracy, making a significant theoretical error unlikely.

Beyond the hyperfine-induced $M1_{\text{hf}}$ contribution, it is also important to consider additional higher order amplitudes, such as the electric quadrupole transition amplitude (E2). It is induced by hyperfine interaction mixing of the $nS_{1/2}$ state with the $n'D_{3/2}$ and $n'D_{5/2}$ states. Since these states have the same parity, requiring $(-1)^l$ to change only by $\Delta l = 0, \pm 2$, this mixing enables the otherwise forbidden $S \rightarrow S$ E2 transition. In addition to this, the atomic quadrupole moment arising from the interaction between atomic electrons and the nuclear moments contributes to E2, which breaks the spherical symmetry of the ground-state electron cloud in alkali-metal atoms. The transition operator[60] for $7S - 8S$ transition is

$$T_0(nS - n'S) = a_1 S \cdot \epsilon \times k + ia_2 (S \times I) \cdot (\epsilon \times k) + ia_3 [(S \cdot \epsilon) \cdot (I \cdot k) + (S \cdot k) \cdot (I \cdot \epsilon)] + ia_4 [(I \cdot \epsilon) \cdot (I \cdot k) + (I \cdot k) \cdot (I \cdot \epsilon)], \quad (4.17)$$

where S and I are the electronic and nuclear spin operators, $\mathbf{k}$ and $\boldsymbol{\epsilon}$ are the momentum and polarization of the photon involved in the transition, the a_i values are real and experimentally determined. The atomic quadrupole moment consists of contributions from the electronic cloud, the nucleus, and hyperfine interaction (HFI) corrections. Equation(4.17) has the first term as the normal M1 amplitude, the second term is the HFI-induced M1 amplitude, the

third term is the electric quadrupole amplitude, and the last term is the interaction of the electron with the electric quadrupole of the nucleus. The contribution from E2 could be significant. There is a very limited literature on estimates of E2 amplitude in ^{211}Fr . The individual estimates of these a_i terms are not easy to calculate [60], so does the term a_3 is. We do not have any literature value of a_3 for Fr. However, the ratio a_4/a_3 [60] can be evaluated because the nuclear quadrupole moment is known.

$$\frac{a_4}{a_3} = \frac{2Q/a_0^2}{I(2I-1)g_S g_I \alpha^2}, \quad (4.18)$$

where g_S, g_I are the electronic and nuclear g-factors in Bohr magnetons ($1\mu_N = I \cdot \frac{m_e}{m_p} \mu_B$), a_0 is Bohr magneton, Q is the Quadrupole moment of Fr nucleus ($Q = -0.19(3)$ barn [67]). Although the nuclear quadrupole moment in ^{211}Fr is very large in comparison to Cs [67], the a_4 term contribution is 137 times smaller than a_3 . Only a_3 term has to be considered.

In heavy alkali-metal atoms, the HFI dominates because the electronic contribution vanishes due to spherical symmetry. The correlation effects increase the hyperfine-induced quadrupole moments in heavy atoms by up to two orders of magnitude [68] due to significant contribution from core-screening effect. When the magnetic-dipole HFI term dominates, for nuclear spin $I \geq \frac{3}{2}$, the ratio of HFI-induced quadrupole moments between the two hyperfine levels is given by [68]:

$$\frac{Q_{F=I+1/2}^{\text{HFI}}}{Q_{F=I-1/2}^{\text{HFI}}} = -\frac{2I(2I+1)}{(2I-1)(2I-2)}. \quad (4.19)$$

Following Eq. 4.19, the upper hyperfine ground state possesses a larger electric quadrupole tensor moment, it is more sensitive to electric field gradients. Consequently, transitions involving this state, with which LB is associated with, may experience larger E2-induced energy shifts. The potential contribution of the E2 interaction warrants further quantitative analysis; at present, it is not possible to determine whether it accounts for the observed discrepancy, although it remains a plausible candidate. Notably, according to Bouchiat, when linearly polarized light is used and the magnetic sublevels are equally populated, the

electric quadrupole (E2) amplitude does not contribute to the transition. Our result for $M1_{hf}$ is approximately 3.1 times larger than $(M1_{hf})_{\text{theo}}$, assuming the theoretical β is accurate. The a_3 term needs to be estimated by theorists to better understand its contribution to Equation (4.17). These observations underscore the importance of accurately determining a_3 in order to fully explain the enhancement observed in $M1_{hf}$ for heavy atoms like ^{211}Fr . A theoretical estimate of this term could help reconcile the remaining differences between theoretical predictions and experimental results. In the following section, we will discuss a few other possible effects in the following section that might have played a role in seeing this discrepancy in $M1/\beta$ value with LB.

Polarization Misalignment Affects β and $M1$: From Equation (1.42), it follows that $\vec{E}_S \times \vec{\epsilon}$ is maximized when $\vec{E}_S$ and $\vec{\epsilon}$ are perpendicular, i.e., when $\theta = 90^\circ$, since the magnitude is given by $|\vec{E}_S||\vec{\epsilon}|\sin\theta$. The alignment and linear polarization state of blue beam are defined by the Glan-Thompson prism just before the PBC input. The M1 amplitude is independent of the electric field and therefore unaffected by misalignment of $\vec{\epsilon} \perp \vec{E}$, whereas β is sensitive to such misalignment. Any deviation from perfect perpendicularity reduces the transition rate for β , potentially lowering the observed rate at higher electric field values compared to expectations. From Equation (4.6), in an extreme case where β is halved, the equation yields a wider (less steep) parabola, increasing the $M1/\beta$ ratio by double. However, the observed discrepancy in our result shows the opposite: our measured $M1/\beta$ ratio is lower than expected, inconsistent with the theoretical prediction for $M1_{hf}$ and β .

A potential factor affecting the polarization of linearly polarized blue light is birefringence, which can originate from two main sources: the viewport window on the science chamber (SC), through which the blue light enters the cavity, and the coatings on the PBC mirrors. A common cause of birefringence is mechanical stress in the mirror substrates. An applied voltage to the piezos can further modify the birefringence by altering the stress state or orientation of the PBC mirror coatings. The outer surfaces of both cavity mirrors are anti-reflection coated to minimize etalon effects between the SC viewport substrate and the

mirrors, which helps reduce unwanted optical interference. Assuming that $\vec{E}_S$, $\vec{B}$, and $\vec{\epsilon}$ are aligned along the x -, y -, and z -axes respectively, the incoming blue light is polarized along $\vec{\epsilon} = \hat{z}$, and its wavevector $\vec{k}$ reverses direction (from $\vec{k}$ to $-\vec{k}$) upon retro-reflection inside the cavity. The interference term between the Stark-induced amplitude and the magnetic dipole (M1) amplitude ideally vanishes under these conditions due to the symmetry of the setup. A deviation from this ideal condition arises when birefringence in the PBC mirrors modifies the polarization state of the reflected light. This modification can prevent full cancellation of the interference term, leading to a residual Stark-M1 contribution. The resulting systematic effect can be significant if the birefringence introduces even a small phase shift or rotation in polarization. The Equation (4.5) can be rewritten more specifically as

$$R_{7S \rightarrow 8S} \propto \left| \left[\beta(\vec{E}_S \times \vec{\epsilon}) \cdot \langle F' | \vec{\sigma} | F \rangle + \langle F' | M1(\vec{k} \times \vec{\epsilon}) \cdot \vec{\sigma} | F \rangle \right] \right|^2, \quad (4.20)$$

where $\vec{\epsilon} = \epsilon_x \hat{x} + \epsilon_y \hat{y} + \epsilon_z \hat{z}$ represents the electric field of the light, under the effect of birefringence, inside the PBC. The spin operator $\vec{\sigma}$ has already been projected on to the states, the Pauli matrices disappear into $\langle F' | \vec{\sigma} | F \rangle$, and the cross product term becomes

$$\vec{E}_S \times \vec{\epsilon} = E_S \hat{x} \times (\epsilon_x \hat{x} + \epsilon_y \hat{y} + \epsilon_z \hat{z}) = E_S(\epsilon_z \hat{y} - \epsilon_y \hat{z}), \quad (4.21)$$

and the M1 term becomes

$$A_{M1} = M1 \cdot k(\epsilon_x \sigma_y - \epsilon_y \sigma_x). \quad (4.22)$$

The substitution of these expressions into the interference term, combined with the use of the Pauli matrix identity

$$\sigma_i \sigma_j = \delta_{ij} I + i \epsilon_{ijk} \sigma_k, \quad (4.23)$$

yields

$$R_{7S \rightarrow 8S} \propto 2\beta E_S M1 \cdot k \left[\epsilon_z^* \epsilon_x + \text{Im} \left(-\epsilon_z^* \epsilon_y \sigma_z + \epsilon_y^* \epsilon_x \sigma_x - \epsilon_y^* \epsilon_y \sigma_y \right) \right]. \quad (4.24)$$

A nonzero interference term like this becomes sensitive to the relative phases between ϵ_z and ϵ_x, ϵ_y , highlighting how the polarization distortion could impact the measurement. The overall interference is also influenced by the angle of incidence of the beam and any change in polarization upon reflection, both of which are affected by birefringence in optical components. We observed that a 10° change in the Glan-Thompson angle varies the amount of light entering the cavity by less than $(4 \pm 0.1)\%$. The optimal Glan-Thompson setting to obtain vertical polarization of the blue beam in the PBC is within $\pm 2^\circ$. We measured the polarization of the output from the locked PBC and found that the degree of linear polarization varies by less than 0.3%, while the azimuthal angle changes by approximately 10–20%. This indicates a slight rotation of the polarization plane, likely due to birefringence in the cavity mirrors. For such polarization rotation, the interference term in Eq. 4.24 is affected. Given the sensitivity of this term to polarization, a 10–20% rotation is estimated to reduce $R_{7S \rightarrow 8S}$ by roughly 5%. We suspect that any intrinsic birefringence lies below the measurable sensitivity of our polarimeter. Considering that our cavity gain is significantly lower than that used by Wieman, the mirrors in our setup would need to exhibit substantially higher birefringence to produce a comparable effect.

4.6 Result

Using the theoretical value of β for entry 2 of Table 4.2, we obtain $M1_{\text{rel}} = (131.8 \pm 5.5) \times 10^{-5} \frac{\mu_B}{c}$, which agrees with Safronova’s theoretical prediction to within $\sim 1\sigma$. However, when considering β and $M1_{\text{hf}}$, the situation is different. Although birefringence in the cavity mirrors may introduce a small polarization rotation, it is unlikely to explain the discrepancy between the experimentally determined values of β and $M1_{\text{hf}}$ and their theoretical predictions. Because β is theoretically well constrained, we instead examined how much $M1_{\text{hf}}$ would need to deviate to reconcile the results. Using the theoretical β , the inferred value is $M1_{\text{hf}} = (-10.8 \pm 2.2) \times 10^{-5} \frac{\mu_B}{c}$, which differs from the theoretical prediction by about 3.3σ . This significant deviation suggests that the theoretical estimate of $M1_{\text{hf}}$ may need to

be revised. However, to verify the reliability of our measurement method, we will remeasure $M1/\beta$ in the same way as we did in 2024 and/or perform a new data-taking cycle in which the trapping period will be increased from 1 ms to about 2 ms. This would ensure that, if there is any magnetic-sublevel dependence of atoms in the MOT due to photon bursting, the longer trapping period can reset the MOT. Another approach would be to reduce the burst window from 400 μs to 200 μs , thereby reducing the momentum imparted to the atoms in the MOT and, consequently, minimizing any position-shift effects that could affect the overall overlap between the MOT and the blue beam. It was tried out in beamtime in January 2026, the data analysis is under process.

4.7 Summary

To summarize, the $M1/\beta$ measurement was performed on both $\Delta F = \pm 1$ transitions using the upper-burst and lower-burst detection techniques. From Fig. 4.12, it is clear that the $M1/\beta$ value for the $\Delta F = -1$ transition obtained using the lower-burst detection method requires further investigation or additional experimental checks. In contrast, the $M1/\beta$ value for the $\Delta F = +1$ transition has remained very consistent across multiple beamtimes. Using the measured $M1/\beta$ value for the $\Delta F = +1$ transition together with the theoretically predicted values of $M1_{\text{hf}}$ and β , we obtain

$$M1_{\text{rel}} = 132(6)_{\text{expt}} \times 10^{-5} \mu_B/c,$$

which agrees with Ref. [69] within 1σ . The lower-burst measurement for the $\Delta F = -1$ transition therefore warrants additional scrutiny.

CHAPTER 5

DEVELOPMENTS

The quadrupole magnetic field for MOT is always on in our present experimental setup. Since the field is spatially varying, there is no specific defined direction of magnetic field in the science chamber. This leads to having an unpredictable m_F -level dependence or an uneven distribution of atoms in MOT. As mentioned previously, for measurement of $M1/\beta$, the m_F -sublevel distribution of trapped atoms is not needed because both $M1$ and β have the same dependence on the m_F sublevel. If we assume the atoms to be equally distributed among the Zeeman sublevels, then the interference term averages to zero because of opposing contributions from symmetric and equally populated m_F sublevels. However, the distribution of trapped atoms in the MOT must be precisely known for future measurements, such as E1-M1 interference which has m_F dependence [70]. It requires spin-polarizing the atoms in the MOT which leads to a known population distribution of atoms in the Zeeman sublevels. It can be done by optically pumping the atoms in a specific $|F, m_F\rangle$ state. The angular momentum coupling coefficients for the angular dependence of the 7S-8S transition between the extreme m_F -levels of 7S and 8S are the strongest [see Appendix A]. This suggests that having the majority of trapped atoms in the extreme magnetic sublevels will yield the largest signal from the highly forbidden 7S-8S transition. This chapter covers the future plans of pumping atoms to the extreme magnetic sublevel, proposed implementation of the set-up on SC for pumping atoms, dealing with magnetic fields, and other important prospects.

Optical Pumping (OP): Kastler (1950) [71] introduced the concept of optical pumping for alkali atoms, a process that creates atomic ground-state polarization. This happens because of two factors: 1) the presence of a dark state, where atoms cannot interact with the polarized optical pumping (OP) beam due to selection rules. The presence of a dark state

plays an important role in terminating the pumping process due to $\Delta m = \pm 1, 0$ angular momentum selection rule. Because once atoms fall into the dark state, they would no longer absorb the pumping light, and hence, the pumping cycle of atoms terminates. 2) Ground-state polarization relaxes more slowly than the excited state. Because the ground state is not subjected to radiative decay or rapid spontaneous emission, polarized atoms can maintain their polarization for a longer time. In our system, two processes are used to populate a single $|F, m_F >$ hyperfine ground state: the first step is hyperfine pumping where we put atoms into the desired hyperfine level, F . The second step is Zeeman pumping where the population is transferred into a specific magnetic sublevel m_F within that hyperfine level.

How will we do it with ^{211}Fr :

- (1) We first trap atoms in a single hyperfine ground state, $7S_{1/2}(F = 5)$, following the standard francium MOT procedure described in Chapter 2.
- (2) The quadrupole magnetic field is then switched off.
- (3) At this stage, the unwanted hyperfine ground state has been depleted, leaving the atoms in the desired hyperfine ground state, though still unevenly distributed among its Zeeman sublevels. A small uniform bias magnetic field is then applied together with the optical pumping beam to define the quantization axis.
- (4) A branch of the trapping laser is frequency shifted to be resonant with the optical pumping transition $7S_{1/2}(F = 5) \rightarrow 7P_{3/2}(F' = 5)$. A σ^+ polarized light, propagating along the quantization axis, excites transitions with $\Delta m = +1$, thus transferring population towards the stretched state $7S_{1/2}(F = 5, m_F = 5)$. By repeated absorption of σ^+ light and subsequent decay via $\Delta m = 0, \pm 1$ (dashed blue lines in Fig. 5.1), the atoms in the MOT become spin-polarized within $\sim 100 \mu\text{s}$, gradually accumulating in the dark state $F = 5, m_F = 5$, where they no longer interact with the pumping light, as shown in the simulation in Fig. 5.1. The repump beam (clean-up beam shown in Fig. 5.1(a)) remains on to prevent population from accumulating in the lower hyperfine ground state and to ensure that all atoms remain in the $F = 5$ level during Zeeman pumping.

Define the quantization axis with B field: A quantization axis is defined with the presence of a dc magnetic field along with spin to which the trapped atoms will be aligned. The interaction of the external magnetic field with nuclear spin (I), the total electronic angular momentum (J) of atom and the total angular momentum of atom, F, governs the Zeeman effect on hyperfine structure. For small magnetic fields, the interaction of J·B is

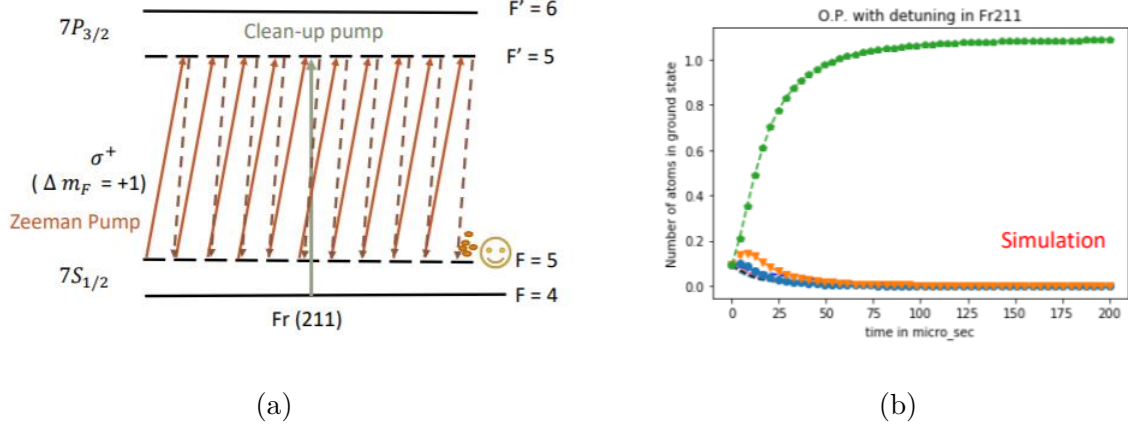


Figure 5.1: (a) shows optically pumping the atoms with σ^+ polarized light into $7S_{1/2}(F = 5, m_F = 5)$. (b) shows the simulated estimate of pumping of atoms in the extreme m_F sublevel shown by the green curve. The orange curve shows the transfer of population of via the top excited state of $7P_{3/2}$. The x-axis is in μs . The simulation shows that $\sim 100\mu s$, most of the population will be pumped into extreme m_F sublevel. The green curve shows the atoms going to $+m_F$ of $F = 5$, and the other curves with different colors show the evolution of atoms on most of the m_F sublevels shown on left side of this figure.

smaller than the I·J interaction resulting into the interaction of J·B as a small perturbation effect. This perturbation shifts the energies of ' $2m_F + 1$ ' sub-levels written below by [72]:

$$\Delta E_{nSFm_F} = g_F \mu_B B m_F, \quad (5.1)$$

where g_F is the Lande-g factor, $\mu_B = 1.399624$ MHz/G. For large magnetic fields, the interaction of J·B dominates the interaction between I·J resulting into mixing of different F states into each other. A uniform magnetic field across the atomic sample is desirable, and the optical pumping beam requires the magnetic sublevels to remain degenerate so that all m_F states can be addressed simultaneously. To satisfy all these conditions, Zeeman pumping

on the $7S_{1/2}(F = 5) \rightarrow 7P_{3/2}(F' = 5)$ transition is sufficient to produce a frequency shift of $\Delta E = (0.28 \text{ MHz/G})m_F B$. For a 5 G field, this results in adjacent magnetic sublevels being separated by approximately one-fifth of the D2 linewidth.

To define the quantization axis for polarized atoms, two steps are required: turning off the quadrupole field for the MOT and turning on a static homogeneous magnetic field to define a B field direction for OP. The timing of the full cycle of trapping and data collection cycle requires switching the B field on a sub-millisecond scale.

5.1 Coils to play with

The quadrupole coil for the MOT has a square cross-section with 0.245" side, a hollow core with a radius of approximately 0.055 m and 48 turns. It has a glass-fiber coating. It has an inductance of about 162 μH and a resistance of 31 m Ω . The self inductance of a circular coil made up of rectangular cross section conductor using Grover's formula can be estimated below [73], p 136:

$$\mathcal{L} = 4\pi a N^2 \left[\log \left(\frac{8a}{R} \right) \cdot \left(1 + \frac{3R^2}{16a^2} \right) - \left(2 + \frac{R^2}{16a^2} \right) \right], \quad (5.2)$$

where R is the geometrical mean distance of the circular coil, is approximately proportional to $0.2235(\alpha + \beta)$ where α, β is the length and breadth of rectangle. a is the mean radius of turns, and N is the total number of turns.

The diameter of bias coil (with insulator coating) is ~ 0.044 " and gauge of coil (AWG) is 18 for all bias coils. With no original record of the number of turns and the coils tightly wound on the frame, making it hard to count, it was necessary to estimate the number of turns. I estimated the number of turns by measuring the resistance. The inductance was measured with two methods: 1) by using an inductance meter and 2) by determining the resonant frequency, f , from an LC tank circuit. The inductance, $\mathcal{L}$, was then calculated using the

Table 5.1: Properties of Bias Coils on Science Chamber

Coils	Dimensions (L(cm) $\times$ B(cm)), Spacing (cm)	Resistance, R (Ω)	Inductance, $\mathcal{L}$ (mH)	Estimated number of turns
(X) Up(Down) coil	42.5 $\times$ 30.5, 32.5	3.0 (3.2)	8.7 (8.5)	85 - 95
(Y) Left(Right) coil	34 $\times$ 29, 40.5	1.1 (1.1)	1.5 (1.5)	25 - 35
(Z) Front(Back) coil	40.5 $\times$ 32.5, 34	2.4 (2.5)	5.8 (5.9)	75 - 85

formula.

$$f = \frac{1}{2\pi\sqrt{\mathcal{L}C}}, \quad (5.3)$$

where C is a known capacitor and f is the measured resonant frequency. I measured the coil resistance, R , and used the standard resistance per unit length of 18 AWG wire, denoted as R'/L' (in $\text{m}\Omega/\text{m}$), to calculate the total length of wire in the coil:

$$\text{Total length of coil} = \frac{R \cdot L'}{R'}. \quad (5.4)$$

Assuming the wire loops are rectangular with vertical length L and horizontal base B , the number of turns, N , can be estimated as:

$$N = \frac{\text{Total length of coil}}{2(L + B)}. \quad (5.5)$$

The estimated number of turns is shown in the last column of the table above. I have emphasized the importance of precisely controlling the coils, the inductive load, by switching them on and off at sub-millisecond timescales. It is crucial for future experiments, such as the α/β measurement that requires optical pumping, to modulate the DC magnetic field throughout the experimental cycle. In particular, the field must be switched on and off dynamically at specific stages, with optical pumping achieved by adjusting the current through the coils at the appropriate times.

5.1.1 Switching off the B field

According to Biot-Savart's law, the magnetic field B is directly related to the current I flowing through the coil. When the magnetic flux through the coils changes over time, it induces

a voltage $V(t)$, as described by Faraday’s law of induction. This voltage causes a current to flow in the closed loop. The direction of the induced current opposes the change in magnetic flux, as stated by Lenz’s law, and is given by $V(t) = -\mathcal{L}\frac{dI}{dt}$. Inductance $\mathcal{L}$ depends on the coil’s properties, such as geometry, gauge, and number of turns (shown in equation (5.2)), which in turn affect the induced voltage and the rate of current change, $\frac{dI}{dt} \propto \frac{dB}{dt}$.

The inductance $\mathcal{L}$ of our bias coils is quite large, where the decay time of current I is measured to be around sub-ms, which is too long for operations like trapping and optical pumping. $\mathcal{L}$ can be reduced by decreasing the number of turns, bringing the coils closer together, or reducing their dimensions and gauge. However, the coils cannot be placed much closer than their current spacing or made significantly smaller, except by placing them inside the chamber. This option, as shown in Fig. 5.5 with green and blue coils, is not feasible due to tight geometric constraints. There is a risk of interference with the 1” trap beams and the base of the power buildup cavity. In addition, mounting the coils inside the chamber and water-cooling them would be challenging. In addition to this, the science chamber is slightly metallic and the presence of copper gaskets creates some closed paths for current circulation. These currents, known as eddy currents [74], are distributed throughout the science chamber because the thick and large stainless steel walls of the chamber and the copper gaskets contribute similarly, and can interfere with the switching of the magnetic field, either by causing imprecise timing or gradient strengths.

We decided to control the switching of the existing coils by replacing the ‘Agilent 3614’ and ‘Kepco’ power supplies for the bias and quadrupole coils, respectively, as their slow frequency response made them unsuitable for the sub-millisecond switching needed for optical pumping (OP). The suitable replacement is Matsusada supplies that operate as ultra-high-speed bipolar power supplies (bipolar amplifiers) operating from DC to 200 kHz. The DOS25-80 model provides ± 25 V, ± 80 A, and a maximum power of 2 kW to drive each quadrupole field, while the DOS60-33 model provides ± 60 V, ± 33.3 A, and 2 kW to control a pair of bias coils.

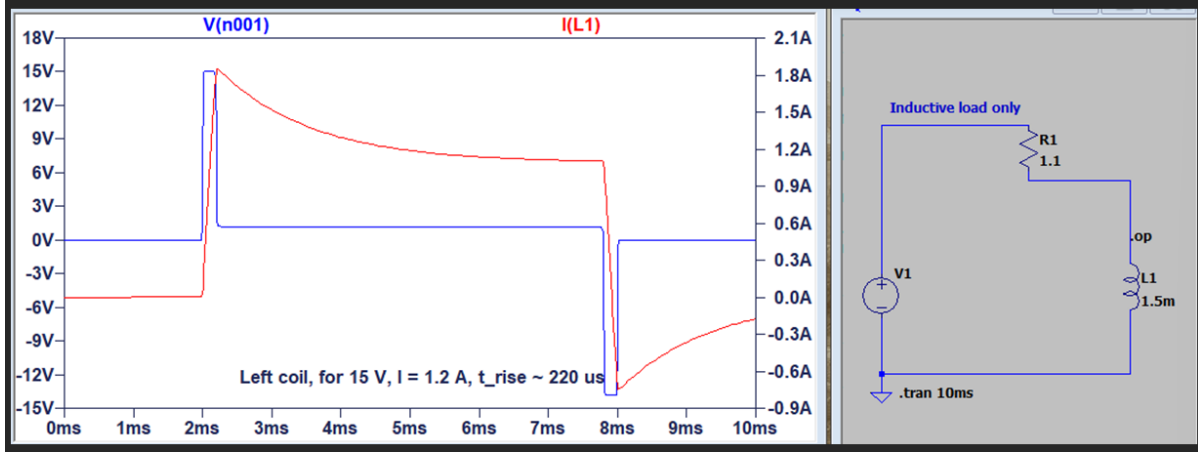


Figure 5.2: Estimation of switching behavior of bias coil with LTSpice by drawing a mock diagram of the RL circuit free from any eddy current. We load a data file with the desired waveform to estimate the real-time behavior of RL circuit made up of bias coil and quadrupole coil. The estimated response and the current flowing in the circuit are in moderate agreement with the LTSpice models. However, in real systems, the presence of eddy currents affect the switching behavior.

To power the Matsusada power supplies which were bought from Japan, they were first certified by a third-party company (Intertek) to ensure that they meet the Canadian Standards Association (CSA) requirements. The DOS 25V-80A power supply is connected to the quadrupole coils using gauge 4 cables and lug connectors. The DOS 60V-33.3A power supply will be used for the gauge 18 bias coils, which typically draw around 15 A. To comply with TRIUMF safety standards, a protective measure is required to prevent the current from exceeding 15 A. Therefore, a slow-blow fuse rated for 10 A will be installed to blow if the current through the bias coils exceeds 10 A, ensuring safety by preventing over-current conditions. The Matsusada power supplies will operate in constant voltage mode, using a custom arbitrary waveform to compensate for delays in current rise and decay in the coils controlling the magnetic field (B-field). The waveform is loaded onto an SRS345 function generator via the Arbitrary Waveform Composer Software. The output from the function generator is connected to the Matsusada power supply's programmable input channel using a BNC cable. To counteract eddy currents, a brief negative overshoot is applied near the polarity inflection point or B-field turn-off. Since eddy currents are proportional to the rate



Figure 5.3: shows the switching of the quadrupole coil with DOS 25 V, 80 A. The shaped input signal with an overshoot and undershoot is given to the Matsusada. The exponential feature in the input signal is to counteract the exponential response due to eddy currents. The blue waveform in this figure is the arbitrary waveform, and the yellow curve is the current flowing through the quadrupole coil observed from the monitor channel of the power supply.

of current change in the coil, the overall decay curve exhibits a multi-exponential response [75], influenced by both eddy currents and coil inductance. The waveform's undershoot, along with an exponential adjustment, counters the exponential response during switching. The goal is to achieve ideal square-wave switching of the B-field, as delays would disrupt the polarization state during optical pumping. To achieve this, we adjust the exponential component of the waveform.

The top waveform (blue) in Fig. 5.3 has an overshoot close to the maximum allowed voltage (25 V and 60 V). The width of the overshoot determines the switching time, after which the voltage drops below zero and then increases exponentially. This exponential rise compensates for eddy currents, stabilizing the voltage at a DC value to maintain the current. To turn off the current, the voltage drops to the opposite extreme (-25 V, -60 V), quickly

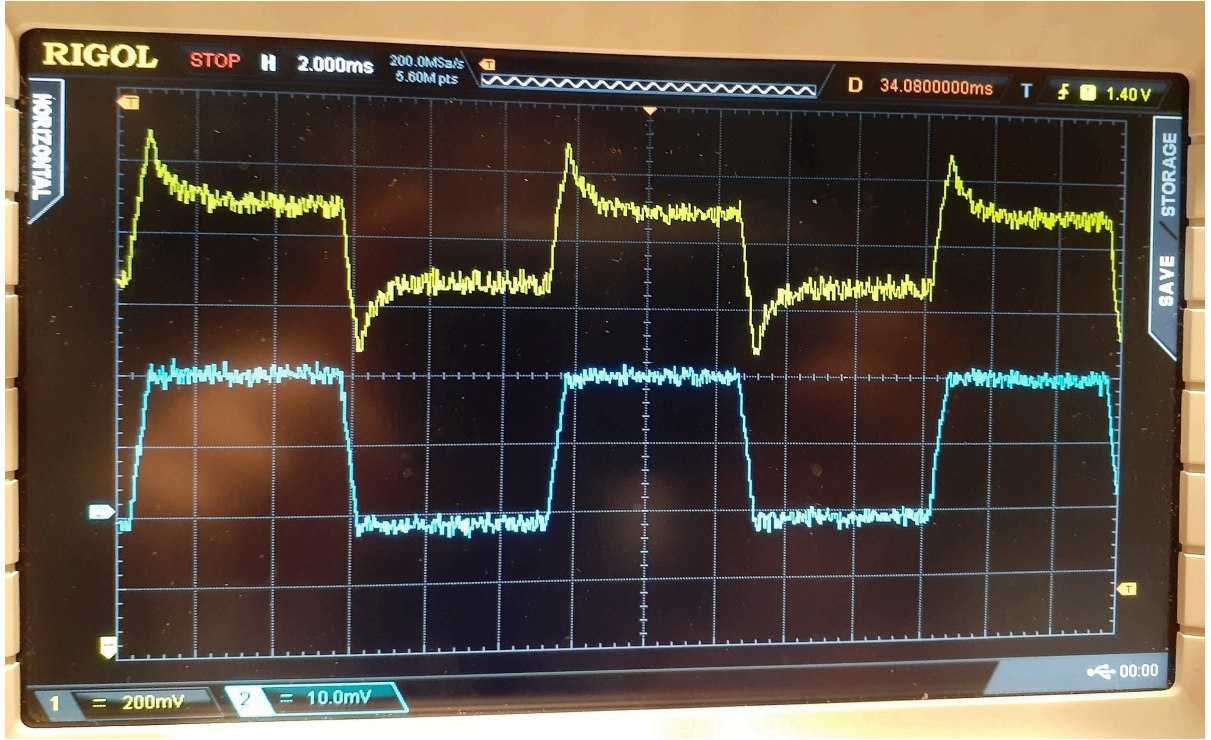


Figure 5.4: shows the response of the magnetic field when the arbitrary waveform, as shown in previous figure, goes through the quadrupole coil. The yellow curve shows the magnetic field which follows the switching behavior of current (shown in previous figure) in the center of the coil. The blue curve shows the magnetic field at the approximate position of MOT when measured on the outer side of the science chamber. It can be seen that switching action is fast on the order of $500 \mu s$.

increasing the current to overcome coil reactance and then stabilizing at the required current to generate a ~ 5 G B field. These levels can be adjusted to run current in both positive and negative bipolar values, such as in configurations for quadrupole coils (e.g., anti-Helmholtz or Helmholtz). This waveform allows precise control of the B-field with a switching time of about $\sim 500 \mu s$.

5.2 Detection of quality of optical pumping

Once the atoms are optically pumped into the extreme magnetic sublevels, we need to check the quality of optical pumping or to see how much of the total population still remains in the wrong m_F states of the hyperfine ground states. The method we want to implement to

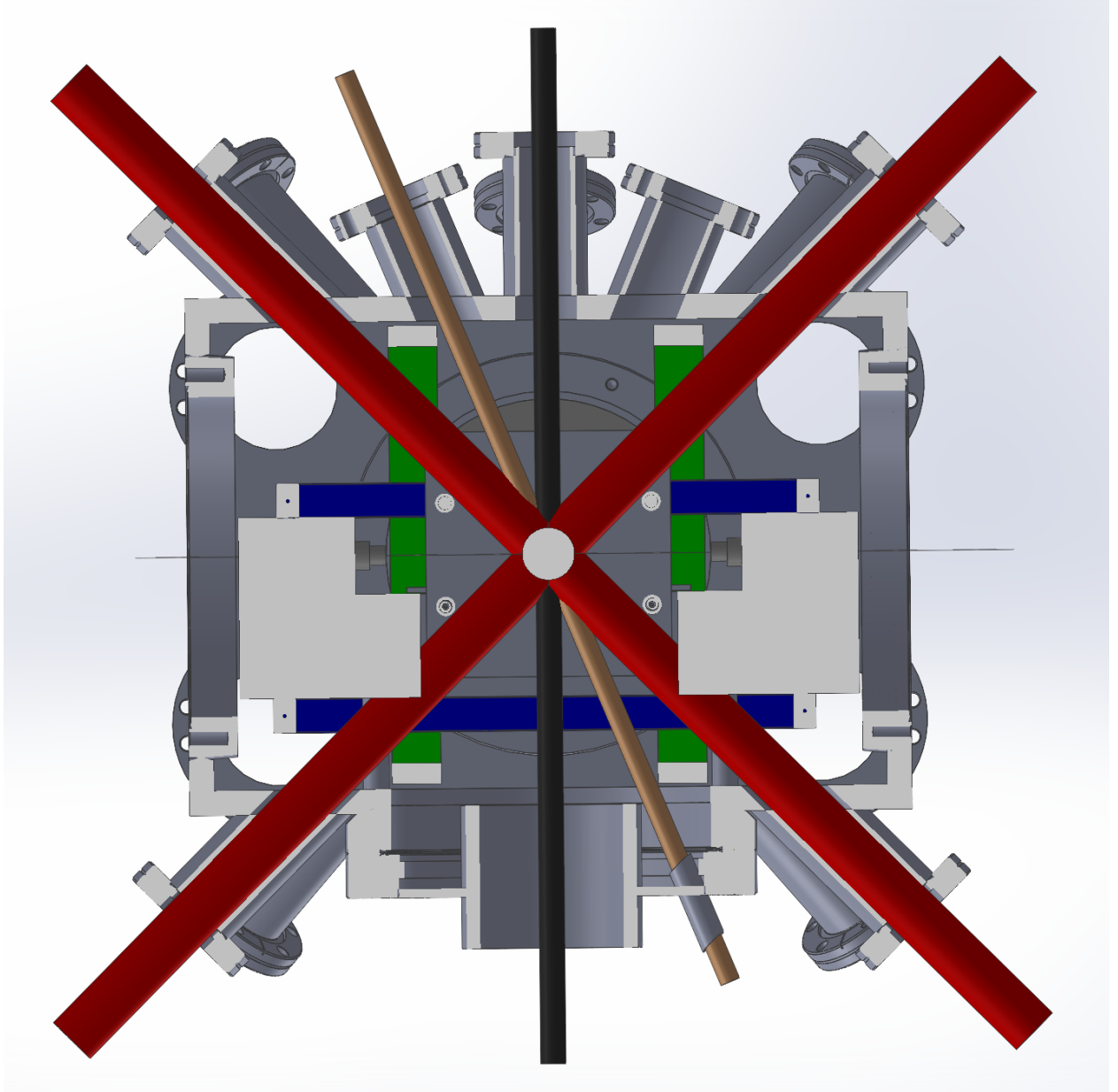


Figure 5.5: A sectional view of the science chamber with coils (in case they are mounted inside), showing a proposed solution to shine the OP beam at 22.5° , along with the modified recessed bottom flange with an extended nozzle to mount optics on the bottom side of the science chamber for retro-reflection of the OP beam.

measure the quality of OP is that where disruption to the system could be minimized. We want to measure in-situ measurement of population distribution of atoms on m_F sublevels, cover the full trapped atomic sample, and we do not want to open the vacuum chamber. There are different ways to explore to check the quality of optical pumping:

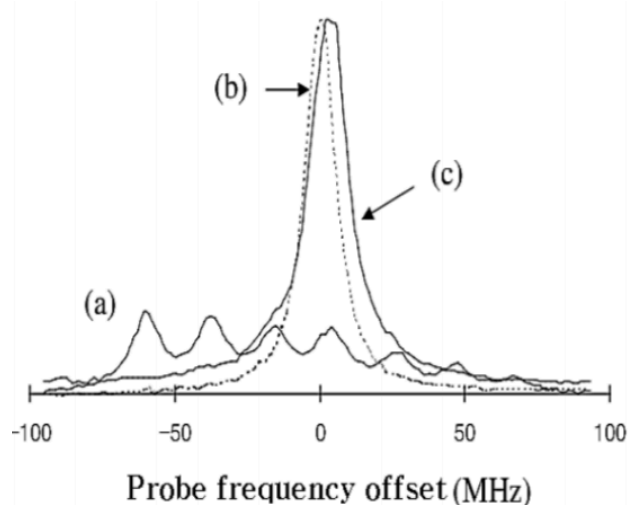


Figure 5.6: Image taken from reference [76] for OP in Cs. A probe frequency scan for MOT shows the distribution of atoms in $|6S_{1/2}, F = 3, m_F\rangle$ in Cs. a) shows the uneven distribution of atoms in sample without optical pumping ($B = 30$ G). b) the sample when the quantization field, $B = 0$, the m_F are non-degenerate, and c) is the case when all atoms are pumped into one m_F and the $B = 30$ G. We aim to see similar peak for with and without optical pumping all our trapped atoms in MOT for Fr and Rb.

1) Increase in B-field – One simple way to isolate Zeeman sublevels is by increasing the B-field so that the m_F levels are well separated. For example, applying a 20 G field shifts adjacent sublevels by 28 MHz. This separation ensures that only atoms from one Zeeman sublevel undergo the transition at a time. By scanning the probe laser across all m_F sublevels, we would get spectra like Fig.5.6(a) without optical pumping and Fig.5.6(c) with optical pumping. The main limitation on increasing the B-field is the ability to switch it off quickly. Higher B-fields generate stronger eddy currents, which can exceed the voltage limits of our supplies and make the fast switching more difficult.

2) Raman spectroscopy - It is a two-photon process where one photon is absorbed and the second one is emitted in a stimulated way in lambda configuration on transition between magnetic sublevels hyperfine ground states via an excited hyperfine state, $7P_{3/2}$. It offers a precise method to evaluate the quality of optical pumping, as the transition linewidth between m_F sublevels can be very narrow (on the order of kHz). The condition to be fulfilled

to completely transfer the population with Raman transition is

$$\Omega_{\text{eff}} \times \tau = \pi, \quad (5.6)$$

where Ω_{eff} is the effective rabi frequency, τ is the Raman pulse duration. The Raman laser intensity, interaction time τ , and detuning from the $7P_{3/2}$ state are chosen to satisfy the π -pulse condition: $\Omega_{\text{eff}} \times \tau = \pi$. The strength of the Raman transition depends on the Clebsch–Gordan coefficients. Based on our estimates, we need a Raman beam intensity of about 30 mW/cm^2 , a detuning of around 800 MHz, and an interaction time of about $130 \mu\text{s}$. The two Raman frequencies will connect the $7S_{1/2}, F = 4$ and $F = 5$ hyperfine ground states. To maintain the phase coherence, both frequencies should come from the same laser source. A laser beam can be placed halfway between the hyperfine ground states; however, buying another laser to drive Raman transitions is both expensive and cumbersome, as it requires maintaining the laser frequency lock, additional frequency stabilization, more optics, and a new setup. Instead of using another laser placed between the two ground states, we propose using an electro-optic modulator (EOM) to generate sidebands.

Using one of the existing lasers (either trap or repump laser) with a 45 GHz sideband generated by an EOM could work, but high-power commercial EOMs with such a bandwidth are difficult to find. This is the main challenge in producing the Raman beams. However, EOspace offers EOM with center frequency of 25 GHz that can generate ± 2 order sidebands, which can be used to provide the phase-coherent light needed for Raman transitions.

5.2.1 Optical pumping launch

To optically pump atoms, a high-quality polarized light beam and retro-reflection are necessary. The transition needs to cycle, and the trapped atoms must experience retro-reflected radiation pressure to avoid excessive momentum transfer from the OP beam. However, the geometry of the science chamber limits the available viewports. Some viewports do not di-

rectly line up with the MOT, and shining the beam through them would result in obstruction by electric field plates or spacers.

One possible solution is to use the 22.5° viewport, as shown in Fig. 5.5. However, there is no available viewport for retro-reflection. One option could be to modify the recessed bottom flange of the chamber with an extended viewport, but this could partially block or clip the OP beam. Additionally, any diffracted light could add unwanted background noise to the detector since the light collection optics are located near the bottom flange.

If the 22.5° viewport is used for the OP beam path, the spin polarization axis of the atoms must be aligned with it during the measurement phase. This means that after OP, the spin-polarized atoms need to be adiabatically rotated by 22.5° to keep the electric field, magnetic field, and $\vec{k}$ vector mutually perpendicular. Our assessment is that this is a challenging task, especially switching the B-field considering the time constraints of the experiment.

To avoid the adiabatic rotation of the magnetic field for spin-polarized atoms, if the 22.5° viewport is used for launching the OP beam, another option considered was to tilt the entire science chamber by 22.5° . This would ensure the magnetic field stays mutually perpendicular to the electric field and $\vec{k}$. However, tilting the chamber introduces significant complications. It would make it challenging to position all the optics for the power buildup cavity, both inside and outside the chamber, at any angle. This would also make routine operation of the chamber more difficult. As a result, this idea was ultimately discarded.

An alternative idea is to share the vertical push beam path with the OP beam. However, the challenge is to retro-reflect the OP beam while blocking the push beam. The detection system in our setup is at the bottom, leaving no room for retro-reflection. We consider placing a $\varnothing 0.5''$ mirror and a small commercial shutter above the first light collection lens in the recessed viewport, but there is not enough vertical space for both. Moving the detection system down slightly could have helped, but access to the mirror for alignment or a mini-motorized stage is still difficult. Moreover, there are limited compact, lightweight, and small-volume motorized stage options. Thus, we abandon this idea.

The most feasible solution with minimal changes to the existing setup is to launch the OP beam using the same path as the push beam, but with modified detection optics to allow retro-reflection of the OP beam. The detection optics consist of both a mirror at 45° and a lens with a hole, allowing the OP beam to be retro-reflected, as shown in Fig. 5.7. Figure 5.8

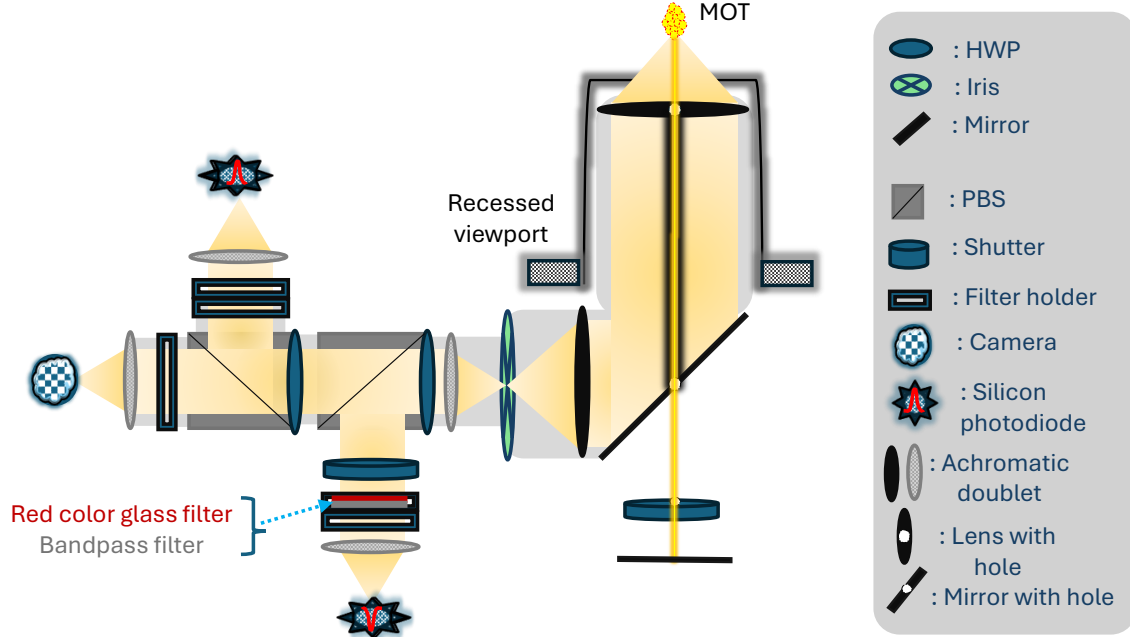


Figure 5.7: Schematic of new detection system with a mirror and a lens with a hole at the bottom to accommodate the retro-reflection of the OP beam.

shows the plan to launch the OP beam alongside the push beam. The OP tower has two launchers on top, enabling the launch of the OP and push beams along the same vertical path, though only one beam is active at a time. The setup must allow for launching two separate beams, where changes to one beam do not affect the other. This requires a solution that accommodates different intensities, polarizations, frequencies, and spatial alignments for each beam. The OP beam will be circularly polarized, while the push beam is currently linearly polarized. Both beams will be combined using a non-polarizing beam splitter, which will preserve the polarization state of the incident light. The frequency of the OP and push beams will be shifted, and they will be turned on and off by driving each beam with a separate AOM.

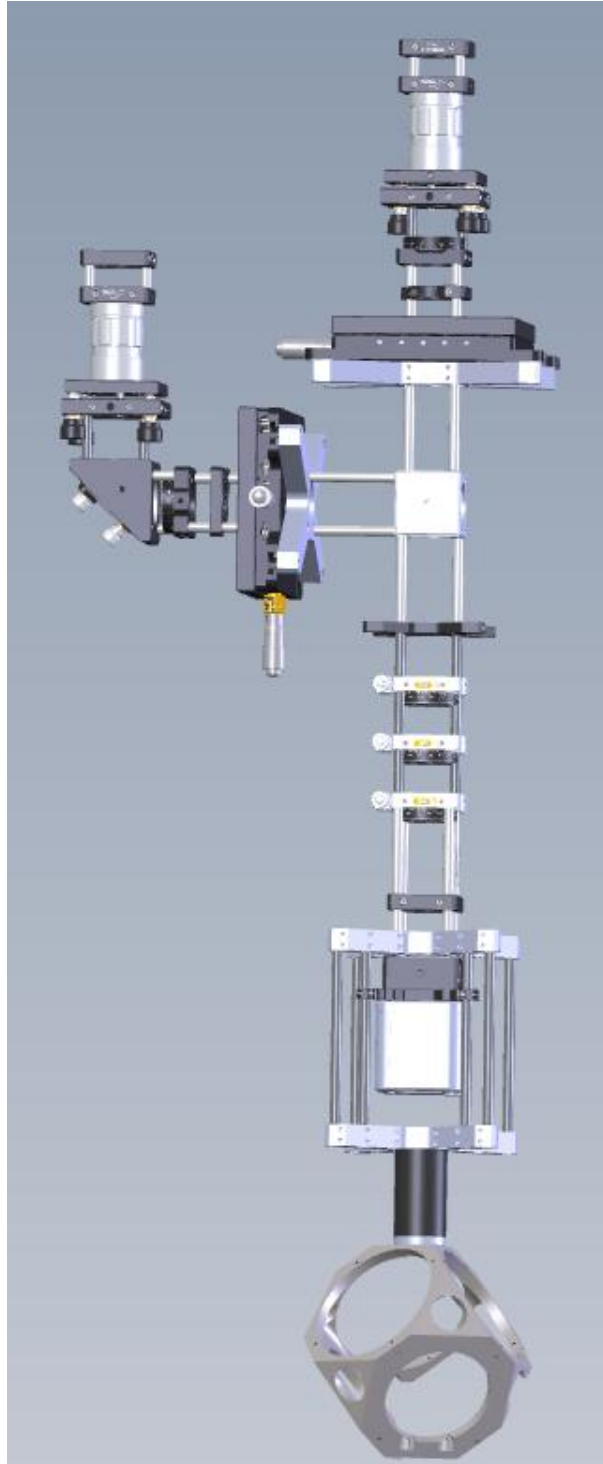


Figure 5.8: shows a CAD model to show the assembly to launch optical pumping and push beam. The assembled optical pumping tower will sit on top of capture trap.

5.2.2 Details of OP tower optics

The main goal was to ensure that the launch of both beams and all optical components are centered with respect to both the upper chamber and the science chamber. The OP tower will first be tested on a test bench to analyze and understand the polarization of both beams. Once the testing is complete, it will be replaced with the push beam launcher, minimizing disruption to the existing setup. The other important part of the OP tower is to correct, control and analyze the polarization of OP beam.

Control over the beams: Both beams will be launched using fiber couplers and their beam sizes can be adjusted with a 2X beam expander (GBE02-B), mounted on a kinematic mount compatible with the 30 mm-cage system. This setup allows for precise angular alignment of the beams. The push beam will be reflected by a 90° angle using a right-angle kinematic mirror mount. A 1" cage rotation mount will be installed on both beams to hold a waveplate, allowing polarization adjustment if needed, followed by an iris. All mounts are designed to hold 1" optics.

An important component is the translation stage, which provides two-axis linear translation for the beam, with a 1" hole to allow clear passage of both OP and push beams. This stage is lightweight to ensure that the OP tower remains manageable. I also designed an adapter plate to make a compatible connection between the translation stage and the cage cubes, accommodating the mounting holes of both optical pieces with different hole configurations.

Control over the polarization: To control the polarization of the optical beam, a combination of half and quarter waveplates can be used to achieve any desired polarization state. The polarized light passes through the optical surface of the glass cell. However, we cannot directly measure the birefringence caused by mechanical stress in the glass without disassembling the apparatus. Instead, we will use a mock glass cell to analyze the polarization quality and determine how to compensate for the effect with a stack of waveplates. This approach will help us understand how to analyze the polarization quality and determine how

to optimize it in order to restore the polarization to the desired state [77]. The stress-induced birefringence will then be corrected in situ. In the future, during long-term operations, manually controlled waveplate mounts will be replaced with motorized optical mounts. This will be especially useful during beamtime when the apparatus is radioactive and cannot be manually adjusted.

Analysis of polarization: The polarization quality will be analyzed using the Thorlabs PAX1000IR1 polarimeter. We also need to measure the polarization quality on the bottom side of the science chamber before the beam is retro-reflected by the mirror. Once the OP tower is mounted on upper chamber, there should be enough space such that the polarimeter can be easily swapped in and out. It will be securely mounted with a CDA1 clamp, allowing for easy swapping without disturbing the rest of the system.

Machine shopped adaptors: Some of the adapters have been machine fabricated for easy mounting. One adapter features a circular O-ring with a slightly extended part, ensuring a snug fit (not too tight or too loose) into the exo-skeleton cage mount on the glass cell. Another adapter creates a suitable connection between the cube mount and the translation stage. Additionally, an adapter has been made to connect the 30 mm cage optics with the 90 mm space, allowing easy removal of the polarimeter. Another adapter will be used to mount the cube mount on the opposite side of the railing. The connection between the OP tower and the exo-skeleton on the glass cell will help with alignment, while other connection from cube mount will ensure that the OP tower is securely mounted and rigid. The adaptor CAD designs are in appendix A.

OP beam: The optical beam will be derived from the trap laser, with an acousto-optic modulator (AOM) applying a frequency shift of around 90 MHz to bring the OP beam into resonance with the cycling transition. The beam intensity will be near the saturation intensity, with a beam diameter of approximately 5 mm.

5.3 Summary

The current MOT setup uses a ($162\ \mu\text{H}$) and a ($\sim 5\ \text{mH}$) bias field that remain on continuously. To spin-polarize the atoms as part of the progression toward APV, the quantization axis must be defined, which requires switching both the quadrupole and bias fields on and off. Progress has been made on switching the bias field by driving the coils with arbitrary waveforms that include controlled overshoot and undershoot, along with exponential edge shaping, to compensate for eddy-current-induced lag. The existing push-beam optical path will be shared with the future optical-pump (OP) beam, allowing the OP beam to be retro-reflected. Apertured optics and a mechanical shutter will block the push beam when needed to prevent its retro-reflection, while a bottom-mounted mirror will retro-reflect the OP beam. Initial OP-quality detection will use a large bias field so individual Zeeman peaks can be resolved. Later, a more efficient detection method will be adopted depending on how frequently OP quality needs to be verified. Implementing OP is the next stepping stone toward APV measurement.

CHAPTER 6

CONCLUSION

Since 2012, the FrPNC lab has made significant progress in the development of apparatus, experimental techniques, all aimed at the ultimate goal of the FrPNC project, which is to observe atomic parity violation in francium. The work so far has included the reliable operation of the neutralizer, coating the inside of the glass cell, robust trapping and transfer of neutral francium atoms from the capture chamber to the science chamber, installation of transparent electric field plates, water-cooled Helmholtz coils to provide the magnetic field, frequency stabilization of the 7S–8S spectroscopy laser, installation of the power buildup cavity, and development of the burst fluorescence technique.

These achievements enabled us to study the highly forbidden 7S–8S transition in francium in a magneto-optically trapped (MOT) sample containing millions of francium atoms within a volume of 1 mm^3 , at micro-kelvin temperatures for tens of seconds. Trapping the francium atoms in a MOT allows us to make use of the minute amounts of Fr much more efficiently compared to atomic beam methods.

A recent addition of a power build-up cavity in the science chamber enhanced the intensity of the 507 nm light by approximately 2000 times at the interaction region, boosting the sensitivity of the apparatus to increase the signal-to-noise ratio and enabling the observation of transitions with an oscillator strength of 10^{-13} , namely the magnetic dipole (M1) transition. To improve the detection efficiency, an additional burst technique was implemented, in which atoms undergoing the 7S–8S transition are cycled on two cycling transitions on the D2 line of francium to generate a burst of photons. The burst technique effectively compensates for the cumulative losses of the entire detection system arising from the detector’s quantum efficiency, the fraction of light reaching the detector, the limited solid angle, and other constraints – such as the use of neutral density filters to protect the detector – not by altering

these factors, but by generating a sufficiently strong signal to overcome them. As a result, the detection efficiency was improved to nearly 0.5, meaning that approximately one photon is detected for every two atoms that undergo the 7S–8S transition. Overall, with the power build-up cavity and the burst technique, we achieved an improvement in detection rate of approximately six orders of magnitude. The pure M1 transition amplitude at 0 V electric field was approximately 20 kHz for this thesis work.

The work presented in this thesis focuses on determining β and $M1_{rel}$ by measuring the hyperfine dependence of $M1/\beta$ between different hyperfine transitions. Our measurement technique for determining $M1/\beta$ is robust and allows us to fully exploit the cycling transitions on the D2 line using the burst technique.

The next major step in the near future will be to optically pump the atoms into specific magnetic sublevels. This will provide another way to extract β by measuring α/β and will serve as a strong test of theory. It will also enable the measurement of the E1-M1 interference, which will play a significant role in searching for asymmetries in the fluorescence signal of the total transition rate under reversals of the electric field, magnetic field, and magnetic sublevels m_F to $-m_F$.

There is a shutdown at TRIUMF until mid-2027, which will not allow any work with francium during this period. However, this time will be used to work with rubidium to establish the optical pumping technique, attempt to observe the M1 transition in rubidium, and upgrade the data acquisition system. These upgrades include improved measurements of atom number and the 507 nm light, enhanced stability of the probe beams, and ensuring that no additional stray light leaks into the detection window during data taking. The FrPNC project has a promising future, with the goal of observing the first interference-type fluorescence signal in francium in 2028.

Appendices

APPENDIX A

ANGULAR MOMENTUM COUPLING COEFFICIENTS ($C_{Fm}^{F'm'}$) FOR FRANCIUM,

$$I = \frac{9}{2}$$

$$C_{5m'}^{5m'} = \frac{m'}{5}$$

$$C_{5m'-1}^{5m'} = \frac{-\sqrt{(6-m')(5+m')}}{10}$$

$$C_{5m'+1}^{5m'} = \frac{+\sqrt{(6+m')(5-m')}}{10}$$

$$C_{5m'}^{4m'} = \frac{+\sqrt{25-m'^2}}{10}$$

$$C_{5m'-1}^{4m'} = \frac{-\sqrt{(5-m')(6-m')}}{10}$$

$$C_{5m'+1}^{4m'} = \frac{-\sqrt{(6+m')(5+m')}}{10}$$

$$C_{4m'}^{5m'} = \frac{+\sqrt{25-m'^2}}{5}$$

$$C_{4m'-1}^{5m'} = \frac{+\sqrt{(4+m')(5+m')}}{10}$$

$$C_{4m'+1}^{5m'} = \frac{+\sqrt{(4-m')(5-m')}}{10}$$

$$C_{4m'}^{4m'} = \frac{-m'}{5}$$

$$C_{4m'-1}^{4m'} = \frac{+\sqrt{(4+m')(5-m')}}{10}$$

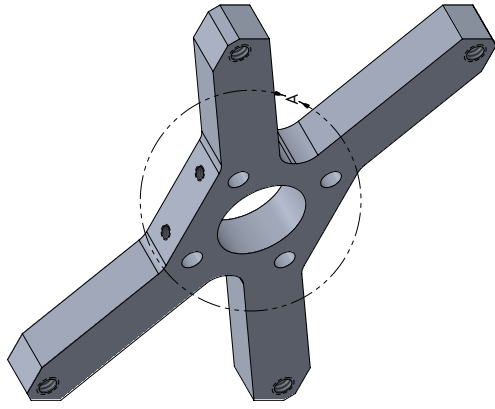
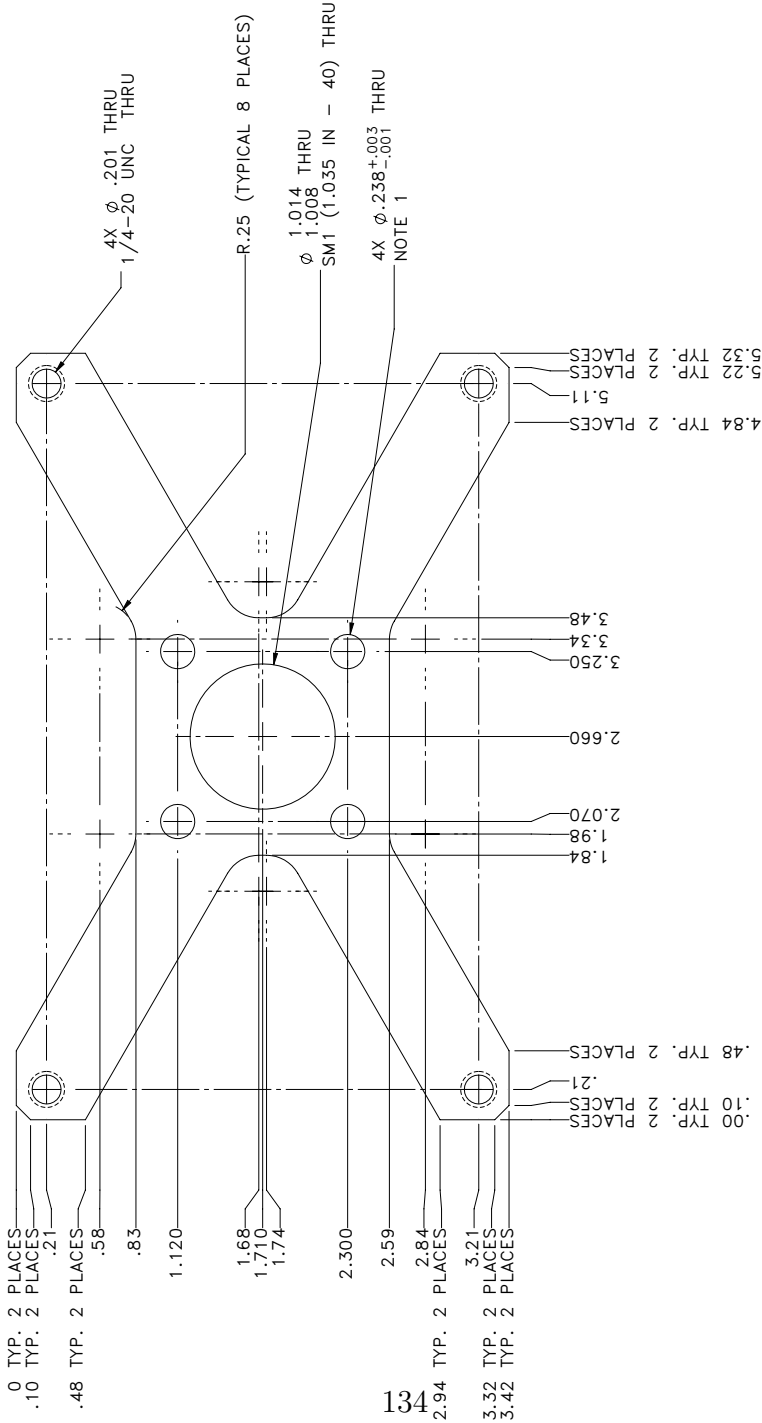
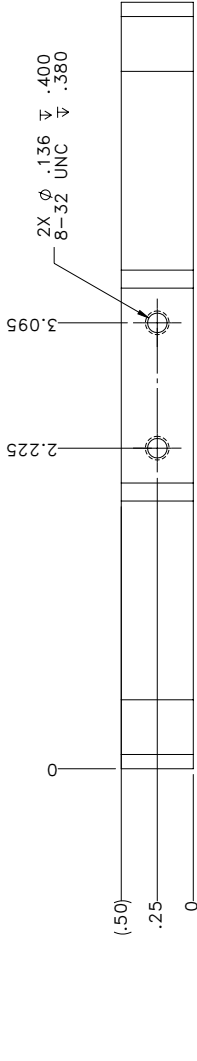
$$C_{4m'+1}^{4m'} = \frac{-\sqrt{(5+m')(4-m')}}{10}$$

These are the angular momentum coupling coefficients for Fr. The transition between the extreme m_F -sublevels have the largest C.G. factor.

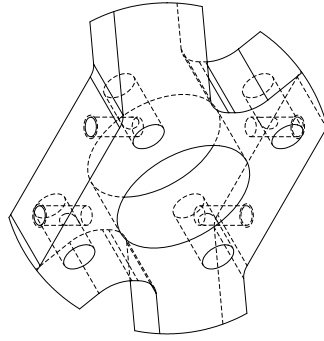
APPENDIX B

ADAPTOR FOR OP TOWER

NOTES:
1. SLIDING FIT FOR 6.00MM ROD



ISOMETRIC VIEW
NOT TO SCALE



DETAIL A

ITEM	DESCRIPTION	MATERIAL	QTY.
1	5.320[135.1] X 3.420[86.9] X 0.50[12.7]	Aluminum	1

DIMENSIONS QUOTED ARE FINISHED DIMENSIONS. NO ALLOWANCE HAS BEEN MADE FOR MANUFACTURE. INTERPRET DIMENSIONING AND TOLERANCING AS PER ASME Y14.5-2009. REMOVE ALL BURRS AND SHARP EDGES.

Canada's particle accelerator centre
Centre canadien d'accélération des particules
4004 Westbrook Mall, Vancouver, BC, V6T 2A3, Canada



ADAPTER FOR XY TRANSITION STAGE OP
ISAC 1
FRANCIUM

SCALE	3:2	DWG NO.	ILE3118	SIZE	SHEET	REV
DATE	2024-01-05			C	OF	1

PERMIT TO PRACTICE	REV	DATE	ECO-7011	CHANGE CONTROL APPLIES	Shaun Georges	AT
TOLERANCES UNLESS OTHERWISE SPECIFIED				REVISION DESCRIPTION	BY	APPD

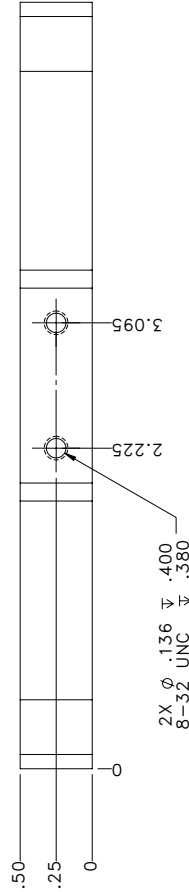
DECIMALS X ±.1
XX ±.01
XXX ±.005
ANGULAR .XXX ±1°
SURFACE FINISH 63 μ inch

ALL DIMS IN INCHES

DESIGNED Anina Sharma	SUB-ASSY
DRAWN Anina Sharma	ASSEMBLY
CHECKED Shaun Georges	
WPN #	
TRACKING #	0-321

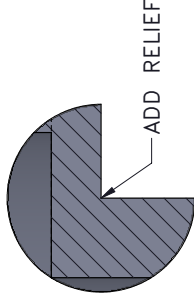
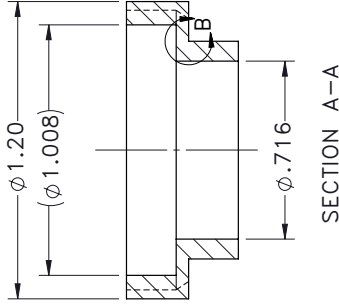
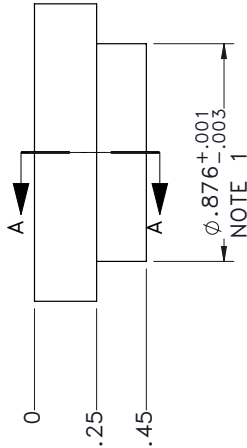
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NOTES:

1. SLIDE FIT REQUIRED IN MATING PARTS $\phi 0.880$ " POCKET.
2. MACHINE PRE-DRILL ϕ .



DETAIL B
SCALE 8 : 1

SECTION A-A

ISOMETRIC VIEW



Canada's particle accelerator centre
Centre canadien d'accélération des particules
4004 Westbrook Mall, Vancouver, BC, V6T 2A3, Canada

OPTICAL PUMPING CIRCULAR ADAPTER
ISAC 1
FRANCIUM

SCALE	2:1	DWG NO.	ILE3120	SIZE	B	SHEET	1	REV	A
DATE	2024-01-05					OF	1		

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DESIGNED	Shaun Georges	SUB-ASSY
DRAWN	Anima Sharma	ASSEMBLY
CHECKED	Shaun Georges	
WPN #		
TRACKING #	0-321	

PERMIT TO PRACTICE	REV	DATE	ECO-7011-	CHANGE	CONTROL	APPLIES	Shaun Georges	BY	APPD
TOLERANCES UNLESS OTHERWISE SPECIFIED									

DECIMALS	.X	± .1
	.XX	± .01
	.XXX	± .005
ANGULAR		± 1°
SURFACE FINISH		63 μ inch

ALL DIMS IN INCHES

APPENDIX C

E FIELD VS TRANSITION RATE

These show the quadratic dependence of the transition rate as a function of the electric field when for $\Delta F = \pm 1$ with Lorentzian profile for 2024 beamtime.

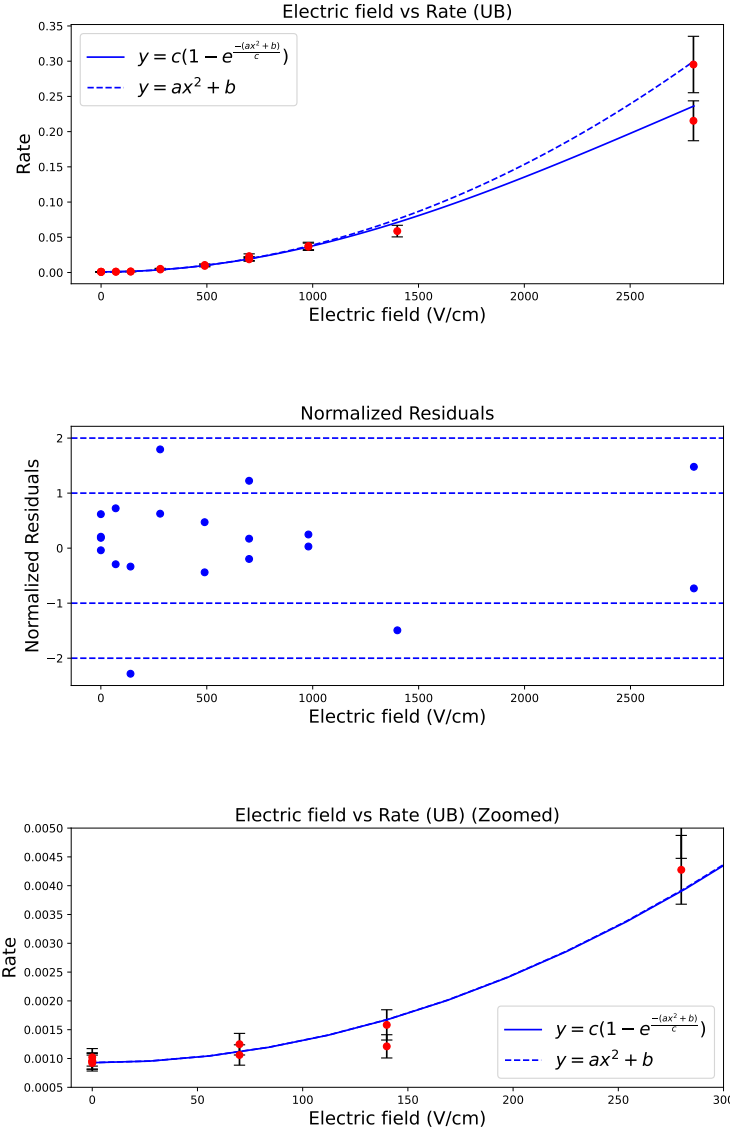


Figure C.1: Quadratic dependence of transition rate as a function of electric field on $\Delta F = +1$, $7S_{1/2}(F = 4) \rightarrow 8S_{1/2}(F' = 5)$ with Lorentzian model. The rate on y axis is in (KHz/atom). The values of fitted parameters from solid line are $a = (3.8 \pm 0.2) \times 10^{-8}$, $b = (9.3 \pm 0.6) \times 10^{-4}$, $c = (6.0 \pm 3.1) \times 10^{-1}$. $M1/\beta$ is 155.9 ± 6.7 V/cm.

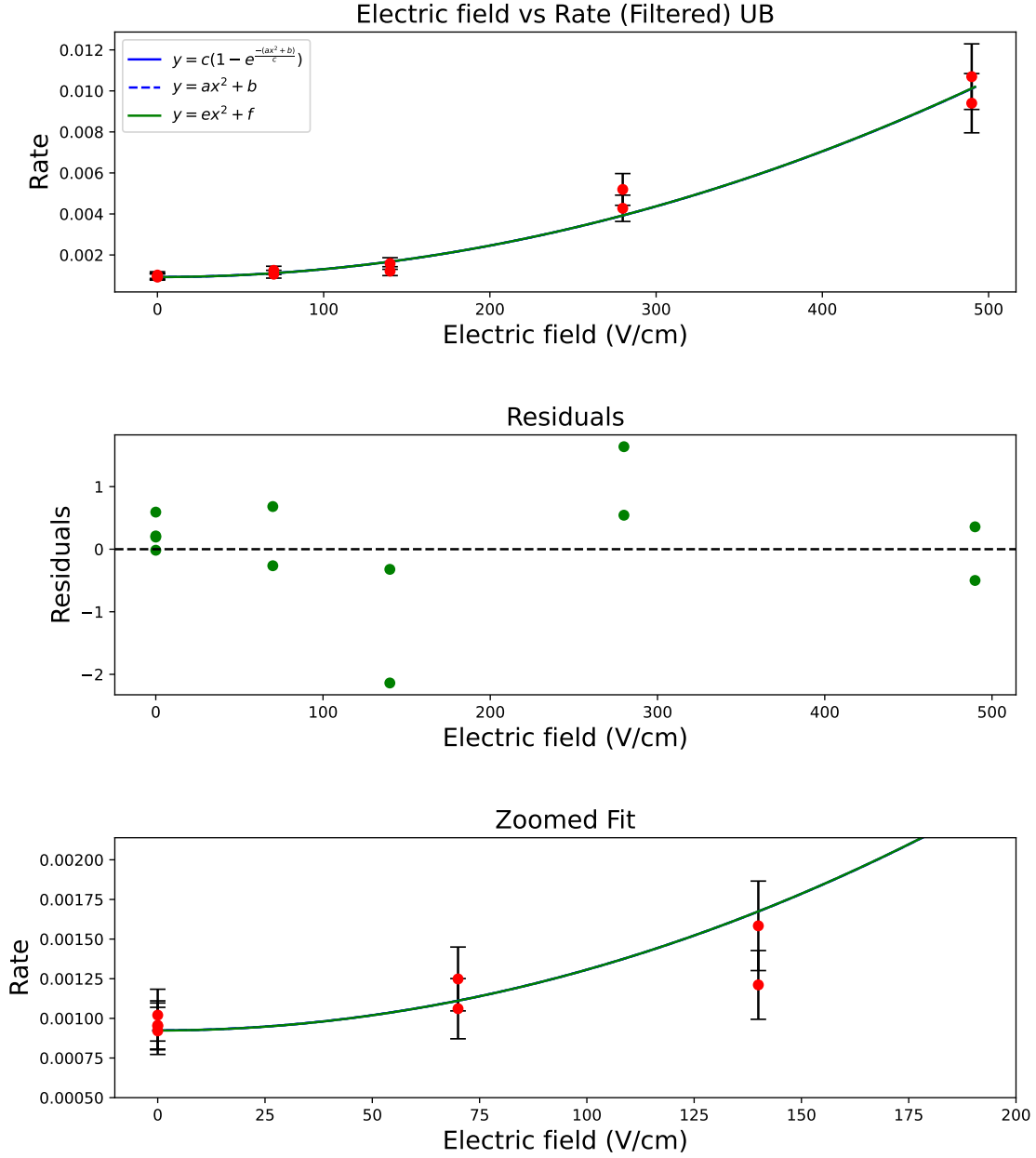


Figure C.2: Quadratic dependence of transition rate as a function of electric field on $\Delta F = +1$, $7S_{1/2}(F = 4) \rightarrow 8S_{1/2}(F' = 5)$ with Lorentzian model. The rate on y axis is in (KHz/atom). The values of fitted parameters from solid line are $a = (3.8 \pm 0.4) \times 10^{-8}$, $b = (9.2 \pm 0.6) \times 10^{-4}$, $c = 1.6 \times 10^{+2} \pm 6.6 \times 10^{+1}$, $e = (3.8 \pm 0.4) \times 10^{-8}$, $f = (9.2 \pm 0.6) \times 10^{-4}$. $M1/\beta$ is 155.2 ± 9.0 V/cm.

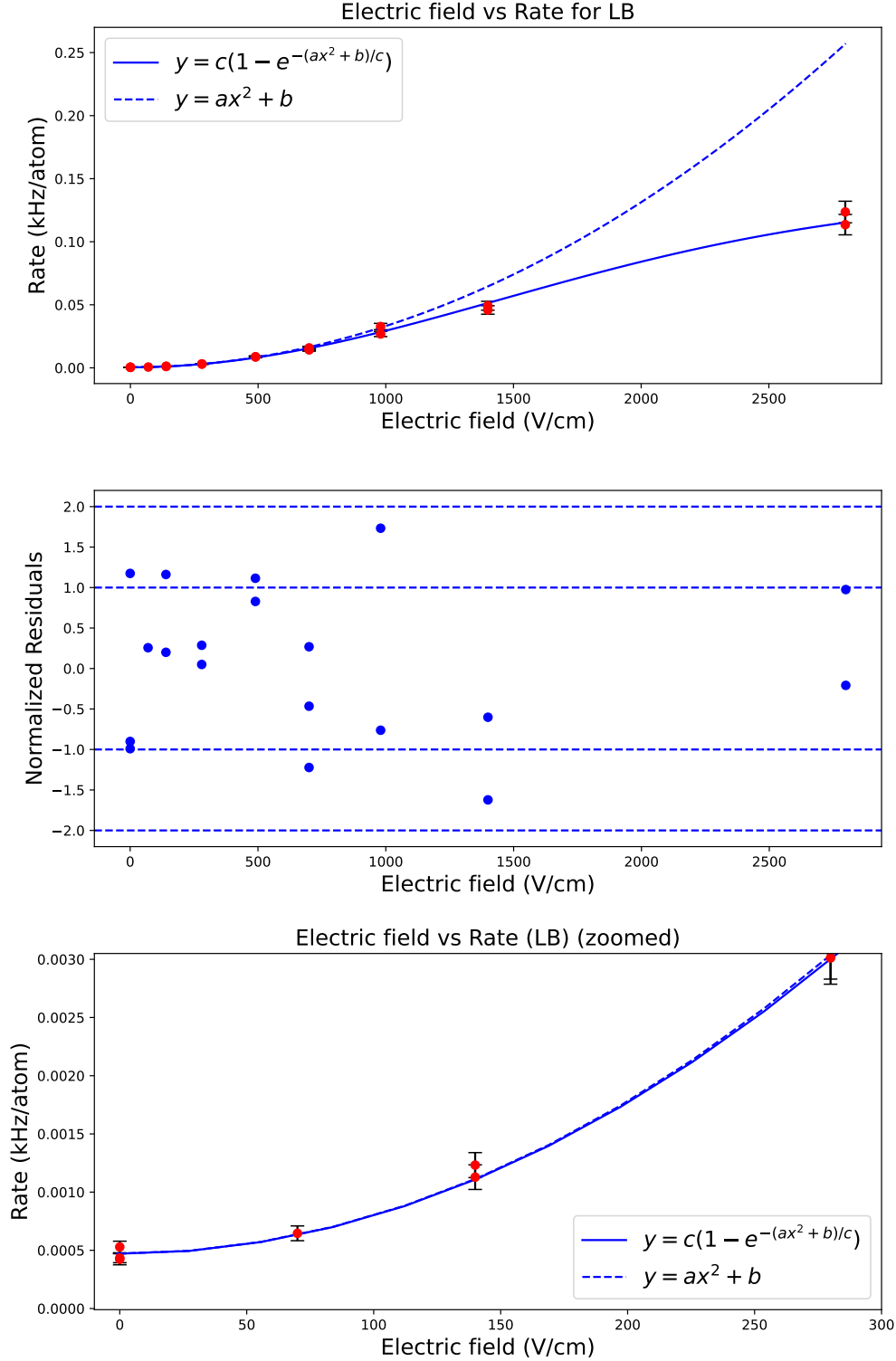


Figure C.3: Quadratic dependence of transition rate as a function of electric field on $\Delta F = +1$, $7S_{1/2}(F = 4) \rightarrow 8S_{1/2}(F' = 5)$ with Lorentzian model. The rate on y-axis is in (kHz/atom). The values of the fitted parameters from the solid line are $a = (3.3 \pm 0.0) \times 10^{-8}$, $b = (4.7 \pm 0.2) \times 10^{-4}$, $c = 1.4 \times 10^{-1} \pm 1.1 \times 10^{-2}$. $M1/\beta$ is 120.0 ± 3.4 V/cm.

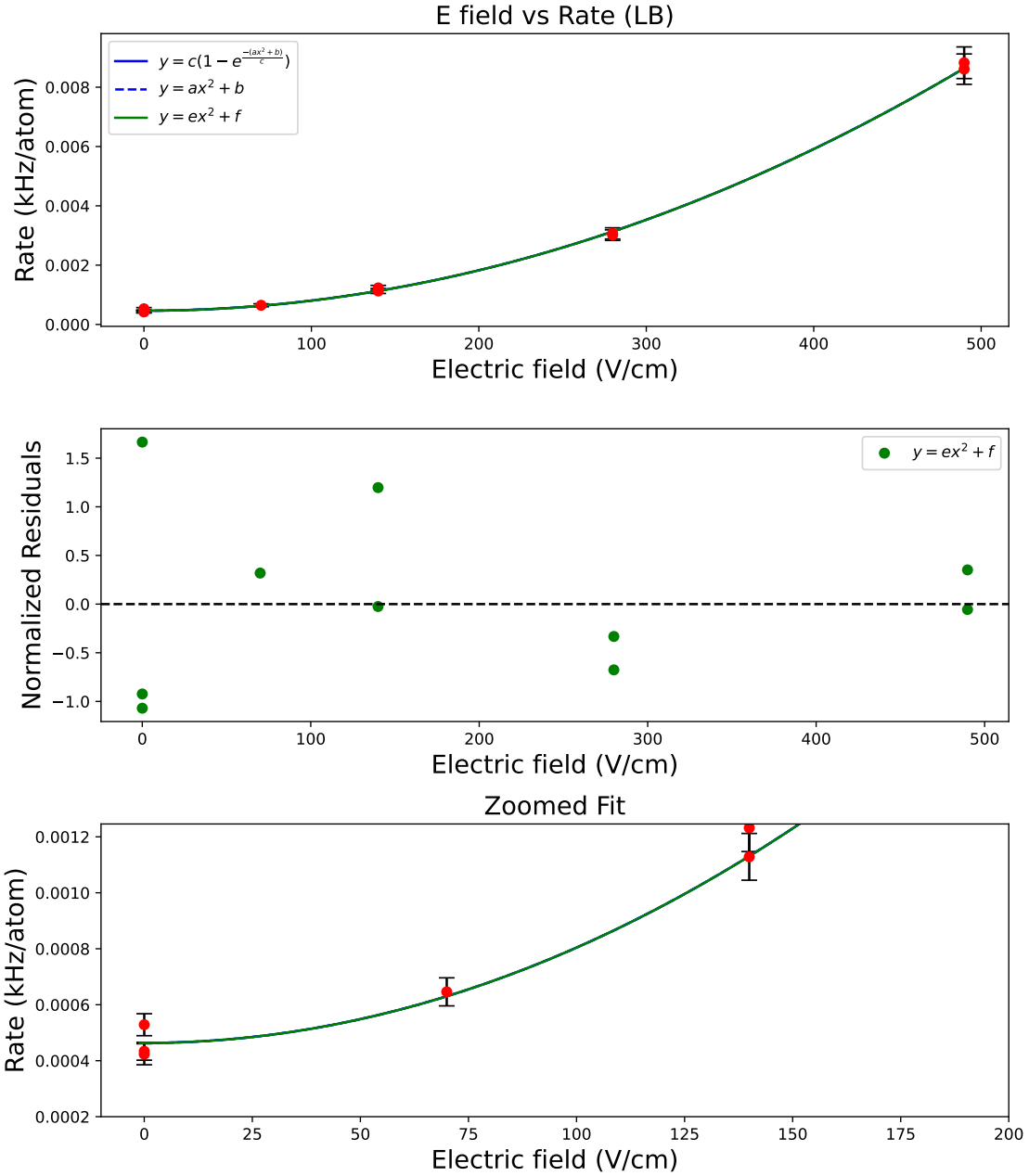


Figure C.4: Quadratic dependence of transition rate as a function of electric field on $\Delta F = +1$, $7S_{1/2}(F = 4) \rightarrow 8S_{1/2}(F' = 5)$ with Lorentzian model. The rate on y-axis is in (kHz/atom). The values of fitted parameters from solid line are $a = (3.4 \pm 0.1) \times 10^{-8}$, $b = (4.6 \pm 0.2) \times 10^{-4}$, $c = 1.5 \times 10^{+1} \pm 9.9 \times 10^{+2}$, $e = (3.4 \pm 0.1) \times 10^{-8}$, $f = (4.6 \pm 0.2) \times 10^{-4}$. $M1/\beta$ is 116.6 ± 3.0 V/cm.

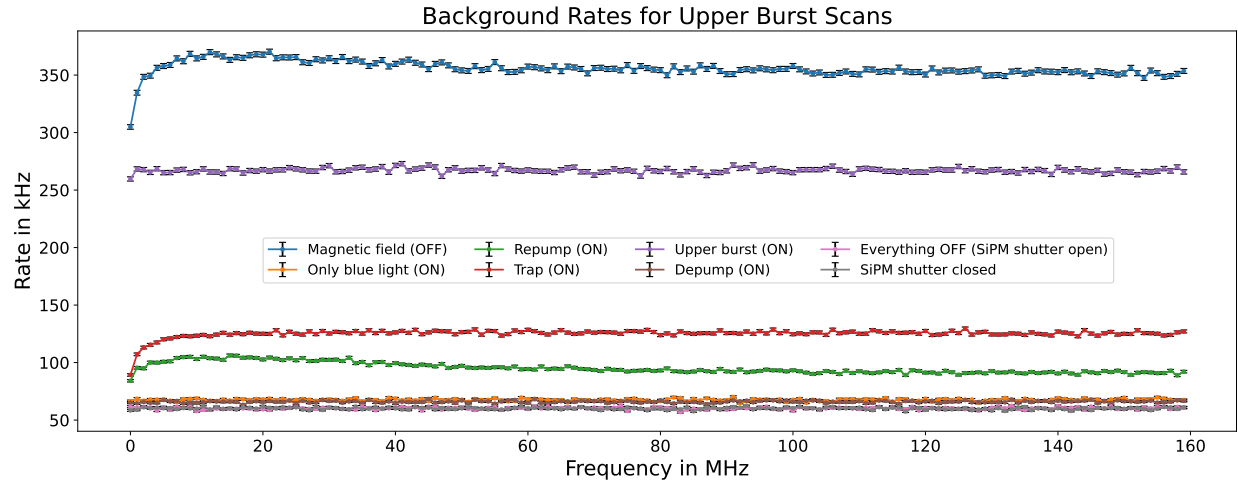


Figure C.5: Backgrounds for UB.

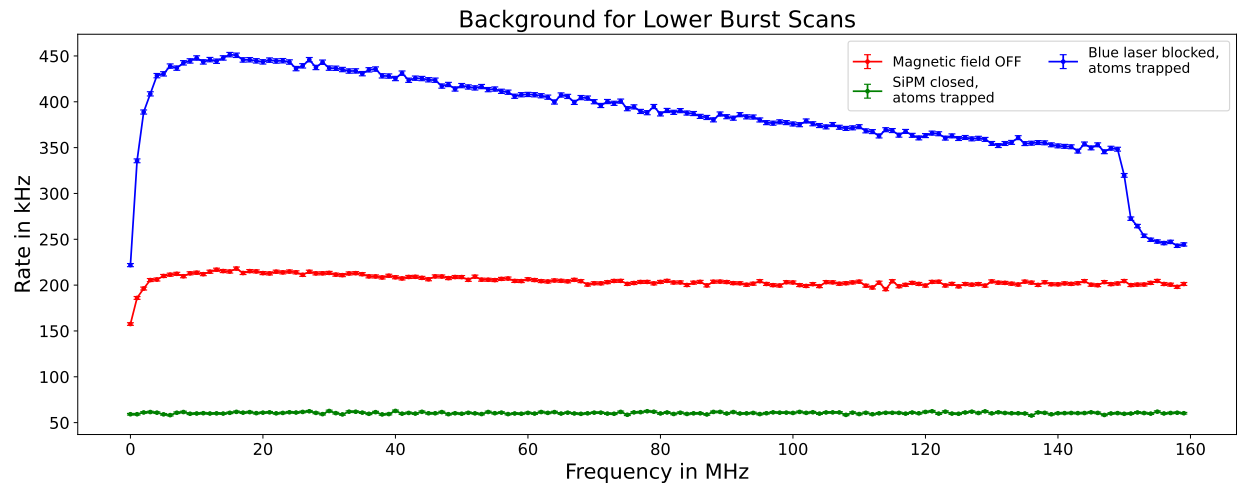


Figure C.6: Backgrounds for LB.

Figure C.7: Information on scan number, electric field, peak area and other numbers on upper burst.

	<i>Electric field (+/- V)</i>	<i>Scan No.</i>	<i>Atom No.</i>	<i>Atom No. Error</i>	<i>Lifetime(s)</i>	<i>Lifetime Error(s)</i>	<i>FWHM Lorentzian (MHz)</i>	<i>FWHM err Lorentzian</i>	<i>FWHM Vog (MHz)</i>	<i>Peak area (Voigt)</i>	<i>Peak area err (Voigt)</i>	<i>Peak area (Lorentzian)</i>	<i>Peak area err (Lorentzian)</i>	<i>Red Ch2 (Lorentzian)</i>	<i>Red Ch2 (Voigt)</i>	<i>Rate (Vog)</i>	<i>Rate err (Vog)</i>	<i>Rate (Lor)</i>	<i>Rate err (Lor)</i>
0	0.0	34.0	388202.0	2727.0	15.6	0.27	5.2	0.23	0.23	341.09	22.31	371.16	14.46	1.89	1.8	0.0009	0.0001	0.001	0.0001
1	100.0	35.0	404251.0	3139.0	14.79	0.27	5.05	0.23	0.23	465.34	29.71	504.68	19.55	2.52	2.4	0.0012	0.0001	0.0012	0.0001
2	4000.0	36.0	439292.0	9752.0	10.76	0.43	7.63	0.17	0.17	114577.95	1487.36	129713.14	2314.94	138.18	23.48	0.2608	0.0147	0.2953	0.017
3	700.0	37.0	396426.0	3732.0	14.18	0.3	5.71	0.22	0.22	2899.52	106.07	3726.72	123.64	35.74	11.62	0.0073	0.0005	0.0094	0.0006
4	2000.0	38.0	327990.0	5023.0	17.07	0.69	5.59	0.19	0.19	16350.42	458.82	19279.84	541.86	122.37	44.0	0.0499	0.003	0.0588	0.0035
5	1000.0	41.0	424205.0	4744.0	14.12	0.35	5.24	0.19	0.19	8393.32	310.31	9863.56	291.89	147.87	81.38	0.0198	0.0013	0.0233	0.0014
6	1400.0	43.0	424544.0	4480.0	13.01	0.29	5.76	0.21	0.21	12964.33	324.86	15998.96	482.5	154.74	34.9	0.0305	0.0017	0.0377	0.0022
7	200.0	44.0	389957.0	3280.0	16.76	0.37	6.17	0.33	0.33	385.92	36.2	472.22	23.28	1.75	1.49	0.001	0.0001	0.0012	0.0001
8	400.0	46.0	386227.0	3184.0	14.86	0.29	5.32	0.18	0.18	1653.0	53.6	2006.47	60.1	15.84	6.66	0.0043	0.0003	0.0052	0.0003
9	0.0	47.0	364130.0	2242.0	14.99	0.22	4.82	0.21	0.21	316.47	19.95	371.51	14.22	2.94	2.31	0.0009	0.0001	0.001	0.0001
10	1000.0	50.0	331647.0	6533.0	17.87	0.97	5.29	0.2	0.2	5315.25	212.09	6216.63	192.14	15.72	9.69	0.016	0.0011	0.0187	0.0012
11	0.0	52.0	389229.0	3289.0	13.57	0.25	4.76	0.19	0.19	347.93	17.35	370.41	12.15	1.17	1.13	0.0009	0.0001	0.001	0.0001
12	0.0	54.0	385945.0	3184.0	14.48	0.27	4.79	0.22	0.22	314.24	20.09	355.44	13.75	2.29	2.0	0.0008	0.0001	0.0009	0.0001
13	4000.0	55.0	379615.0	4645.0	12.11	0.29	6.73	0.18	0.18	71103.39	914.15	81764.48	1757.7	693.6	98.26	0.1873	0.0099	0.2154	0.012
14	200.0	56.0	352259.0	3282.0	16.66	0.4	5.41	0.3	0.3	443.27	35.33	557.79	27.17	2.44	1.79	0.0013	0.0001	0.0016	0.0001
15	1400.0	60.0	348956.0	3992.0	15.16	0.41	5.7	0.21	0.21	10269.86	278.08	12747.5	388.23	116.26	28.84	0.0294	0.0017	0.0365	0.0022
16	100.0	61.0	339589.0	3112.0	15.38	0.34	5.13	0.28	0.28	306.08	25.48	360.3	17.72	2.29	2.0	0.0009	0.0001	0.0011	0.0001
17	700.0	62.0	347740.0	3222.0	15.62	0.35	5.2	0.19	0.19	3196.25	140.74	3719.07	111.94	28.71	19.17	0.0092	0.0006	0.0107	0.0006
18	400.0	63.0	378784.0	3767.0	13.99	0.31	5.47	0.19	0.19	1377.19	60.9	1619.51	48.89	8.08	5.16	0.0036	0.0002	0.0043	0.0003
19	1000.0	64.0	333689.0	3169.0	15.87	0.38	5.59	0.18	0.18	5399.91	144.4	6588.71	180.52	97.09	27.73	0.0162	0.0009	0.0197	0.0011

Figure C.8: Information on scan number, electric field, peak area and other numbers on lower burst.

	<i>Electric field (V_{eff} V)</i>	<i>Scan No.</i>	<i>Atom No.</i>	<i>Atom No. Error</i>	<i>Lifetime(s)</i>	<i>Lifetime Error(s)</i>	<i>FWHM Lorentzian (MHz)</i>	<i>FWHM err Lorentzian</i>	<i>FWHM Vog (MHz)</i>	<i>Peak area (Vogt)</i>	<i>Peak area err (Vogt)</i>	<i>Peak area (Lorentzian)</i>	<i>Peak area err (Lorentzian)</i>	<i>Red Chi/2 (Lorentzian)</i>	<i>Red Chi/2 (Vogt)</i>	<i>Rate (Vog)</i>	<i>Rate err (Vog)</i>	<i>Rate (Lor)</i>	<i>Rate err (Lor)</i>
0	1000.0	3.0	437599.0	4272.0	11.57	0.21	5.18	0.16	5.63	5766.83	213.16	6567.6	171.02	55.75	35.82	0.0132	0.0008	0.015	0.0009
1	0.0	4.0	408735.0	3564.0	10.85	0.17	4.36	0.26	4.57	169.06	0.0	216.07	11.12	1.92	1.46	0.0004	0.0	0.0005	0.0
2	700.0	7.0	429627.0	4370.0	10.81	0.2	4.9	0.17	5.32	3357.72	151.19	3791.45	111.15	59.95	45.3	0.0078	0.0005	0.0088	0.0005
3	2000.0	8.0	402958.0	3899.0	10.64	0.18	5.58	0.14	6.19	17946.05	468.97	19844.18	412.88	98.09	56.84	0.0445	0.0025	0.0492	0.0027
4	200.0	9.0	381607.0	3618.0	11.9	0.22	5.05	0.25	5.17	378.42	0.0	470.12	20.44	2.69	1.84	0.001	0.0001	0.0012	0.0001
5	400.0	10.0	351353.0	3154.0	12.63	0.23	4.9	0.18	5.25	946.91	47.21	1078.45	34.34	5.71	4.41	0.0027	0.0002	0.0031	0.0002
6	4000.0	11.0	531706.0	6717.0	7.57	0.13	6.38	0.1	7.15	60370.93	494.12	65708.52	837.99	168.99	28.14	0.1135	0.0059	0.1236	0.0066
7	100.0	12.0	411921.0	3812.0	11.36	0.2	4.69	0.3	4.85	213.75	20.72	266.11	14.83	1.96	1.63	0.0005	0.0001	0.0006	0.0
8	1400.0	13.0	333400.0	4753.0	10.24	0.25	5.2	0.18	5.71	9743.94	376.7	10921.33	302.12	77.78	55.1	0.0292	0.0019	0.0328	0.0019
9	1000.0	14.0	338647.0	4269.0	12.17	0.31	5.02	0.14	5.52	4395.91	150.37	4824.46	112.2	30.44	22.79	0.013	0.0008	0.0142	0.0008
10	700.0	15.0	404140.0	3948.0	11.41	0.21	5.0	0.16	5.45	3134.12	132.18	3479.19	95.84	27.13	21.42	0.0078	0.0005	0.0086	0.0005
11	0.0	16.0	380724.0	3640.0	12.3	0.24	4.71	0.27	4.94	129.13	11.77	165.15	8.3	1.07	0.82	0.0003	0.0	0.0004	0.0
12	200.0	17.0	394209.0	2688.0	11.87	0.16	4.8	0.27	5.02	343.51	0.0	444.85	22.85	3.02	2.14	0.0009	0.0	0.0011	0.0001
13	400.0	20.0	403222.0	3399.0	12.31	0.21	5.0	0.15	5.37	1066.25	44.66	1214.47	33.1	4.76	3.39	0.0026	0.0002	0.003	0.0002
14	1400.0	21.0	412765.0	4682.0	11.34	0.24	5.34	0.18	5.82	9688.83	388.74	11079.3	316.52	82.45	53.84	0.0235	0.0015	0.0268	0.0016
15	1000.0	22.0	435414.0	3725.0	11.09	0.17	5.14	0.17	5.57	6038.25	216.76	6900.56	187.86	68.8	41.59	0.0139	0.0009	0.0158	0.0009
16	2000.0	23.0	460747.0	5274.0	11.04	0.23	5.8	0.17	6.34	18720.52	539.92	21148.28	496.49	145.81	76.33	0.0406	0.0024	0.0459	0.0026
17	4000.0	24.0	510911.0	6707.0	8.19	0.16	6.96	0.15	7.65	51892.4	704.64	58020.04	1032.47	249.45	60.54	0.1016	0.0054	0.1136	0.0062
18	0.0	25.0	434185.0	3817.0	11.82	0.2	4.45	0.35	5.25	130.19	6.55	183.69	12.83	2.7	2.06	0.0003	2e-05	0.0004	4e-05

APPENDIX D

DIPOLE FORCE EFFECT ON PBC

The intensity gradient of the electric field polarizes the atom creating an optically induced dipole moment. The dipole force [78] on atom in a gradient of E is $F_{\text{dip}} = -\frac{1}{2}\alpha \nabla E^2$ where α is the polarizability of the atom. The electric field of laser beam drives the electric dipole transitions and the magnetic field of blue laser beam drives the magnetic dipole transition. It is known that in the standing wave, the electric and magnetic fields are out of phase by 90° [79] So, when the blue beam enters the cavity, it bounces inside the cavity, the electric and magnetic fields of the beam have opposite sign upon reflection from cavity mirrors. It implies that the node for one mode of field behaves as anti-node for the another, meaning the intensity of electric field of laser is maximum at node (say) but magnetic field intensity will be minimum at the same node. Thus, the high intensity spots for electric and magnetic field of blue beam are separated by $\lambda/2$. The interaction dipole potential [80] of the optically induced dipole moment in E is below:

$$U_{\text{dip}}(r) = -\frac{3\pi c^2}{2\omega_0^3} \left(\frac{\Gamma}{\omega_0 - \omega} + \frac{\Gamma}{\omega_0 + \omega} \right) I(r), \quad (\text{D.1})$$

where ω_0, ω is the atomic resonance, frequency of the standing wave respectively. The dipole potential depends on intensity, detuning $(\omega - \omega_0)$, and driving frequency ω . This intensity gradient of the blue beam interacts dispersively with the optically induced dipoles creating an optical dipole force. Being dispersive in nature, the polarizability changes the sign of the gradient force on the atoms gets attracted toward and outward from the region of high intensity field regions of blue beam. $\omega_0 < \omega$, $U_{\text{dip}}(r)$ becomes positive making the dipole force pulls the atoms into high intensity region of blue beam and vice-versa. In other words for our normal scenario, our blue beam 507 nm is blue-detuned to 718 nm, 817 nm transitions

in Fr. It makes the $\omega_o > \omega$ making $U_{\text{dip}}(r)$ negative, repulsive in nature. The atoms will be pushed out of the high intensity region (nodes), decreasing E1 rate but increasing the M1 rate. It implies that the rate of excitation of atoms is position dependent.

If the atoms are uniformly distributed, the overall excitation rate would be same for both E1 and M1 components. However, for our case, there come a few scenarios where the effect of the dipole force should be considered on trapped atoms in cavity: 1) The standing wave could move the atom as the cavity length is adjusted to maintain the periodic boundary conditions. 2) to avoid photoionization, in data taking period, trapped atoms see alternatively the trapping beams with the standing wave for blue beam. This alternate on/off happens at 1 ms means the trapped atoms do not intend see the same amount of blue light from the point light begin to fill in the cavity to the end of the blue beam on-time window. In addition to this, the standing wave of the blue light beam inside the cavity has a waist inside the cavity which means the intensity distribution of blue beam inside the cavity is not uniform.

In 2021 measurement of $M1/\beta$ measured the excitation rate of pure M1 at zero electric field vs different light levels inside the cavity. If there is a strong dipole force effect, then, with increasing blue light levels, the linear behavior of light levels vs the excitation rate of M1 should show deviations. No deviation was noticed meaning the dipole force effect is negligible.

However, to minimize the position dependence of excitation rate of M1, the blue beam could be phase modulated by creating sidebands at one free spectral range of PBC. The time averaged Bessel function inside the cavity gives the intensity distribution. It is widely known that on frequency modulating the carrier beam (blue beam) going to the cavity, the sidebands will be created at integer multiples of modulating frequency. The modulation index creates the amount of power in the sideband relative to the carrier beam. To have the uniform intensity distribution, we want the power distribution to follow [81]:

$$2J_1^2 = J_0^2 + 2J_2^2, \quad (\text{D.2})$$

where J_0 represents the Bessel function for the carrier beam, and the first, second order sidebands (J_1, J_2) respectively. To generate sidebands separated by one free spectral range (FSR) of the power build-up cavity (PBC), we employ a second electro-optic modulator (EOM), model PM9-VIS-937 from QUBIG (referred to as EOM2 in Fig. 3.1). This modulator is driven by a commercial driver (QDG10-A2, QUBIG), which integrates a high-frequency signal generator and an RF power amplifier, providing stable operation over a frequency range of 0.85–1 GHz. The EOM itself is resonant at 937 MHz, with a tunable range from 668.1 to 987.8 MHz and a bandwidth of 2.48 MHz. It features a $3 \times 3 \text{ mm}^2$ aperture, and the crystal is anti-reflection coated for the 360–650 nm wavelength range to minimize optical losses. The features of EOM work for both 496 nm and 507 nm. To optimize sideband generation, the modulation index is set to 1.2 rad, satisfying the condition outlined in Eq. (D.2). The sidebands are monitored by scanning the PBC, achieved by modulating the voltage applied to a fast piezo actuator mounted on one of the cavity mirrors. This configuration enables precise control and observation of the cavity response during modulation. Once the piezo scanning frequency is high enough for the TEM_{00} mode to appear stable on the oscilloscope, we adjust the RF power of EOM2 to tune the height of the first-order sidebands. Our goal is to generate sidebands at 937.5 MHz, corresponding to one FSR of the PBC. Since these sidebands overlap with the carrier, it is difficult to verify if their amplitude satisfies the condition in Eq. (D.2). To address this, we first generate sidebands slightly red- and blue-detuned from the carrier. Lorentzian fits are applied to these sidebands, and the area under each curve is calculated. A spline is then fitted to these values to estimate the exact sideband height at the carrier frequency.

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