

Laser-Ablated Beryllium Ions for Cold Antihydrogen in ALPHA

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Abstract

One of the best ways to study antimatter is to investigate antihydrogen, the bound state of an antiproton and a positron. Antihydrogen atoms do not exist naturally and must be synthesized in the lab by merging carefully-prepared plasmas of positrons and antiprotons. If the atoms are created in a magnetic trap like the one used by the ALPHA experiment at CERN, then a fraction of the coldest atoms remain trapped, while the rest escape and annihilate on the trap walls. The trapped atoms may then be probed using microwaves or lasers to make high-precision comparisons with hydrogen.

Increasing the trapping rate would allow us to perform precision measurements on antihydrogen in a shorter period of time and with better systematics. Particle simulations indicate that by sympathetically cooling positrons using laser-cooled beryllium ions, we have the ability to improve the antihydrogen trapping rate by up to two orders of magnitude. This thesis describes the effort to design and qualify a beryllium ion source that is compatible with the extreme environment of the ALPHA trapping apparatus. To produce the ions, pulsed laser ablation of a beryllium target is investigated and the ion plume is characterized. By carefully choosing the laser parameters, a plume with a suitable number of ions and kinetic energy distribution can be created for subsequent trapping and laser-cooling in a Penning trap. This thesis reports the successful demonstration of the ion source, in particular its compatibility with the requirements of the ALPHA experiment. The ion source has been installed into the main apparatus at ALPHA and is expected to be commissioned early next year.

Declaration

The author confirms that this work has not previously been accepted in substance for any degree and is not being concurrently submitted in candidature for any degree.

Muhammed Sameed
Swansea, September 24, 2017

A handwritten signature in black ink, consisting of a stylized 'S' followed by a horizontal line.

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To my family.
In loving memory of my grandparents.

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Chapter 1

Introduction

1.1 Matter-Antimatter Asymmetry and Antimatter Gravity

The Standard Model of Particle Physics is one of the most successful scientific theories of modern times. It explains, to an extremely high degree of accuracy, how fundamental particles in nature interact with each other, mediated by the electromagnetic, weak and strong forces. It successfully incorporates the principles of quantum mechanics and special relativity into its framework, and makes predictions that can be tested with our present-day technology.

The Standard Model asserts that at the moment of the creation of our universe, also known as the Big Bang, equal amounts of matter and antimatter were created. Collisions between these particles resulted in the creation of energy in the form of photons, which again converted to matter and antimatter particles. The Universe was thus in a state of constant flux, containing matter-antimatter as well as photons. As the Universe expanded and cooled down, most of the matter and antimatter would have annihilated, leaving us with an ocean of background energy and vast emptiness. When we look at the Universe today however, we see that the cosmos is dominated by matter [1], in the form of stars, galaxies, and interstellar dust, with barely any an-

timatter [2, 3]. This stark contrast from what the Standard Model predicts has confounded scientists for many decades now, demanding serious effort dedicated towards explaining this flagrant discrepancy.

Presently, our best understanding of the laws of physics indicates that there are four known fundamental interactions in the universe: electromagnetism, weak interaction, strong interaction, and gravity. While the Standard Model does an excellent job at accurately describing the first three of these interactions, it does not incorporate gravity into its framework. As such, it does not provide any prediction on how antimatter particles interact gravitationally with matter particles. Various arguments have been presented that insist that the gravitational response of antimatter to matter should be exactly the same as that between matter and matter [4–6]. A handful of scientists believe otherwise however and have formulated modified theories that incorporate gravitational repulsion of antimatter [7–9]. An accurate response of antimatter to gravity is yet to be measured [10].

1.2 Antihydrogen Experiments

The existence of antimatter was first theorized by Paul Dirac in 1931 [11]. Shortly thereafter, Carl Anderson in 1933 published the discovery of the first antimatter particle, the positron, in a cloud chamber experiment he set up to investigate cosmic rays [12]. Both of these individuals went on to win Nobel Prizes for their contributions. It was only after the construction of the Bevatron Accelerator that the antiproton (1955) [13] and antineutron (1956) [14] were also discovered. In the decades to follow, physicists have carefully measured some of the properties of antimatter particles, for example the positron-to-electron mass ratio [15] and the positron g-factor [16].

The prospect of creating and studying neutral antihydrogen atoms motivated the establishment of the Low Energy Antiproton Ring (LEAR) at CERN in Geneva, Switzerland [17]. Approved in 1980 [18], LEAR was built in less

than two years in 1982 and continued operations until 1996 [19]. In 1995, the PS196 experiment at LEAR became the first experiment to successfully create and detect nine atoms of antihydrogen [20], beating the E862 experiment at Fermilab in Chicago, USA, which also successfully created 57 antihydrogen atoms in 1997 [21]. In both of these experiments, the anti-atoms were produced at relativistic speeds, making them unfeasible for trapping and precision experiments.

While experiments like PS196 were focusing on making antihydrogen atoms in flight, a group of physicists at LEAR were working in parallel to capture, cool, and store antiprotons with the aim of creating cold, trappable antihydrogen atoms [22–26]. The success of these pioneering experiments laid down the foundation for the construction of the Antiproton Decelerator (AD), a dedicated ultra-low energy antiproton facility designed for experiments with cold antiprotons and trapped antihydrogen atoms [27]. Approved in 1997, the AD was completed in 2000 and started operating in the same year. ATHENA [28], ATRAP [29] and ASACUSA [30] were the first experiments to receive antiprotons from the AD. These experiments were later joined by ACE [31], ALPHA [32], AEGIS [33], GBAR [34] and BASE [35]. Presently, only five out of these eight experiments are operational: ATHENA dissolved in 2005 and some of its members teamed up with new collaborators to form ALPHA; ACE completed its operations at the end of 2013; and GBAR is not expected to start before 2018.

With an increasing number of experiments at the AD, an additional accelerator known as ELENA (Extra Low ENergy Antiprotons) is being commissioned that will slow down the antiproton beam even further before delivering it to the experiments. As a result of the reduced energy of the antiprotons, the fraction of antiproton capture by the individual experiments is expected to increase by a factor of ~ 100 . ELENA is expected to start delivering antiprotons to GBAR in 2018, with the remainder of the experiments being connected in 2019 and 2020.

1.3 Thesis Outline

This thesis is roughly divided into two parts. The first part (Chapters 2 - 5) outlines the basics of how antihydrogen atoms are created and trapped in ALPHA. The main components of the ALPHA apparatus are described in Chapter 2, in particular the particle sources and particle traps. Chapter 3 explores the various plasma manipulation techniques that have been developed over the years that enable precise control of antiproton and positron plasmas for antihydrogen formation. In Chapter 4, the steps that lead to the creation of antihydrogen atoms are discussed in detail. This chapter also explores the fate of the atoms once they have been formed. Chapter 5 is a key component of this thesis. It discusses how antihydrogen formation and trapping is influenced by antiproton and positron plasma temperatures, and motivates the remainder of this thesis.

With the knowledge that plasma temperatures affect antihydrogen formation and trapping, the second part of this thesis (Chapters 6 - 11) explores the potential of sympathetically cooling positrons in ALPHA in order to increase the antihydrogen trapping rate. Chapter 6 contains a review of previous work carried out to sympathetically cool positrons, and discusses simulations that highlight its feasibility in ALPHA with beryllium ions. In Chapter 7, I simulate one-dimensional laser-cooling of beryllium ions in ALPHA and present the results. Laser ablation of metals is briefly reviewed in Chapter 8 and the relevance of ion generation using laser ablation is examined considering the constraints imposed by the ALPHA trapping environment. The design and construction of a laser-ablated beryllium ion source is described in Chapter 9. A stand-alone apparatus to test and qualify the ion source is also described in this chapter. In Chapter 10, laser ablation measurements with aluminum and beryllium are analyzed. Finally, in Chapter 11, the integration of the beryllium ion source into ALPHA is discussed, and future ion-positron experiments are outlined.

1.4 Author Contributions

The ALPHA experiment consists of approximately 50 individuals from ~ 20 different academic institutions. This makes the collaboration small enough that a graduate student quickly becomes involved with all aspects of the experiment. I was no exception, and I spent most of 2014, 2015 and 2016 taking shifts to run ALPHA, running experiments, debugging issues, maintaining cryogenics and developing hardware. During the first half of 2014, I spent a significant portion of my time upgrading the quench-protection system for the superconducting magnets used in ALPHA for trapping antihydrogen. In 2016, I was assigned as “Run Coordinator” for 3 weeks, responsible for the operation and maintenance of ALPHA and the implementation of its physics program.

During my time at ALPHA, the progress in antihydrogen physics has resulted in some significant achievements. In particular, I was involved in the following:

- Improving the limit on the charge of antihydrogen [36]
- Observing the 1S-2S transition in trapped antihydrogen [37]
- Observing the hyperfine spectrum in antihydrogen [38]
- Accumulation of antihydrogen atoms [39]
- Developing SDREVC to stabilize electron and positron plasma parameters [40]
- Adiabatic cooling of antihydrogen atoms [41]

Although I spent a significant amount of time on these experiments, they were not the main focus of my PhD and are not described in this thesis.

As part of a smaller group within ALPHA, I lead the development, commissioning and testing of a beryllium ion source that will be used to sympathetically cool positrons in ALPHA. This work is discussed in Chapters 6 through

11. It was guided by my supervisor Prof. Niels Madsen and his postdoc Dr. Daniel Maxwell, from whose assistance I have greatly benefited.

Chapter 2

ALPHA Apparatus

The aim of the ALPHA collaboration is to perform precision measurements on trapped antihydrogen. In order to achieve this, we must be able to produce, confine and manipulate charged particles, namely positrons and antiprotons from which antihydrogen atoms are made, and electrons that assist in the antihydrogen formation process. Additionally, we must have the capacity to produce and trap sufficient numbers of antihydrogen atoms that make experimental investigation viable. Consequently, developing techniques that result in improved antihydrogen trapping is an ever-present theme in all of ALPHA's activities.

In this chapter, I describe the main components of the ALPHA apparatus that pertain to the production of positrons, electrons and antiprotons in ALPHA. I also describe the Penning-Malmberg traps that confine these particles, and the neutral magnetic trap that holds the cold antihydrogen atoms. Finally, I describe the detectors that are used to identify and characterize electrons, positrons, antiprotons and antihydrogen atoms in ALPHA.

2.1 Theory

In the presence of static electric and magnetic fields, charged particles experience a Lorentz force given by the equation

$$\vec{\mathbf{F}} = q(\vec{\mathbf{E}} + \vec{\mathbf{v}} \times \vec{\mathbf{B}}) \quad (2.1)$$

By choosing a suitable configuration of electric and magnetic fields, charged particles may be confined in a small region of space; such a device is referred to as an “electromagnetic trap”. One such configuration, known as a Penning trap, consists of a pure quadrupole electric field and a superposed uniform axial magnetic field. The electric potential and the axial magnetic field are given by

$$\phi(r, z) = \frac{V_0}{2d^2} \left(z^2 - \frac{r^2}{2} \right) \quad (2.2)$$

$$\vec{\mathbf{B}} = B_0 \hat{z} \quad (2.3)$$

Here d is a geometric parameter of the electrodes that provide the electric potential. Typically, hyperbolic electrodes are used to provide a pure quadratic electric potential, and a solenoid magnet produces the uniform axial magnetic field. In such a trap, the charged particle trajectory is described by a bounce motion in the axial direction, and a superposition of two circular motions (magnetron and cyclotron motion) in the radial direction, as shown in Figure 2.1.

When the axial confining potentials are quadratic, charged particles follow a simple harmonic trajectory in the axial direction. The angular frequency of oscillation ω_z is given by

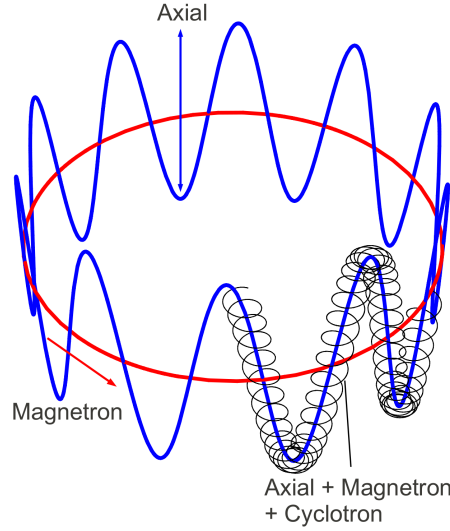
$$\omega_z = \sqrt{\frac{qV_0}{md^2}} \quad (2.4)$$

where m and q are the mass and electric charge respectively of the particle.

In the absence of any radial electric potential, charged particles would simply exhibit cyclotron motion in the radial direction, with the cyclotron angular frequency ω_c given by

$$\omega_c = \frac{qB_z}{m} \quad (2.5)$$

Figure 2.1: A sketch of the motion of a charged particle in an ideal Penning trap. Along the axis, the particle exhibits a bounce motion (blue). In the radial direction, the center-of-mass motion is a circular (magnetron) orbit (red). On top of that, the particle also revolves around the magnetron orbit, known as cyclotron motion (black). From [42]



The quadratic electric potential modifies this trajectory however and induces charged particles to exhibit a combination of cyclotron and magnetron motion. The modified cyclotron angular frequency ω'_c is given by

$$\omega'_c = \frac{\omega_c}{2} + \sqrt{\frac{\omega_c^2}{4} - \frac{\omega_z^2}{2}} \quad (2.6)$$

while the magnetron angular frequency ω_m is given by

$$\omega_m = \frac{\omega_c}{2} - \sqrt{\frac{\omega_c^2}{4} - \frac{\omega_z^2}{2}} \quad (2.7)$$

It should be noted that the magnetron frequency is real only in the case when

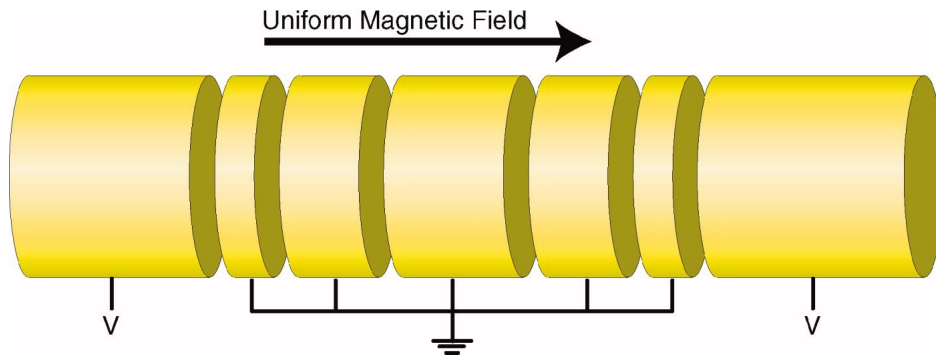
$$\omega_c > \sqrt{2}\omega_z \quad (2.8)$$

When the trapping potential is such that $\omega_c < \sqrt{2}\omega_z$, the magnetic field is

unable to overcome the effect of the radial quadratic potential that pushes charged particles away from the trap center. As a result, particles are no longer confined radially and the trap is no longer stable.

Equation 2.8 provides the stability criterion for ideal Penning traps with hyperbolic electrodes. In practice, hyperbolic electrodes are usually only employed for precision measurements with single or few particles. When working with larger numbers, hyperbolic electrodes would hinder particle loading; instead, a stack of open-ended cylindrical electrodes is used as shown in Figure 2.2. Such a device, known as a Penning-Malmberg trap, has axial and radial motion slightly modified from that in an ideal Penning trap, but particle confinement is still achievable with the application of the correct potentials on the end electrodes, and the presence of a sufficiently strong magnetic field.

Figure 2.2: A simplified sketch of a Penning-Malmberg trap. The trap is composed of a stack of hollow cylindrical electrodes on which electric potentials can be applied, and a uniform axial magnetic field provided typically by a solenoid magnet. From [43]



The Catching Trap (Section 2.4.4), Atom Trap (Section 2.3) and Positron Accumulator (Section 2.4.1) all contain Penning-Malmberg traps for the confinement of electrons, antiprotons and/or positrons. In each case, a strong solenoid provides the magnetic field for radial confinement, and electrode potentials confine the particles axially. The specific geometries of these traps are described later in this chapter.

Neutral particles, such as antihydrogen atoms, do not experience a Lorentz force and therefore cannot be confined in a Penning trap. Antihydrogen atoms formed by mixing positron and antiproton plasmas would therefore reach the trap walls shortly after formation and promptly annihilate. Luckily, antihydrogen atoms, like hydrogen atoms, can be trapped magnetically due to their non-zero magnetic dipole moment. The antihydrogen magnetic dipole moment arises from the vector sum of the following

- the intrinsic magnetic moment of the positron
- the magnetic moment due to the orbital motion of the positron around the antiproton
- the intrinsic magnetic moment of the antiproton

In the ground state, the antihydrogen magnetic moment comes primarily from the positron magnetic moment, which is approximately equal to 1 Bohr magneton, $1\mu_B = \frac{e\hbar}{2m_e}$.

A magnetic dipole moment $\vec{\mu}$ in the presence of an external magnetic field $\vec{B}$ has a potential energy U given by the equation

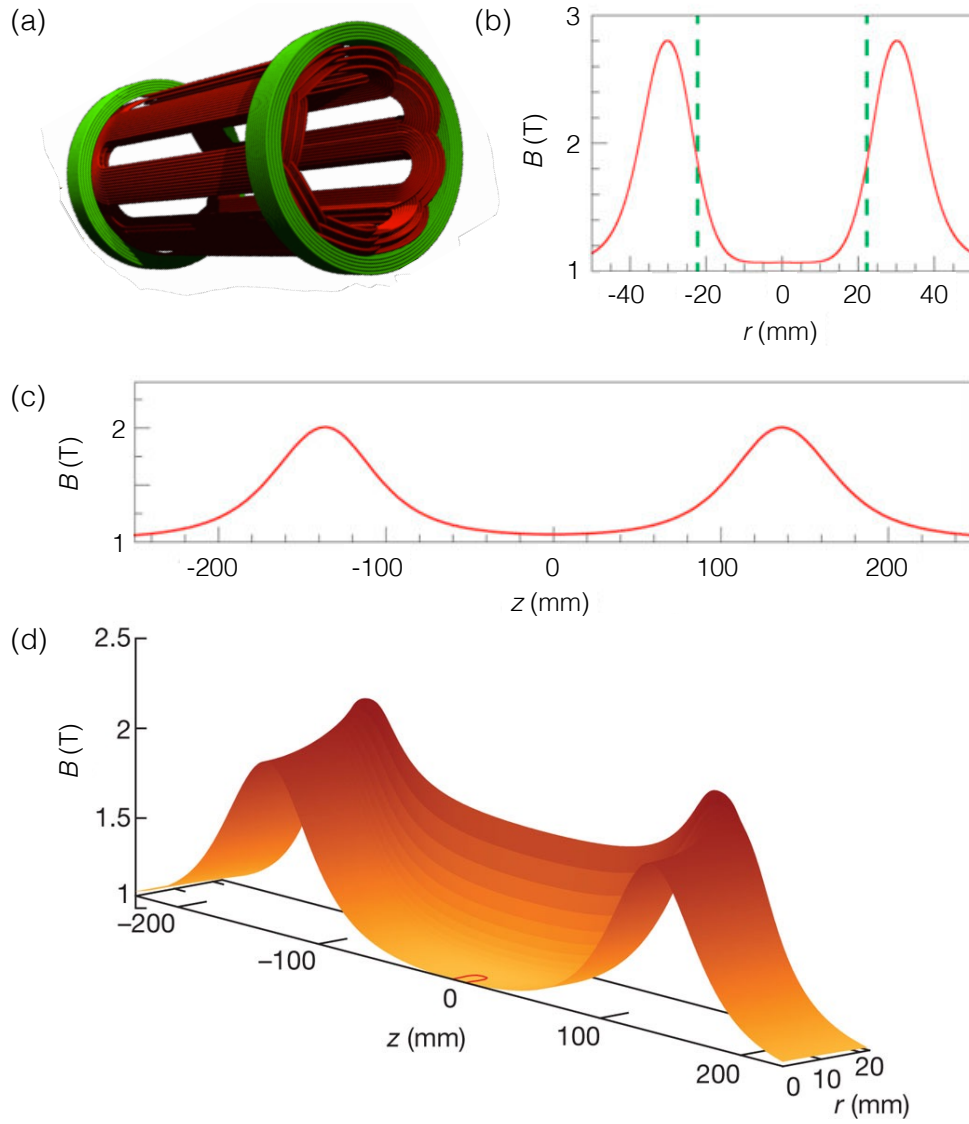
$$U = -\vec{\mu} \cdot \vec{B} \quad (2.9)$$

In the case of antihydrogen in the ground state, the magnetic dipole may either be aligned or anti-aligned with the external magnetic field. We have

$$U = -\vec{\mu}_{\bar{H}} \cdot \vec{B} = \pm\mu_B B \quad (2.10)$$

If there is a magnetic gradient, moments aligned with the magnetic field ($\mu_{\bar{H}} = +\mu_B$) move towards regions of higher magnetic field, and moments anti-aligned ($\mu_{\bar{H}} = -\mu_B$) move towards regions of lower magnetic field. These are called “high-field seekers” and “low-field seekers” respectively. By configuring a magnetic field with a magnetic minimum in the center of the trap, the atoms with a positive magnetic moment will experience a force towards the center. One such magnetic configuration is created by having two coaxial

Figure 2.3: (a) A schematic of a magnetic minimum trap consisting of two mirror coils (green) and an octupole (red). (b) The radial magnetic field profile at ($z = 0$ mm) due to the octupole magnet. (c) The axial magnetic field profile (at $r = 0$ mm) due to the two mirror coils. (d) The three-dimensional magnetic field profile of the full magnetic trap with a magnetic minimum in the trap center.



magnetic mirror coils at each end of the trap to provide axial confinement, and a multipole magnet in between them to confine the atoms radially. Figure 2.3 shows such a magnetic trap with two mirror coils and an octupole magnet. When the magnets are fully energized, a magnetic minimum is created in the trap center. Any high-field-seeking antihydrogen atoms that are created within the trap would drift away from the trap center and eventually fall out of the trap. Any low-field-seeking atoms will remain confined as long as their energies are below what is required to overcome the trapping fields. This ability to confine low-energy and low-field-seeking atoms forms the basis for trapping antihydrogen in the ALPHA apparatus.

2.2 Apparatus Overview

Figure 2.4: A simplified schematic of the upgraded ALPHA apparatus. Antiprotons coming from the AD (left of the schematic) are caught by the Catching Trap. Positrons generated by the radioactive sodium source are accumulated in the Positron Accumulator. Both particle species are transferred into the Atom Trap where they are processed and subsequently mixed to form antihydrogen. The apparatus includes ports for lasers to access the trapping volume.

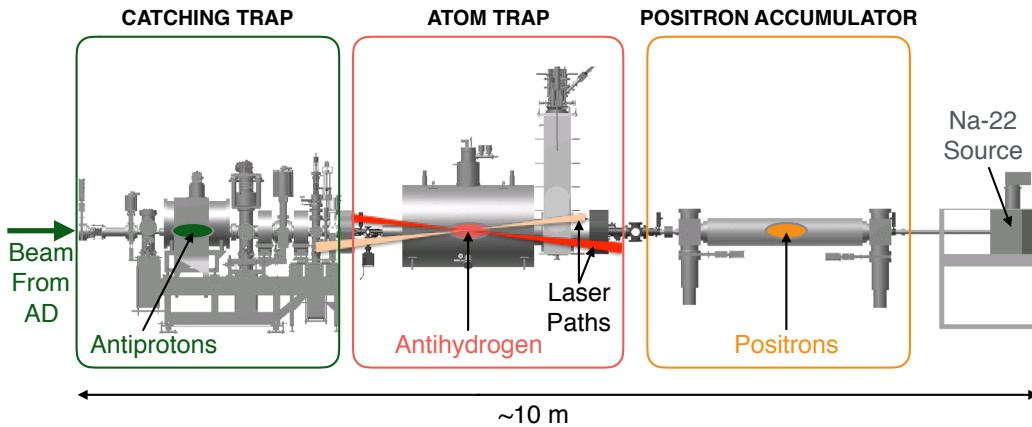


Figure 2.4 shows a simplified layout of the ALPHA apparatus in its current form. The ALPHA apparatus was originally designed with antihydrogen production and trapping in mind [32]. Having achieved that objective in 2010 [44], it was modified to include several important upgrades:

1. Construction of an independent Catching Trap for antiprotons
2. Installation of vacuum windows to provide optical access to the antihydrogen trapping region
3. Construction of a new cryostat for the superconducting magnets in the Atom Trap
4. Upgrade of the quench protection system and magnet control software for the trap magnets

As a result of these upgrades, the ALPHA apparatus effectively split into three independent sections:

- Catching Trap: to capture, accumulate and manipulate antiprotons
- Atom Trap: to recapture, process and mix positrons and antiprotons for antihydrogen production
- Positron Accumulator: to capture and accumulate positrons

The modular nature of these traps means that each component can be accessed independently for easy operation and maintenance. Additionally, positrons and antiprotons can be accumulated in their respective traps while measurements are carried out in the Atom Trap. The ability to operate the three traps in parallel reduces the dead-time between consecutive antihydrogen production cycles, allowing for physics measurements to be made in less time and with reduced sensitivity to slow drifts.

In Section 2.4.4, the process of capturing antiprotons in the Catching Trap using high-voltage electrodes is described in detail. One of the main motivations of constructing an independent Catching Trap was to avoid the use of high voltage in the Atom Trap. The electronic noise associated with high-voltage results in unwanted heating of charged plasmas and is detrimental for the production of antihydrogen. With the Catching Trap upgrade, the ALPHA Atom Trap is shielded from these undesirable effects.

The other main motivation for constructing the Catching Trap was to create space in the Atom Trap area for the installation of laser windows. The optical access results from the provision of four unobstructed laser paths that go through the center of the Atom Trap. Each laser path may be used independently, making them available for laser-spectroscopy, laser-cooling of atoms, laser-cooling of ions, or any other laser use one might imagine.

The modular nature of the Catching Trap will also be beneficial once ELENA becomes operational. With the expected factor of ~ 100 increase in the number of antiprotons captured due to the lower energy of the ELENA beam, the Catching Trap may be used as an antiproton accumulator or reservoir, making significantly more antiprotons available for antihydrogen experiments.

The ALPHA apparatus, in particular the Atom Trap, has several superconducting magnets as part of the antihydrogen production and trapping equipment. The niobium-titanium wire used for these magnets is only superconducting at very low temperatures. For this reason, the ALPHA trap magnets are kept in a bath of liquid helium which has a temperature of 4.2 K. Liquid helium is a precious commodity however, and upgrading the Atom Trap cryostat was necessary in order to reduce the amount of helium required. Prior to the upgrade, the Atom Trap helium consumption during normal operation was in excess of 30 liters per hour. After the upgrade however, the nominal consumption is only 18 – 20 L/h. The cryostat and trap upgrades have also resulted in an improved trap environment: the trap electrodes reach temperatures of ~ 7 K compared to 15 K previously. This contributes towards an improved trap vacuum as well as colder positron and electron plasmas.

When the upgraded ALPHA apparatus came online in 2013, only four of the eight Atom Trap magnets were initially commissioned (the solenoid of the Recatching Trap, the two end mirror coils of the Mixing Trap, and the octupole; see Figure 2.7). In early 2015, the remaining four magnets were commissioned, and a significant portion of the quench protection system and the associated magnet control software was updated. This was followed by

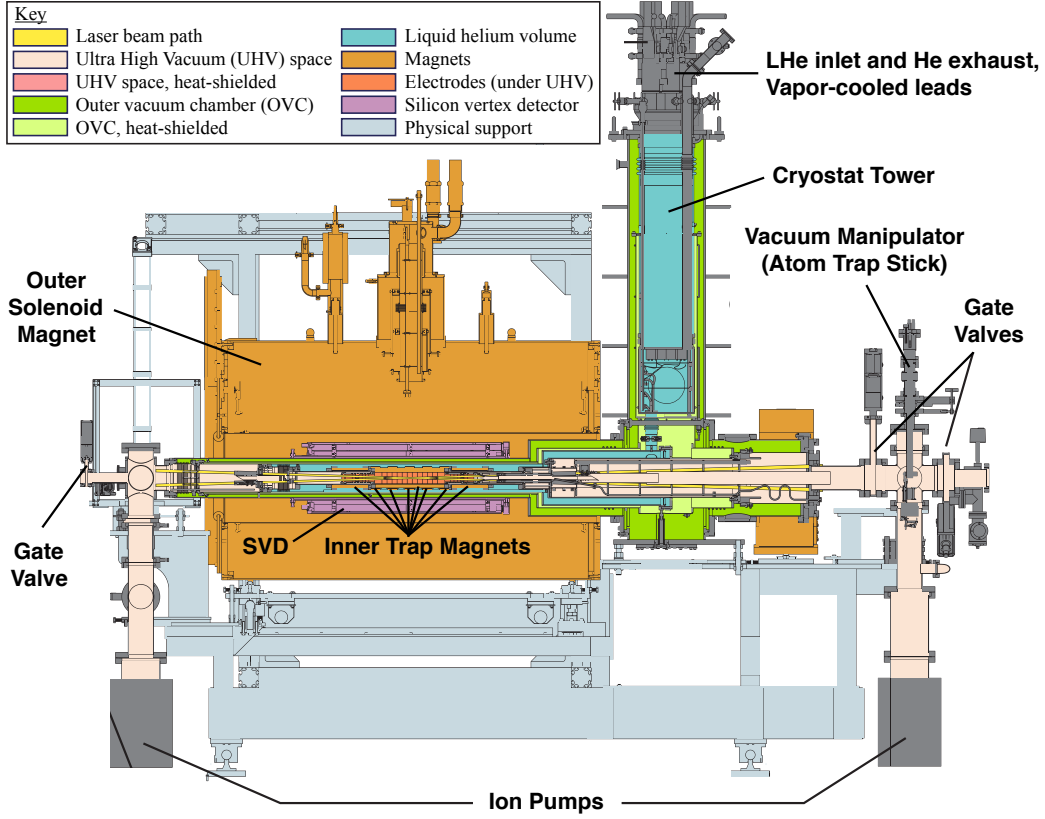
further modifications that now allow for an improved stabilization of the magnet currents (to within 0.5 A in 10 minutes) and a better control of the on-axis magnetic field, an important step towards precision antihydrogen spectroscopy with microwaves and lasers. The availability of the three central mirror coils also enables a flattening of the central magnetic minimum field of the antihydrogen trap, resulting in a larger region of magnetic field homogeneity in the center of the trap. The upgrade also makes available the booster solenoid (see Section 2.3, Figure 2.7) at the positron end of the Atom Trap, allowing for better plasma manipulation techniques to be deployed on positrons.

2.3 The ALPHA Atom Trap

The central region of the ALPHA Apparatus where antihydrogen atoms are created and trapped is known as the Atom Trap. Figure 2.5 shows a detailed schematic of the Atom Trap region. Several key features of the Atom Trap are immediately visible:

1. The antihydrogen trapping region is kept under an Ultra-High Vacuum (UHV) with the assistance of vacuum pumps, vacuum gauges and gate valves at each end of the Atom Trap. UHV conditions are necessary in order to have stable charged particle plasmas and avoid particle loss due to collisions with background gas particles. Under UHV conditions, the antiproton and positron annihilations with background gas particles are negligible. The antiproton and positron plasmas can therefore be stably loaded and prepared and subsequently mixed to form antihydrogen atoms. The UHV also ensures that the rate of annihilation of antihydrogen atoms with background particles is small.
2. A liquid helium cryostat, which is effectively split into a horizontal and a vertical section, surrounds most of the central UHV space. The vertical section, also known as the “cryostat tower”, houses the liquid helium inlet and helium vapor exhaust system at the top of the tower.

Figure 2.5: Detailed layout of the Atom Trap. All components are drawn to scale.



The current leads for the superconducting trap magnets, also located at the top of the tower, are cooled by the exiting helium vapor. The superconducting magnets themselves are located in the horizontal section of the cryostat, which is sandwiched between the UHV space and the outer-vacuum chamber (OVC). This is also shown in Figure 2.6. The liquid helium keeps the trap magnets cold (at 4.2K), while the OVC insulates the colder inner structures from the rest of the apparatus.

3. A Silicon Vertex Detector (SVD) radially surrounds the antihydrogen trapping region outside the OVC (see Figure 2.6). The SVD is the main detector used to reconstruct the location and timing of antihydrogen annihilations in the trapping volume. It is discussed further in Section 2.5.4.

4. One of the largest features of the Atom Trap is a solenoid magnet that surrounds the horizontal section of the cryostat as well as the SVD. The 1 T field from this superconducting magnet forms the magnetic basis of the Penning-Malmberg trap where antiprotons and positrons plasmas are prepared and mixed to form antihydrogen.

Figure 2.6: A schematic cross-section of the Atom Trap showing different volumes at the core of the antihydrogen trapping region. Adapted from [42].

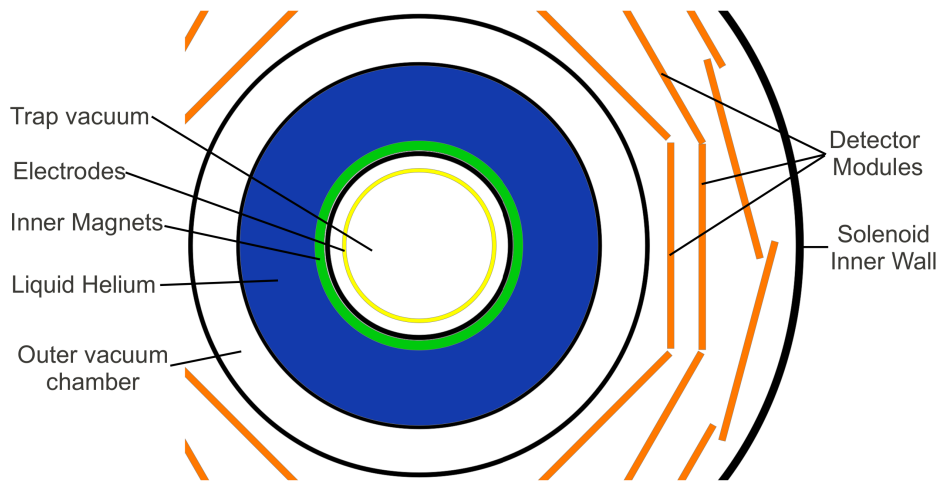


Figure 2.7: This diagram shows a simplified version of the ALPHA Atom Trap. The Penning-Malmberg trap is formed from the electrodes (gold) and a 1 T solenoid magnet (not shown). The five mirror coils (green) and the octupole (grey) create the magnetic trap. Also shown are two booster solenoids (blue) that assist in antiproton and positron preparation. Adapted from [42]

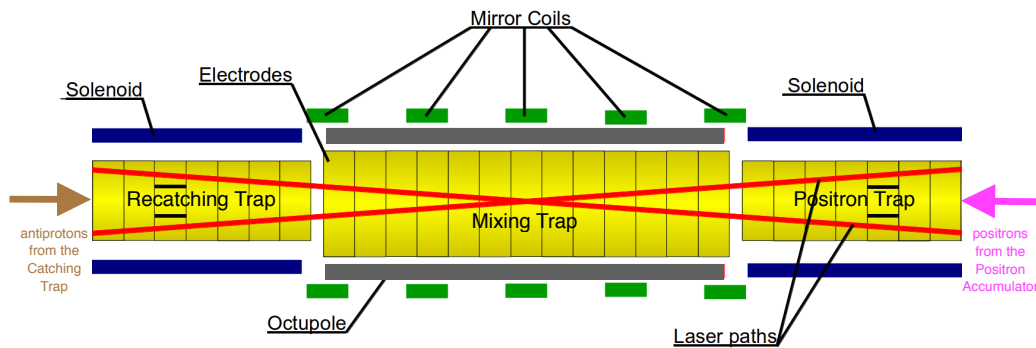


Figure 2.7 shows a schematic of a Penning-Malmberg trap superimposed with a magnetic trap, a representation of what is currently used in ALPHA. The

Penning-Malmberg trap is effectively divided into three parts: the Recatching Trap which receives antiprotons from the Catching Trap, the Positron Trap which collects positrons from the Positron Accumulator, and the Mixing Trap, where antiprotons and positrons are mixed together to form antihydrogen. Such a system offers a great deal of flexibility to trap and manipulate the positrons and antiprotons for antihydrogen production. If the magnetic trap (at least two mirror coils and the octupole) is energized at the time of formation, antihydrogen atoms with the correct spin and sufficiently low energy remain confined in the trap as well.

When trapping antihydrogen, the trap depth ΔU is simply given using Equation 2.10

$$\begin{aligned}
 \Delta U &= U_{max} - U_{min} \\
 &= \mu_{\bar{H}} (B_{max} - B_{min}) \\
 &= \mu_B \Delta B \\
 &= \frac{e\hbar}{2m_e} \Delta B \\
 &= (0.67k_B) \Delta B
 \end{aligned} \tag{2.11}$$

where ΔB is the difference between the highest and lowest magnetic field strength (in tesla). This means that a magnetic field difference of 1 tesla produces a trap depth (in units of kelvin) of only 0.67 K, i.e. antihydrogen atoms with an energy less than $0.67k_B$ will remain trapped, while the rest will overcome the confining magnetic potential and annihilate. The typical ALPHA trap depth is 0.5 K.

2.4 Plasma Sources

As mentioned previously, ALPHA uses positron and antiproton plasmas and mixes them together to make antihydrogen. In addition, electrons are also used in antiproton preparation processes discussed in Chapter 3. All of these plasmas require an independent source that can generate plasmas controllably, repeatably and reliably.

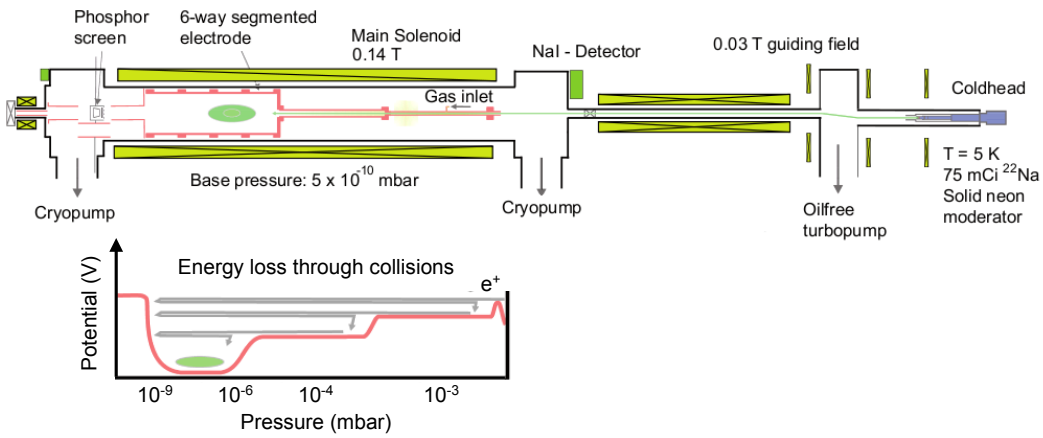
2.4.1 Positron Accumulator

Positrons in ALPHA originate from a source of ^{22}Na which undergoes radioactive decay



Na-22 has a half-life of 2.6 years, and therefore emits a reproducible amount of positrons on a day-to-day basis. The emitted positrons typically have an energy range in the hundreds of keV, and therefore need to be slowed down before they can be trapped. At ALPHA, this is accomplished in a multi-stage device known as the Positron Accumulator, a schematic for which is shown in Figure 2.8.

Figure 2.8: A schematic of the Positron Accumulator at ALPHA. Positrons emitted from a radioactive Na-22 source pass through a moderator of solid neon, and accumulate in a three stage Surko-type trap [45]. The path of the positrons is shown in green. The trap consists of a large solenoid magnet (0.14 tesla) with three stacks of hollow cylindrical electrodes of increasing radii at each stage. Nitrogen buffer gas is introduced into the trapping volume which naturally forms a pressure gradient (see lower panel of graphic). By applying electrode potentials as shown, the positrons lose energy as they traverse the trapping region and slowly accumulate in the third stage. In about 150 seconds, the positron accumulator is able to collect up to 80 million positrons. Adapted from [46]



In the first stage, the emitted positrons pass through a layer of solid neon,

which serves as a moderator to slow down the positrons. While most of the positrons annihilate within the moderator volume, a small fraction ($< 1\%$) make it through with significantly reduced kinetic energies (~ 80 eV). These lower energy positrons are magnetically guided into the Penning-Malmberg trap of the accumulator, comprising a 0.14 T magnetic field and 7 cylindrical electrodes.

The first two electrodes are narrow, with a diameter of 12.7 mm each. The third electrode has a diameter of 30.5 mm, and the last four electrodes are the widest, at 200.7 mm each. Pure nitrogen gas is introduced into the trap through a gas inlet near the second electrode. The gas slowly diffuses downstream (to the left). The three-stage electrode stack with varying diameters causes a pressure gradient to build up within the trapping volume. Nitrogen acts as a buffer gas and positrons bouncing axially back and forth in the trap lose energy to the gas molecules. By applying step-potentials to the electrodes (as shown in the lower panel of Figure 2.8), the positrons over time are confined in the third stage of the trap. Since the fourth electrode is 6-way segmented, one may employ the rotating wall technique to compress the positron cloud further (see Section 3.2).

Once accumulated, the positrons may be sent to the Atom Trap following a series of steps. First, the nitrogen gas is pumped out from the Accumulator in order to avoid vacuum contamination of the Atom Trap. Next, the mechanical valve separating the Accumulator from the Atom Trap is opened, and a series of electrode potentials are applied to eject the positron cloud which is guided by magnetic fields from a copper coil to the Atom Trap where they are recaptured.

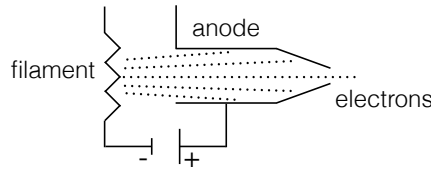
2.4.2 Electron Gun

Electrons are also frequently used in the ALPHA traps for a variety of reasons

- to develop plasma manipulation techniques (See Chapter 4)

- to measure magnetic fields in the trap
- to capture, cool and compress antiprotons (See Chapter 4)
- to diagnose issues with the trapping apparatus

Figure 2.9: Simplified schematic of an electron gun. Current passing through the filament heats it up, resulting in the emission of electrons which are accelerated towards the anode and exit the device through an aperture.



The electrons needed for these purposes are produced by a device known as an Electron Gun, a schematic for which is shown in Figure 2.9. The electron gun contains a barium-oxide filament. When current is passed through this filament, it heats up and electrons are produced via thermionic emission. Applied electric fields accelerate the electrons away from the filament and magnetic fields guide the electrons into the trap. In ALPHA, there are two electron guns currently in use, one for the Catching Trap and one for the Atom Trap.

2.4.3 Antiproton Decelerator

The Antiproton Decelerator (AD) is an accelerator ring at CERN that provides experiments with bunches of antiprotons for experimentation. A beam of energetic protons (26 GeV/c momentum) from the Proton Synchrotron (PS) is incident on a fixed target of high density iridium metal. The resulting high energy collisions create a plethora of particles of a wide energy range. A small fraction of these particles are antiprotons. By using a dipole magnet, antiprotons at 3.5 GeV/c are mass selected and injected into the AD.

Once inside the AD ring, the antiproton beam is decelerated in several stages, first to 2 GeV/c, next to 300 MeV/c, and finally to 100 MeV/c [27]. The

beam is decelerated using resonant RF cavities. By applying an oscillating electric field in these cavities, antiproton motion is opposed and the particles are slowed down. Unfortunately deceleration leads to an increase in the beam's relative momentum spread and requires deceleration to be intermixed with phase-space cooling in order to maintain a narrow beam.

Cooling of the antiproton beam occurs using two main techniques. At the higher GeV energies, stochastic cooling [47] is employed which involves correcting the orbit of antiprotons during each turn in the ring. A pickup circuit detects the displacement of an antiproton from the ideal orbit and applies a correction to its orbit further downstream. A correction applied to one particle may disturb another one however, so the correction amplitude is calculated in such a way that the overall beam size and momentum spread is reduced at each turn. This maintains the narrowness of the beam as it is being decelerated.

At the lower MeV energies, stochastic cooling is less effective and electron cooling [48] is used instead. A beam of cold electrons is overlapped with the antiproton beam. Coulomb interactions between the two species allow the antiprotons to lose energy and cool down. At the end of the process, the electrons are ejected leaving a pure beam of antiprotons.

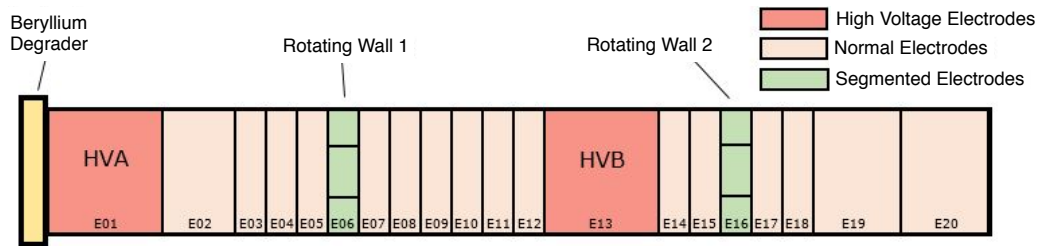
The whole deceleration and cooling cycle takes about 100 seconds, with typically 3×10^7 particles remaining. The 5.3 MeV antiproton beam is collected into ~ 100 ns bunch and ejected to one of several experiments connected to the AD.

2.4.4 Catching Trap

The ALPHA experiment is one of the recipients of the antiproton beam from the AD. These antiprotons need to be caught, cooled and compressed before they can be used for antihydrogen production. As mentioned, the device ALPHA uses to catch and prepare the antiprotons is known as the Catching

Trap, a Penning-Malmberg trap with a 3 tesla solenoidal magnetic field and a stack of 20 electrodes as shown in Figure 2.10.

Figure 2.10: A schematic of the Catching Trap electrodes at ALPHA. There are several types of electrodes: high-voltage electrodes are used to capture the incoming energetic beam of antiprotons; normal electrodes are used for routine plasma manipulations; segmented electrodes are used for Rotating Wall Compression (Section 3.2). Adapted from [49]

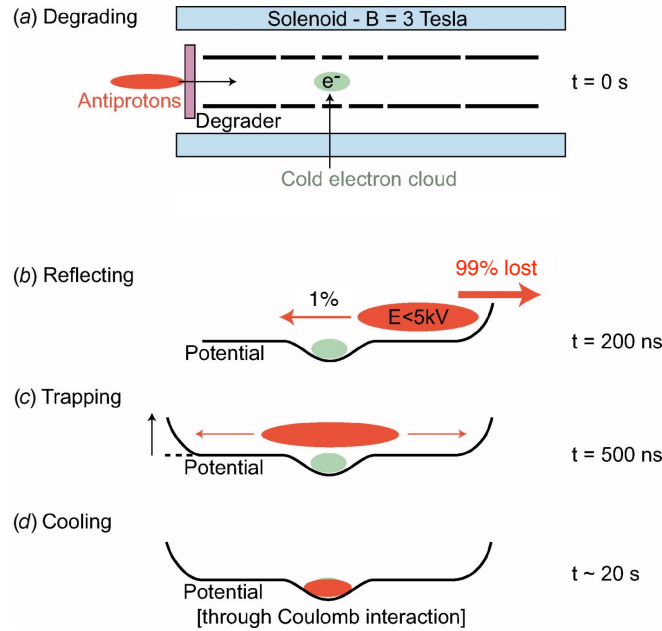


Electrodes E01 and E13 are connected to special high voltage pulsers that can output up to 5 kV. They are labeled HVA and HVB in Figure 2.10 and are used for catching the antiproton bunch arriving from the AD. The steps needed to catch and cool antiprotons are shown in Figure 2.11.

Before the antiprotons arrive, the trap is pre-loaded with a cloud of approximately 85 million electrons. The electrons efficiently cool in the 3 T magnetic field by cyclotron radiation (see Section 3.1) and quickly reach an equilibrium temperature. Electrode HVB is also energized to 5 kV.

At time $t = 0$ s, the antiproton bunch arrives and passes through a 0.2 mm thick set of beryllium and aluminum foils which “degrades” the beam i.e. it drastically slows down the many of the transiting antiprotons via Coulomb collisions. As a consequence, a small fraction of antiprotons end up with an attenuated axial energy < 5 keV. The antiprotons then enter the Catching Trap and travel through the electrode stack until they reach HVB. Approximately 99% of the antiprotons have an energy above 5 keV and are lost, while the rest are reflected back towards the degrader. Before the antiprotons reach the degrader, the second high-voltage electrode HVA is energized

Figure 2.11: An illustration of the steps used to capture the beam of antiprotons in the Catching Trap. (a) 5.3 MeV antiprotons from the Antiproton Decelerator pass through a 0.2 mm thick degrader, resulting in many antiprotons losing energy via collisions with the degrader material. Pre-loaded electrons in the center of the trap lose energy via cyclotron radiation and cool down to ambient temperature. (b) A 5 kV blocking potential reflects the slow-moving antiprotons. The more energetic particles overcome the blocking potential and are lost. (c) ~ 500 ns after the arrival of antiprotons, a second 5 kV blocking potential is erected before the reflected antiprotons escape, thereby trapping them. (d) The trapped antiprotons are cooled by Coulomb collisions with the self-cooling electrons. In 20 seconds, approximately 70% of the antiprotons have thermalized. Adapted from [43].



within a few 100 ns, trapping the antiprotons. Collisions with the pre-loaded electrons cool down the antiprotons, while the electrons are self-cooled by cyclotron radiation. Typically the antiprotons are allowed to cool for 20 seconds before the high-voltage electrodes are turned off, resulting in $\sim 110,000$ trapped antiprotons.

At this moment, antiprotons are ready for further plasma processing, details of which will be discussed in the next two chapters. Some of these processes use special electrodes: E06 and E16, are segmented into six azimuthal sec-

tions in order to implement the Rotating Wall technique (see Section 3.2). Furthermore, custom-made amplifiers are connected to each individual electrode for the application of up to ± 140 V potentials for plasma manipulation and transport.

2.5 Detectors

The process of producing antihydrogen atoms requires antiproton, positron and electron plasmas to have very specific parameters during the plasma preparation and mixing processes. The ability to diagnose the number, size, density and temperature of the plasmas is crucial in order to reliably and repeatably create and trap the anti-atoms. ALPHA employs a number of detectors for this purpose. They are briefly described in the remainder of this chapter.

2.5.1 Faraday Cup

The number of charged particles in an electromagnetic trap can be destructively measured by releasing them onto a Faraday Cup. The Faraday Cup is usually a piece of conducting metal such as aluminum. As the charged particles strike the Faraday Cup, a voltage builds up. By knowing the capacitance of the Faraday Cup and measuring the voltage signal, we can calculate the total charge of the arriving particles, and thereby infer the total number of particles that hit the metal surface.

In ALPHA, the degrader is used as a Faraday Cup to count electrons in the Catching Trap, and a piece of aluminum metal on a vacuum manipulator is used for counting both electrons and positrons in the Atom Trap. These Faraday Cups are only sensitive for electron and positron numbers $> 10^6$. For this reason, they are not used to count the number of antiprotons (typically between 10^4 and 10^5). Instead, external scintillators (see Section 2.5.3) are used to infer the number of antiprotons.

2.5.2 MCP and Phosphor Screen

In ALPHA, a detector composed of a micro-channel plate (MCP) and a phosphor screen is used to obtain the radial profile of charged particle plasmas. An MCP is essentially an array of micrometer-sized channels that behave as electron-multipliers. When a charged particle hits inside of one of these channels, secondary electrons are produced. As these electrons travel along the channel, they may strike the wall again and again, eventually producing a cascade of secondary electrons that exit the back end of the channel as an electron beam. The magnitude of the electron-cascade is a function of the energy of the original charged particle as well as the applied potential difference between the front and back of the MCP. The electrons that leave the MCP then strike the phosphor screen which emits photons that can be imaged by a CCD camera.

The image recorded by the camera mimics the radial density profile of the incoming beam of charged particles. Therefore, by releasing a trapped plasma onto an MCP/phosphor detector, the radial distribution of the charged particles can be obtained. By releasing a charged particle plasma at a Faraday Cup and another similarly-prepared plasma at an MCP, the intensity of the MCP image can be calibrated to give an absolute measure of the total charge as well. In ALPHA, a calibrated MCP is routinely used to count as few as 10^4 electrons or positrons, two orders of magnitude lower than what the Faraday Cup is capable of.

The MCP is also used to obtain the temperature distribution of trapped electron and positron plasmas in ALPHA. By slowly lowering the confining electrode potentials (relative to the axial bounce frequency), the hottest (most energetic) particles are allowed to escape first and hit the MCP, while the colder ones arrive later. The arrival of the electron-cascade at the phosphor screen, which follows the arrival of the plasma at the MCP, is extracted using a custom-designed pickoff circuit. The signal from the phosphor pickoff

is then used to obtain the temperature profile of the plasma.

With the ability to provide information about the number (e^- , e^+), radial distribution (e^- , e^+ , $\bar{p}$) and temperature (e^- , e^+) of the plasma, the MCP detector is an indispensable diagnostic tool to measure the plasma properties in the ALPHA Penning-Malmberg trap. Further details about the use of an MCP and phosphor screen to detect antiprotons, positrons and electrons in ALPHA can be found in Reference [50].

2.5.3 Scintillators

The annihilation of an antiproton with a proton or neutron results in the production of ionizing radiation in the form of charged pions [51]. When ionizing radiation passes through a scintillator, the energy of the radiation is absorbed and re-emitted as photons. Photomultiplier-tubes (PMTs) attached to scintillators can then amplify the light and convert it into a measurable electrical signal. In order to avoid spurious signals from electrical noise, only electrical signals that exceed a certain threshold are used.

One way to minimize the effect of electrical noise is to use scintillators in pairs. A positive count is only registered when an electrical signal above threshold is observed in both scintillators of the pair at the same time. While a pion that passes through one scintillator is likely to pass through the second scintillator as well, the electrical noise in a scintillator is random and therefore simultaneous signals above threshold are unlikely. In this way the signal from a pion is correctly identified while electrical noise is rejected.

In ALPHA, five pairs of scintillators are used, one pair on each side of the Catching Trap, one pair on each side of the Atom Trap, and one pair next to the Atom Trap Stick. In this way, the annihilation of antiprotons along the length of the trap can be monitored. By releasing a trapped antiproton plasma towards the Faraday Cup or MCP, the total number of antiprotons in that plasma can be calculated from the total number of counts registered

by the scintillator pairs. Similarly, by slowly releasing the antiprotons and looking at the time distribution of the counts, the temperature of the antiproton plasma is extracted.

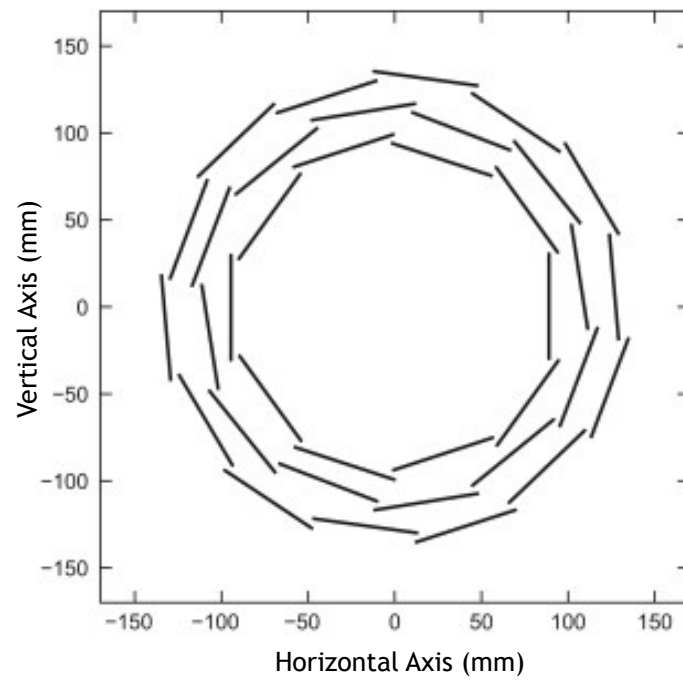
While the MCP is used to ascertain the radial distribution of antiprotons, the scintillators provide the necessary number and temperature information to fully characterize the antiproton plasma.

2.5.4 Silicon Vertex Detector

The detectors discussed so far in this thesis are used to obtain information regarding charged particles only i.e. electrons, positrons and antiprotons. The main detector ALPHA uses to detect antihydrogen atoms is the Silicon Vertex Detector (SVD). The SVD is divided into two halves (left and right), each composed of 36 independent silicon modules and distributed in three layers as shown in Figure 2.12. Each module is composed of 512 silicon strips, half of which run along the length of the module and the other half run across the module. When an antiproton or antihydrogen annihilation occurs in the trapping volume, a pion passing through one of the strips may deposit sufficient energy and register a “hit” in that particular module. When hits are registered in two separate layers, a particle track can be reconstructed.

A sophisticated software algorithm is responsible for reconstructing particle tracks as a result of registered hits in the SVD. When two or more particle tracks intersect, the algorithm reconstructs the estimated location and time of annihilation. For antihydrogen annihilations, the SVD has a vertex reconstruction efficiency of $\sim 68\%$ and a vertex resolution in cylindrical coordinates (r, ϕ, z) of 0.8 cm, 18° , and 0.5 cm respectively. By choosing a suitable set of cuts, false positives from cosmic rays can be suppressed to a rate of 47 ± 2 mHz. Further details of the SVD hardware and software are given in References [46] and [52].

Figure 2.12: A cross-section of the Silicon Vertex Detector showing the distribution of silicon modules in three layers. The inner, middle, and outer layers have 10, 12, and 14 modules respectively. Adapted from [52]



Chapter 3

Plasma Cooling and Manipulation

In the ALPHA experiment, antihydrogen atoms are created by mixing positrons with antiprotons. If one were to directly mix the cloud from the Positron Accumulator with the antiproton beam caught in the Catching Trap, the resultant rate of antihydrogen formation would be extremely low, and any atoms created would be too energetic to be trapped in the ALPHA magnetic trap (see Chapter 5 for details).

One must therefore apply sophisticated techniques borrowed from plasma physics in order to optimize the number, density and temperature of the positron and antiproton plasmas before they are mixed to form low-energy trappable antihydrogen. In this chapter, I outline the processes used at ALPHA for preparing antiprotons and positrons prior to antihydrogen formation. In particular, I describe SDREVC (see Section 3.4), a recent technique developed at ALPHA which has been a mainstay in our hallmark achievement of the first observation of the 1S-2S atomic transition in trapped antihydrogen [37].

3.1 Cyclotron Cooling

Charged particles deflected by an external magnetic field emit electromagnetic radiation perpendicular to the magnetic field, known as cyclotron radiation. The power radiated is given by the Larmor formula [53]:

$$P = \frac{q^2}{6\pi\epsilon_0 c^3} a^2 \quad (3.1)$$

where q is the particle's charge, a is the centripetal acceleration due to cyclotron motion and c is the speed of light. As power is radiated, the particle loses energy and slows down. This process of charged particles losing energy due to cyclotron radiation is known as cyclotron cooling.

By plugging in the centripetal acceleration $a = \frac{Bqv}{m}$ and using the kinetic energy formula $E = \frac{1}{2}mv^2$, one may solve the energy differential equation $\frac{dE}{dt} = -P$ using the previously stated Larmor formula. The resultant time-evolution of the energy of a charged particle undergoing cyclotron cooling can be written as

$$E(t) = E_0 e^{-\frac{t}{\tau_c}} \quad (3.2)$$

$$\tau_c = \frac{3\pi\epsilon_0 c^3}{q^4} \frac{m^3}{B^2} \quad (3.3)$$

The energy of the particle undergoes an exponential decay where τ_c is the characteristic decay time, also known as the cyclotron cooling time [53].

For electrons or positrons in a 1 T magnetic field, we have $\tau_c \sim 3$ s, making cyclotron motion an efficient cooling mechanism in tesla level magnetic fields. Due to the significantly larger mass of antiprotons however, the cyclotron cooling time is on the order of 10^{10} seconds, making the process extremely inefficient for cooling antiprotons or heavier ions. This is the reason why at the time of catching antiprotons from the AD, the Catching Trap is pre-loaded with a cloud of electrons that cyclotron-cool themselves in the 3 T magnetic field of the Penning-Malmberg trap.

In ALPHA, cyclotron cooling of electrons and positrons is crucial in the preparation of antiproton and positron plasmas. The steps during which

cyclotron cooling is used for antihydrogen production are detailed in Chapter 4.

3.2 Rotating Wall Compression

A charged particle plasma in a Penning-Malmberg trap rotates azimuthally about the trap axis due to the $\mathbf{E} \times \mathbf{B}$ drift. If the plasma has radially uniform density, the frequency of rotation f is approximately given by the equation

$$f = \frac{nq}{4\pi\epsilon_0 B} \quad (3.4)$$

where n is the plasma number density, q is the particle charge and B is the axial magnetic field [54]. A plasma with uniform density therefore rotates like a rigid rotor in steady state, with all particles having the same rotational frequency all the way to the edge of the plasma.

Due to the cylindrical symmetry of a Penning-Malmberg trap, the total angular momentum L_θ of a charged plasma is a conserved quantity [55].

$$L_\theta = \sum_j (mr_j^2\omega_j + qA_\theta r_j) \quad (3.5)$$

Here the first term is the mechanical angular momentum and the second term is the canonical angular momentum. In the presence of a strong uniform magnetic field (where $A_\theta = Br/2$), the canonical angular momentum term dominates. The total plasma angular momentum L can then simply be approximated as

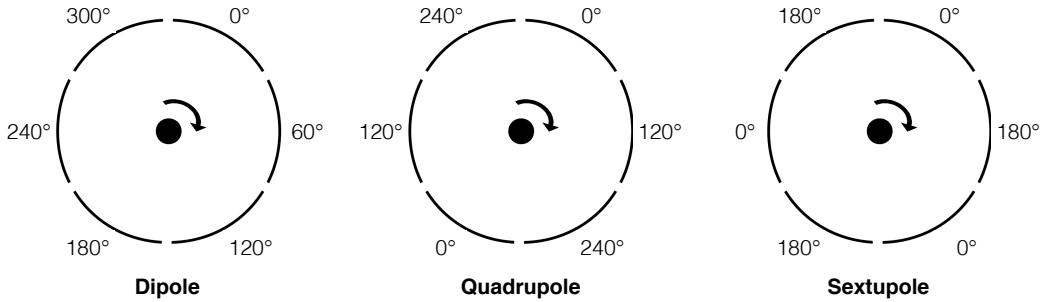
$$L \simeq \sum_j qA_\theta r_j = \frac{qB}{2} \sum_j r_j^2 \quad (3.6)$$

where r_j represents the radial location of the j^{th} particle from the rotation axis. As a result, a particle drifting radially outward would result in other particles moving radially inward in order to conserve angular momentum. As long as the diameter of the confining electrodes is sufficiently large, a plasma rotating in an ideal Penning-Malmberg trap remains perpetually confined.

In practice, Penning-Malmberg traps used in experiments are far from ideal. Magnetic field inhomogeneities, electrode imperfections and collisions with background gas result in the outward radial transport of particles, resulting in a slow plasma expansion. This exerts a torque against the plasma motion and leads to a reduction in the plasma rotational frequency and angular momentum.

The loss in plasma angular momentum may be counteracted by applying a torque that balances the drag. In fact, by applying a torque that exceeds the drag, the angular momentum of the plasma may be increased even further. Since the rotational frequency of the plasma is proportional to the plasma density (Eq. 3.4), the increase in angular momentum results in a plasma that is more dense than before. A plasma may therefore be compressed further using this technique, commonly referred to as Rotating Wall Compression.

Figure 3.1: Sketch of typical electrodes used for Rotating Wall Compression in ALPHA. The electrodes are split into six segments. The relative phase of the electric potential on each segment is labeled. From left to right: dipole, quadrupole and sextupole rotating electric fields. Adapted from [42].



One way to create the required torque is to apply a sine-wave electric potential to an azimuthally segmented electrode, as shown in Figure 3.1. The applied potential V_i on each segment is given by

$$V_i = V_0 \sin(\omega_{RW}t - \theta_i) \quad (3.7)$$

These potentials generate a rotating electric field that applies a torque on the plasma and compresses it. As long as the drive amplitude V_0 is sufficiently large (known as the Strong Drive Regime [56]), the plasma rotation

frequency synchronizes with the applied rotating-wall frequency ω_{RW} . Consequently, a higher rotating-wall frequency results in denser, more compressed plasmas. Rotating Wall Compression is an integral technique for the production of antihydrogen at ALPHA (see Chapter 4).

3.3 Evaporative Cooling

In a collection of particles, the temperature is defined as the average of the kinetic energy of the ensemble. If the most energetic particles are removed from the ensemble, the average kinetic energy of the collection decreases, thus lowering the temperature. The process of overall temperature reduction by the selective removal of energetic particles is known as evaporative cooling (EVC). Evaporative cooling is routinely used in atomic physics to cool a gas of atoms to the ground state and create Bose-Einstein condensates [57].

During the preparation of positron and antiproton plasmas for antihydrogen synthesis, ALPHA routinely employs the evaporative cooling technique in order to reduce the plasma temperature. The charged particles are contained in a shallow potential well in the Penning-Malmberg trap. The electrode potential at one end of the well is then slowly lowered, allowing the most energetic particles to escape along the axis. Since the particles escape in the axial direction, the total angular momentum of the remaining plasma remains conserved (Eq. 3.6). As a result, a hole in the plasma is created which is quickly filled by the adjacent particles via collisions. Conservation of angular momentum then necessitates other particles to move radially outward, resulting in plasma expansion. This technique simultaneously reduces the plasma's particle density and space-charge. This feature of evaporative cooling is exploited in a new technique known as Strong-Drive-Regime-Evaporative-Cooling, or SDREVC (see Section 3.4).

3.4 SDREVC

As mentioned in Section 3.2, a uniform charged particle plasma in a Penning-Malmberg trap rotates like a rigid rotor in steady state. The plasma can be approximated as a solid cylinder of radius r_p , length l_p and uniform charge density $\rho = nq$. On the condition that $l_p \gg r_p$ (for example in the case of a cold ($T \sim 0$), magnetized plasma), the on-axis potential ϕ_p due to the plasma is given by

$$\phi_p = \frac{nqr_p^2}{4\epsilon_0} \left[1 + 2 \ln\left(\frac{R_w}{r_p}\right) \right] \quad (3.8)$$

where R_w is the inner radius of the electrodes of the Penning-Malmberg trap [58].

The SDREVC technique involves performing rotating-wall compression on the non-neutral plasma in the Strong Drive Regime (SDR) while simultaneously evaporatively cooling it (EVC). It is based on the realization that the rotating-wall frequency fixes the plasma density n in Eq. 3.8 above, whereas the shape of the trap well and final electrode potential during EVC fixes the plasma potential ϕ_p . Simultaneous application of both techniques would then uniquely determine the full set of final plasma parameters. Therefore, by applying SDREVC, one can theoretically obtain plasmas of the desired number and density from a variety of initial plasma parameters.

SDREVC was first realized in ALPHA using electron plasmas [58]. Subsequent implementation of the technique with positrons has been successfully used to obtain stable and reproducible positron plasmas for antihydrogen production [37].

3.5 Adiabatic Cooling

In a Penning-Malmberg trap, charged particles are confined axially in an electric potential well. The potential $U(z)$ at the center of the confining well can be approximated as a harmonic potential of the form $U(z) = \frac{1}{2}k_z z^2$.

Under this approximation, a charged particle in such a well experiences an axial restoring force $F_z = -k_z z = -m\omega_z^2 z$ where m is the particle's mass and ω_z is the axial bounce frequency. The total axial energy E_z of the particle is given by

$$E_z = \frac{1}{2}mv_z^2 + \frac{1}{2}m\omega_z^2 z^2 \quad (3.9)$$

For a given energy, the particle traces out an elliptical orbit in phase-space where the area of the ellipse is given by $2\pi \frac{E_z}{\omega_z}$. For slow changes in ω_z , the area of the ellipse remains constant. Therefore, the ratio of the particle's energy to the frequency of oscillation i.e. $\frac{E_z}{\omega_z}$ is an adiabatic invariant [59].

By slowly lengthening the confining potential well, one can adiabatically decrease ω_z . As a result the particle's axial energy and consequently its temperature is reduced [60]. This process is known as adiabatic cooling.

The above is a simple description of one-dimensional adiabatic cooling of a single particle. For low density plasmas, the single-particle approximation holds true since Coulomb interactions can be neglected. In this approximation, the temperature scaling is given by $T \sim \omega_p$ i.e. the temperature varies linearly with the plasma frequency.

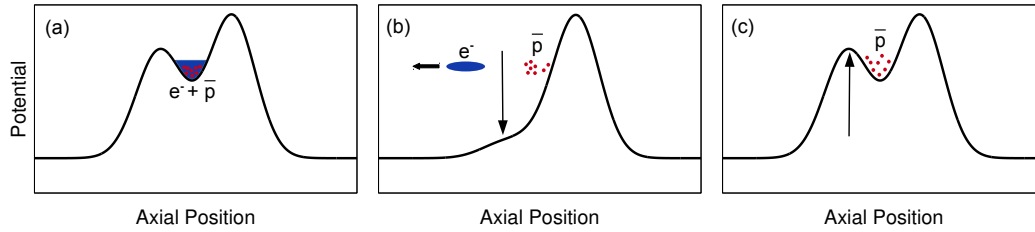
In the case of high densities however, collisions and space-charge effects need to be taken into consideration. A rigorous theoretical three-dimensional treatment [61] predicts a scaling of $T \sim \omega_p^{\frac{4}{3}}$ in such a scenario, which agrees reasonably with experiments performed on antiprotons by Gabrielse *et al.* [62].

The adiabatic cooling of positrons (Section 4.1.2) is now included in the latest antihydrogen production sequence at ALPHA. Furthermore, ALPHA has also adiabatically cooled trapped atoms of antihydrogen, the first ever demonstration of adiabatic cooling in any system of neutral antimatter.

3.6 Electron Kicking

Earlier in Section 2.4.4, we saw that the capture and cooling of the beam of antiprotons from the AD requires the use of a pre-loaded cloud of electrons. The result is a mixed antiproton-electron cold plasma trapped in a potential well. In order to create antihydrogen, the electrons need to be removed from the trap prior to mixing of the antiprotons with the positrons. If the electrons are not removed, it could lead to a loss of positrons due to electron-positron annihilations as well as the formation of positronium. Furthermore, the space-charge of the electrons can heat up the positrons, resulting in a reduced antihydrogen formation and trapping rate (see Chapter 5).

Figure 3.2: A schematic representation of an electron-kick. (a) The mixed antiproton-electron plasma is trapped due to confining electrode potentials. (b) One of the confining potentials is quickly lowered, allowing the lighter electrons to escape. (c) Before the heavier antiprotons can escape, the confining potential is reestablished, leaving a pure antiproton plasma. From [42]



One mechanism of removing electrons from the mixed antiproton-electron plasma utilizes the mass difference of the two species. The procedure, referred to as an “electron-kick”, is illustrated in Figure 3.2. First, a voltage pulse that lowers one of the confining electrode potentials is applied, leaving the constituents of the plasma temporarily unconfined. The lighter electrons have larger velocities and escape rapidly while the heavier antiprotons are still in the trap. If the confining electrode potential is reestablished fast enough (on the order of ~ 100 ns), the antiprotons would not have sufficient time to escape. The antiprotons remain trapped, and a pure antiproton plasma is achieved.

A negative consequence of electron-kicking is the unwanted heating of the antiproton plasma. The quick application of a voltage pulse and the sudden removal of the electron space-charge causes the remaining plasma to heat up. A slightly modified strategy can be devised to minimize this heating. By tuning the amplitude of the electron-kick, one can select a certain fraction of the electrons to be removed instead of the entire electron cloud. The remaining electrons would then cool via cyclotron radiation, and also sympathetically cool the heated antiprotons via Coulomb collisions.

Once the mixed plasma re-cools, the confining potentials are lowered to account for the reduced space-charge after the initial electron removal. Now, an electron-kick of a much smaller magnitude is required, resulting in a much smaller level of plasma heating than before.

In ALPHA, the electron-kicking process is divided into several stages with the duration and amplitude of each kick carefully optimized. The first electron-kick is tuned so that $\sim 10^6$ electrons remain. This kick is immediately followed by rotating-wall compression which both compresses and heats up the plasma. The electrons help to re-cool the hot plasma via cyclotron emission and sympathetically cool the antiprotons. Once the plasma is cold, two more electron-kicks are applied: the initial one leaves just over 10^5 electrons, while the final one gets rid of all of the electrons. The end product is a pure and dense cloud of antiprotons with no remaining electrons and a low temperature.

Chapter 4

Antihydrogen Formation and Trapping

Furnished with the description of the ALPHA trapping apparatus from Chapter 2 and armed with the necessary plasma manipulation techniques from Chapter 3, we are now ready to look at detailed steps involved in synthesizing and trapping antihydrogen atoms.

In this chapter, I outline the sequence of steps in the Catching Trap and Atom Trap that result in trapped antihydrogen at ALPHA. I describe the preparation of positrons and antiprotons and the various ways in which these prepared plasmas may be mixed. Finally, I discuss the fate of the created anti-atoms, the ones that remain trapped and the ones that escape and annihilate.

4.1 Particle Preparation

The plasma manipulation techniques mentioned in the previous chapter offer us an enormous parameter space to work with in our quest to obtain the most suitable plasmas for antihydrogen formation. By carefully studying how each of these techniques affect the number, density, and temperature of the positron and antiproton clouds, we are able to develop a well-defined

process that gives us our desired plasma features.

4.1.1 Preparing Antiprotons

Antiprotons coming from the Antiproton Decelerator are caught and cooled in the Catching Trap as described in Section 2.4.4. Prior to capturing the antiprotons, SDREVC is performed on the pre-loaded electron plasma to fix the density and space-charge parameters: ~ 85 million electrons and a radius of ~ 9 mm. Upon the arrival of the antiproton beam, high voltage (5 kV) electrodes are used for their catching, and the pre-loaded electron plasma enables the cooling of the antiprotons. It takes only 20 seconds to cool $> 70\%$ of the antiprotons to < 5 keV kinetic energy. Lowering the high-voltage at this time ejects the remaining un-cooled antiprotons. Once the antiprotons have been caught and cooled, further manipulations are carried out on the mixed electron-antiproton plasma in order to optimize the antiproton number, density and temperature.

The plasma is compressed by applying the rotating-wall technique. Since the rotating-wall applies a torque and does work on the plasma, this heats up the antiprotons. The presence of electrons is therefore necessary to provide a mechanism for the antiprotons to cool down.

Once compressed to the desired radius, electrons are partially ejected from the trap by applying an electron-kick. The amplitude of the first kick is such that not all of the electrons are removed. Rotating-wall compression is applied to compress the remaining plasma even further. By having an intermediate electron-kick between two plasma compressions, significantly more compressed antiprotons can be achieved. This is because the minimum radial size of the plasma is limited by maximum density achievable with the Rotating Wall technique. Removing a fraction of electrons thereby permits further compression of the mixed plasma.

After the second compression, the remaining electrons are removed with two

more electron-kicks of decreasing amplitude. The first kick removes most of the electrons but also heats up the plasma. The leftover ($\sim 100,000$) electrons help in re-cooling the antiprotons before they themselves are ejected to leave a pure plasma of antiprotons, ready to be transferred to the Recatching Trap (see Figure 2.7). At this stage, we typically have 110,000 antiprotons with a radius of 0.2 mm and a temperature of ~ 400 K.

By quickly ($\sim 2\mu\text{s}$) removing the potential on the confining electrode, the antiprotons are ballistically sent to the Recatching Trap, where they are caught and cooled with pre-loaded electrons in a similar fashion to what was done previously in the Catching Trap. The booster solenoid is energized and the plasma then undergoes similar treatment as before, with rotating-wall compression, electron-kick, another compression, and two more electron-kicks. After the final kick that removes all of the electrons, we are again left with a cold (~ 100 K), compressed (0.4 mm) cloud of $\sim 110,000$ antiprotons.

After de-energizing the booster solenoid, the antiprotons are shuffled from the Recatching Trap to the Mixing Trap where they are evaporatively cooled, leaving $\sim 90,000$ antiprotons ~ 1.0 mm in radius and ~ 40 K in temperature. The antiprotons are now ready for mixing.

4.1.2 Preparing Positrons

While the antiprotons are being prepared in the Catching Trap, positrons are collected for 150 seconds in the Positron Accumulator as described in Section 2.4.1. At the end of the accumulation cycle, the nitrogen gas in the accumulator is pumped out and about 80 million positrons are injected into the Positron Trap (see Figure 2.7).

The positron cloud, with a pulse length of ~ 200 ns, is dynamically caught, with a trapping efficiency of about 10%. The positrons are allowed to cool for 30 seconds in the 3 T magnetic field created by the booster solenoid. Since the transfer to the Positron Trap is a lossy mechanism, the density

and number of positrons reaching the Atom Trap may vary from shot-to-shot. In order to alleviate the effects of this fluctuation further downstream, SDREVC is performed on the positron plasma, with the typical choice of parameters leaving us with ~ 3 million positrons. The positrons are further compressed using rotating wall compression, and further allowed to cyclotron-cool and thermalize with the trap environment before the booster solenoid is de-energized. After waiting 20 seconds in the 3 T magnetic field of the booster solenoid, the positrons nominally reach temperatures of ~ 50 K.

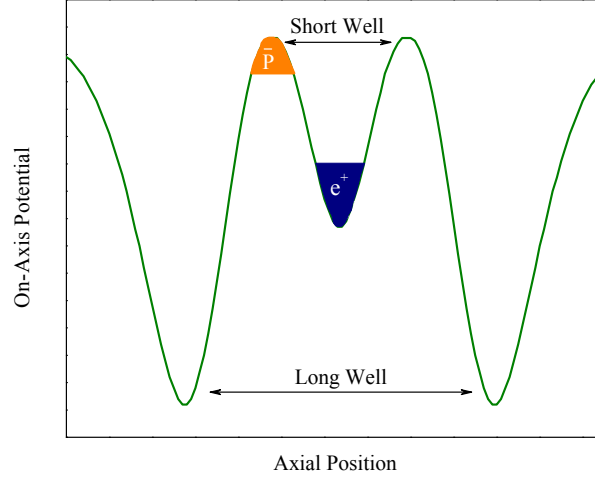
Next, the positron plasma is moved to the center of the Mixing Trap and the neutral atom trap is energized, which has the undesirable consequence of heating up the positrons due to the magnetic field inhomogeneities introduced by the octupole magnet. The positrons are therefore adiabatically cooled and then evaporatively cooled, with the final positron cloud containing ~ 2 million positrons at an average temperature of 18 K. Both the antiprotons and positrons are now ready to be mixed to create antihydrogen.

4.2 Plasma Mixing

Once both antiproton and positron plasmas are ready, they need to be brought together in adjacent wells before mixing can take place. Since the two species have opposite charges, their simultaneous confinement can be accomplished in what is known as a “nested potential well” [63].

Figure 4.1 shows a schematic of a nested potential well typically used at ALPHA. It comprises a long inverted well that can confine negative particles, in particular antiprotons, and a smaller inner well to contain oppositely-charged positrons. This configuration results in the creation of two “side-wells”. In practice, before mixing takes place, antiprotons are kept in one of the side wells while positrons are situated in the middle well. If antiprotons are given sufficient axial energy to overcome the potential barrier, they escape the side well and pass through the positron plasma in the middle. If a travel-

Figure 4.1: Illustration of a nested potential well that confines positrons in a short well and antiprotons in an inverted long well. From [64]



ing antiproton captures a positron, a neutral atom of antihydrogen is created.

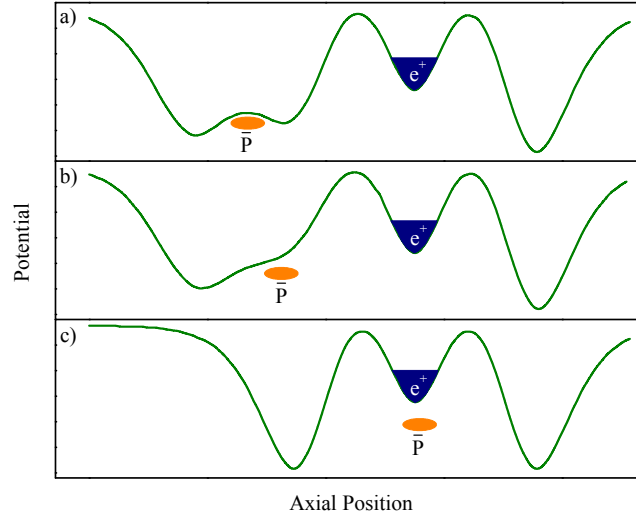
There are several different ways of mixing that result in the formation of antihydrogen in a nested potential well. Each of these mechanisms produces antihydrogen atoms with slightly different properties, in particular slightly different kinetic energy distributions and temperatures.

4.2.1 Simple Mixing

The first cold antihydrogen atoms were synthesized by launching antiprotons into a cold positron plasma [65]. The antiprotons and positrons were kept in adjacent potential wells in a modified nested potential well as shown in Figure 4.2. In this configuration, antiprotons were held at a relative potential higher than the positron plasma barrier. By lowering the antiproton confining potential, the antiprotons were launched into the positron plasma, causing some of the positrons to be bounded to the incoming antiprotons, thus making neutral antihydrogen atoms.

This simple technique produced antihydrogen atoms that were too hot to trap

Figure 4.2: Illustration of a simple mixing sequence. (a) Antiprotons are trapped in a potential well where the potential energy is higher than the barrier of the positron well. (b) Antiprotons are allowed to escape into the nested well where positrons are located. (c) Antiprotons enter the positron plasma to create antihydrogen. From [64]



however [66]. This was primarily because antiprotons were injected with an energy of approximately 15 eV, creating antihydrogen atoms that were too energetic to be trapped in the 0.5 K magnetic trap. In subsequent years, more complex techniques, namely “autoresonant mixing” and “slow merge”, were developed to make colder, trappable antihydrogen atoms.

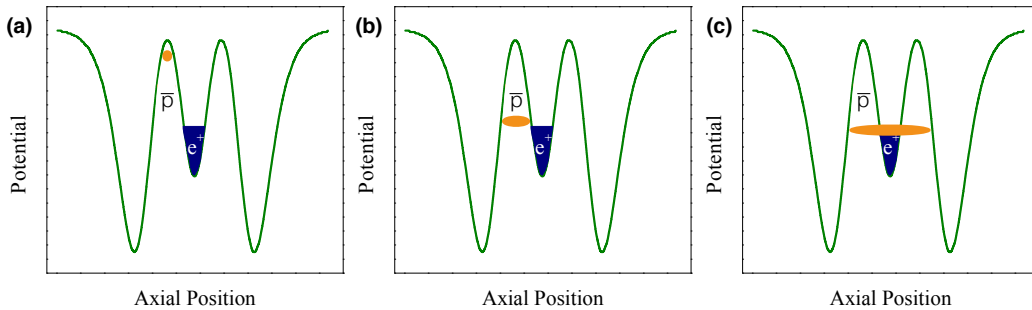
4.2.2 Autoresonant Mixing

Antiprotons trapped in a nested potential well have an axial bounce frequency in the hundreds of kHz range. By applying a swept frequency drive, the axial oscillations can be locked to the drive. When this happens, the oscillation amplitude can be controlled by varying the drive frequency, a phenomenon known as autoresonance [67].

Since the antiprotons in the nested potential well sit next to the positrons, autoresonant excitation can be used to raise the antiproton axial energy and

inject the cold antiprotons into the positron plasma. The process is illustrated in Figure 4.3. Initially, the cold antiprotons are confined in a side well and have axial energy that is too low to overcome the potential barrier. Next, a swept frequency drive is applied, which locks to the antiproton bounce frequency. As the drive frequency is lowered, it increases the axial potential of the phase-locked antiprotons. By slowly sweeping the drive, the potential energy of the antiprotons can be raised until the energy is large enough to overcome the confining potential. At this moment, antiprotons spill over into the positron well and fall out-of-lock with the drive frequency, thereby preventing the particles from acquiring any additional energy from the drive. The antiprotons mix with the positron plasma with very low axial kinetic energies, resulting in the production of cold antihydrogen.

Figure 4.3: Illustration of an autoresonant mixing sequence. (a) Antiprotons are trapped in a side well while positrons are trapped in the middle well. (b) Antiprotons are given axial energy using an autoresonant drive. (c) Antiprotons overcome the potential barrier and enter the positron plasma. Adapted from [64]

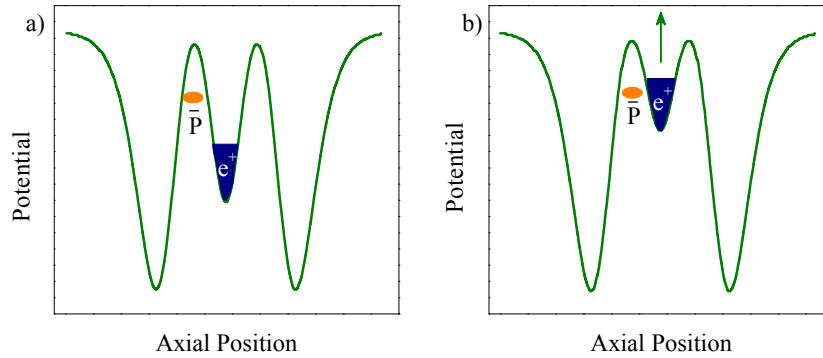


This technique was used by the ALPHA collaboration to successfully create the first trapped atoms of antihydrogen [44]. The success of autoresonant mixing in creating trappable anti-atoms comes from the fact that the cold antiprotons thermalize in the positron plasma before antihydrogen formation begins, leading to higher formation rates and much colder antihydrogen atoms (temperature dependence of antihydrogen formation and trapping is discussed in Chapter 5)

4.2.3 Slow Merge

A variation of the simple mixing technique of making antihydrogen atoms, slow merge (also known as incremental mixing) uses a simple potential manipulation that allows antiprotons to very slowly merge with the positron plasma. Figure 4.4 outlines the slow merge mixing process. As before, the antiprotons and positrons are initially held in two adjacent wells in a nested trap. Next, the positron well is very slowly raised until a point is reached where the most energetic antiprotons start spilling into the positron plasma to create antihydrogen. As the positron well continues to be raised, more and more antiprotons enter the plasma to create more and more antihydrogen.

Figure 4.4: Illustration of slow merge mixing. (a) Antiprotons are trapped in a side well while positrons are trapped in the middle in a typical nested potential well. (b) The relative potential difference between the two confining wells is slowly reduced until the antiprotons have sufficient energy to enter and mix with the positron plasma. From [64]



A peculiarity of the slow merge mixing is that the hottest positrons and antiprotons start mixing first, effectively evaporatively-cooling the remaining plasma. As the process continues, the colder particles start to mix as well. Therefore, one would expect that the coldest and most trappable antihydrogen is created at the latter stages of slow merge. At ALPHA, this feature is hard to measure however since evaporative-cooling due to slow merge is convoluted with lepton heating due to ambient heat loads.

The slow merge process at ALPHA is optimized to maximize overall antihydrogen trapping rate, i.e. the number of trapped atoms per attempt. Slow merge was the technique of choice for antihydrogen synthesis in a variety of experiments carried out by ALPHA in 2016 and 2017 (for example [37, 38]).

4.3 Antihydrogen Trapping

Antihydrogen atoms, like hydrogen atoms, are neutral [36]. As discussed earlier in Section 2.3, they cannot be trapped in a Penning-Malmberg trap, and can only be confined in a magnetic trap. Ground state antihydrogen atoms have a magnetic moment of ~ 1 Bohr magneton, which equates to a trap depth (in energy units of kelvin) of 0.67 K per 1 T well depth. The antihydrogen atoms formed using the plasma mixing techniques described above have an average kinetic energy in the tens of kelvin. Only a small fraction of atoms are cold enough to be trapped in the ALPHA magnetic trap, which has a trap depth of just ~ 0.5 K. In the most recent experiments carried out at ALPHA, as many as 50,000 antihydrogen atoms are created in each trapping cycle (taking into account detector efficiency), but only $\sim 0.02\%$ of the atoms are trapped.

4.3.1 Untrapped Antihydrogen

The vast majority of the antihydrogen atoms created are either high-field seekers (with magnetic moment parallel to the magnetic field), or have a kinetic energy larger than the trap depth. These atoms are not confined by the magnetic fields of the trap and quickly drift to the walls of the electrodes where they annihilate.

A specially designed silicon-vertex-detector (see Section 2.5.4) tracks the path of the annihilation products which can be used to reconstruct the location of the annihilation within the trap volume. The annihilation information of untrapped antihydrogen atoms is an extremely valuable dataset and has

several uses

- the total annihilation count (known as detector “triggers”) is used to estimate the total number of antihydrogen atoms formed in a given trapping cycle
- plasma mixing parameters are tuned against detector triggers in order to optimize the antihydrogen formation rates
- the spatial and temporal distribution of the annihilation vertices is used to diagnose unexpected issues, for example magnetic field inhomogeneities or trap asymmetries that result in annihilation hot spots in specific regions
- the annihilation dataset is used to train a multivariate analysis (MVA) algorithm that detects trapped hydrogen with high signal-to-noise ratio

A more detailed treatment of the silicon-vertex-detector can be found in References [46] and [52].

4.3.2 Trapped Antihydrogen

As long as the magnetic trap remains energized and the trap volume is in a pure vacuum, the low-field seeking antihydrogen atoms below 0.5 K energy remain trapped indefinitely. In practice, the vacuum is not perfect, and trapped atoms may annihilate with background gas particles over time. The atoms can be held for long enough however (up to 1,000 seconds in [68]) to allow for precision experiments [36, 37, 69] to be conducted with trapped antihydrogen.

Since only a few atoms are trapped per cycle, the experiments under investigation need to be repeated several tens or hundreds of times before the required statistical significance is achieved to make a scientific claim. This makes measurements susceptible to shot-to-shot variations as well as slow drifts in the trapping rate, increasing the measurement uncertainty. Increase in the antihydrogen formation and trapping rates would go a long way to

improve existing measurements and enable presently unfeasible experiments to take place in the future. The rest of the thesis is motivated by this desire to have a higher antihydrogen trapping rate at ALPHA and outlines some of the progress that has been made recently towards that end.

Chapter 5

Temperature Dependence of Antihydrogen

In the previous chapter, we have seen that antihydrogen atoms with a kinetic energy larger than the well depth remain untrapped after synthesis and quickly annihilate on the trap walls. In order to increase the trapping rate, one can employ one of two schemes:

1. Increase the well depth of the trap by changing the magnetic configuration or using stronger magnetic fields.
2. Increase the number of atoms created with a kinetic energy smaller than the well depth of the trap.

The first of these two is a technological as well as physics challenge: it requires a more sophisticated quench protection system, better mechanical and thermal anchoring of the trap magnets, and more robust plasma manipulation techniques in the presence of larger magnetic field gradients and inhomogeneities. The progress in this sector is a long-term objective beyond the scope of this thesis and will not be discussed henceforth.

The second scheme appears to be more attainable in the short term: it requires developing colder positron and antiprotons plasmas, and mixing them in a fashion that minimizes heating. In this chapter, we will discuss the theoretical underpinnings of antihydrogen formation, and how the temperature

of the constituent plasmas affects the rate of antihydrogen formation and the antihydrogen kinetic energy distribution that eventually determines the trapping rate.

5.1 Formation Mechanisms

At low energies, making antihydrogen is a simple matter of bringing an antiproton and a positron together, and removing the excess energy between them via a third particle to create a stable antihydrogen atom. This can be achieved in several ways:

$$\bar{p} + e^+ \rightarrow \bar{\text{H}} + \gamma \quad (5.1)$$

(Radiative Recombination)

$$\bar{p} + e^+ + \gamma \rightarrow \bar{\text{H}} + \gamma + \gamma \quad (5.2)$$

(Stimulated Radiative Recombination)

$$\bar{p} + e^+ + e^+ \rightarrow \bar{\text{H}}^* + e^+ \quad (5.3)$$

(Three-Body Recombination)

The first two of these processes, namely radiative recombination (RR) and stimulated radiative recombination (SRR), rely on photons to carry away the excess energy and momentum, allowing the antiproton and positron to bind together and form an antihydrogen atom.

In radiative recombination, the rate of antihydrogen formation depends on the densities of antiprotons and positrons, and the relative-velocity between the two particle species. Looking at the formation process in the rest-frame of an antiproton, the cross-section of formation Γ_{RR} can be written as

$$\Gamma_{\text{RR}} = 3 \times 10^{-17} \left(\frac{4.2}{T_{e^+}} \right)^{1/2} \frac{n_{e^+}}{\text{m}^{-3}} \text{ s}^{-1} \text{ per antiproton} \quad (5.4)$$

where n_{e^+} is the positron density and T_{e^+} is the positron temperature [63]. In a typical ALPHA environment, positron plasmas have a temperature of

20 K and a density of 10^{13} m^{-3} , resulting in $\Gamma_{\text{RR}} \approx 10^{-4} \text{ s}^{-1}$ per antiproton.

In stimulated radiative recombination, the antihydrogen production rate is also influenced by the laser. The presence of electromagnetic radiation can stimulate the positrons to transition to bound Rydberg states. Unlike spontaneous recombination, stimulated recombination is a resonant process as the conservation of energy requires that the photon energy be equal to the sum of the kinetic energy of the incoming positron and the positron binding energy. For antihydrogen, we have

$$\frac{m_e v_0^2}{2} + \frac{E_0}{n^2} = h\nu \quad (5.5)$$

where v_0 is the initial positron velocity in the antiproton rest frame, $E_0 = 13.6 \text{ eV}$ is the ionization potential of atomic (anti)hydrogen, n is the quantum level of the positron in the Rydberg atom, and ν is the photon frequency. Only positrons with velocities within the linewidth of the transition can participate in the stimulated recombination process to a given Rydberg state. The ratio g of induced (W^{ind}) to spontaneous (W^{spon}) recombination probability to a given quantum level is given by [70]

$$g = \frac{W^{\text{ind}}}{W^{\text{spon}}} = u(\nu) \frac{B}{A} \quad (5.6)$$

where

$$u(\nu) = \frac{P}{cF\Delta\nu} \quad (5.7)$$

is the spectral energy density of the radiation field of power P and photon beam cross section area F , and $\Delta\nu$ is the combined spectral width of the radiation field and the natural width of the bound atomic state. The Einstein coefficients for spontaneous and stimulated emission, A and B respectively, are related as follows

$$A = 8\pi \frac{h\nu^3}{c^3} B. \quad (5.8)$$

Using Equations 5.6, 5.7 and 5.8, we have

$$g = \frac{Pc^2}{F\Delta\nu 8\pi h\nu^3} \quad (5.9)$$

Since g is the enhancement factor for induced capture of positrons to a specific bound state, the total enhancement factor G for stimulated radiative

recombination is then given by the ratio of all induced recombination rates r_n^{ind} to all spontaneous recombination rates r_{spon} .

$$G = \frac{r_n^{ind}}{r_{spon}} = g \frac{r_n^{spon}}{r_{spon}} \quad (5.10)$$

The total cross-section of formation Γ_{SRR} is therefore

$$\begin{aligned} \Gamma_{SRR} &= (1 + G)\Gamma_{RR} \\ &= (1 + G) (3 \times 10^{-17}) \left(\frac{4.2}{T_{e^+}} \right)^{1/2} \frac{n_{e^+}}{\text{m}^{-3}} \text{ s}^{-1} \text{ per antiproton.} \end{aligned} \quad (5.11)$$

For positrons at low energies when $E_0 \gg \frac{m_e v_0^2}{2}$, the rate of spontaneous positron capture to a specific quantum level n is proportional to n^{-1} . At low energies, $h\nu \approx \frac{E_0}{n^2}$. Therefore, from Equation 5.9, $g \propto n^6$. Consequently, from Equation 5.10, G will be proportional to n^5 .

At the 4.2 K cryogenic temperatures normally used for antihydrogen formation experiments, Gabrielse *et al.* [63] approximates $G \approx 2 \times 10^{-5} n^5 I$ where I is the laser intensity. For large values of n i.e. longer laser wavelengths, infrared lasers of arbitrarily high intensities are available. The enhancement factor in these cases is primarily limited by the competing photoionization process due to the high laser intensities. For $n < 10$ however (corresponding to wavelengths $< 10\mu\text{m}$), producing high-intensity short-wavelength lasers and focusing them using durable optics is still a technological challenge and G is therefore limited in practice by the achievable laser intensity in the recombination volume.

In three-body recombination (TBR), a spectator positron assists the formation process. This usually happens when two positrons collide in the presence of an antiproton, resulting in one positron losing sufficient energy to bind to the antiproton, while the other flies away.

We can obtain an estimate of the three-body recombination rate by considering a positron plasma of density n_e and temperature T_e in the rest frame of a single antiproton. The classical distance of closest approach b between two colliding positrons is given by

$$b = \frac{1}{4\pi\epsilon_0} \frac{e^2}{k_B T_e}. \quad (5.12)$$

For a single positron in the plasma, the frequency of collisions is characterized by an impact parameter I which is the product of the cross-section of collision and the positron flux:

$$I \sim \pi b^2 n_e v_e \quad (5.13)$$

where $v_e = \sqrt{\frac{k_B T_e}{m_e}}$ is the velocity of the positron.

When a positron collision takes place, the probability P that a spectator positron is present within the collision volume to take away the excess energy is given by

$$P \sim n_e b^3. \quad (5.14)$$

Multiplying Equations 5.13 and 5.14, we obtain the scaling for TBR

$$\Gamma_{\text{TBR}} \sim (\pi b^2 n_e v_e)(n_e b^3) \sim n_e^2 b^5 v_e \sim n_e^2 T_e^{-\frac{9}{2}}. \quad (5.15)$$

Therefore, TBR is expected to be the dominant process of antihydrogen formation with cold and dense positron plasmas. In the presence of a strong magnetic field, the cross-section of formation Γ_{TBR} is given by

$$\Gamma_{\text{TBR}} = C (8 \times 10^{-24}) \left(\frac{4.2\text{K}}{T_{e^+}} \right)^{9/2} \left(\frac{n_{e^+}}{\text{m}^{-3}} \right)^2 \text{ s}^{-1} \text{ per antiproton} \quad (5.16)$$

where C is a function of magnetic field and temperature [71]. For the typical ALPHA positron parameters in a 1 T magnetic field, $C \sim 0.19$ [72]. This

results in a Γ_{TBR} of approximately 10^2 s^{-1} per antiproton, six orders of magnitude higher than Γ_{RR} for the same positron conditions. However, due to the $T_{e+}^{-\frac{9}{2}}$ scaling of TBR, this mechanism of formation is highly suppressed for hot positrons.

The $\Gamma_{\text{TBR}} \sim n_{e+}^2 T_{e+}^{-\frac{9}{2}}$ scaling for three-body recombination is only valid in the steady state when antiprotons are continually present in the positron plasma, allowing the weakly bound antihydrogen atoms sufficient time to collide with a positron and become strongly bound. Robicheaux [72] proposes that in the case of antiprotons transiting through a cloud of positrons without attaining steady state, and weakly bound antihydrogen atoms susceptible to ionization from the trap fields F_{trap} , the TBR scaling with temperature is modified. For an antiproton inside a cold positron plasma with density n_e and temperature T_e , the impact parameter from Equation 5.13 remains the same as before. The susceptibility of weakly bound antihydrogen atoms to ionize due to F_{trap} necessitates that a spectator positron be available for TBR within a distance $R \ll b$ of the antiproton, where R is related to F_{trap} by

$$F_{\text{trap}} \sim \frac{e}{4\pi\epsilon_0 R^2}. \quad (5.17)$$

In this case, the probability of finding a spectator positron within the collision volume to take away the excess energy is simply

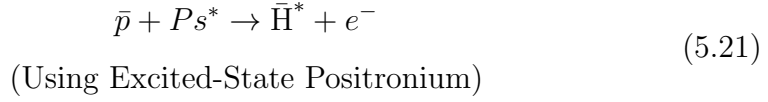
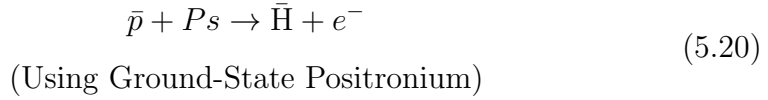
$$P \sim n_e R^3 \quad (5.18)$$

Multiplying this probability in Equation 5.18 with the impact parameter in 5.13 then gives us the new scaling for TBR:

$$\Gamma_{\text{TBR}} \sim (\pi b^2 n_e v_e)(n_e R^3) \sim n_e^2 b^2 R^3 v_e. \quad (5.19)$$

Since R does not scale with temperature, $b \propto T^{-1}$ and $v_e \propto T^{\frac{1}{2}}$, the rate will scale like $\Gamma_{\text{TBR}} \sim n_{e^+}^2 T_{e^+}^{-\frac{3}{2}}$. In this scenario, the temperature dependence is not as profound as in the steady state situation.

Instead of using a cloud of cold positrons, one may also use ground state or excited state positronium to make antihydrogen:



Antihydrogen production using positronium relies on the antiproton capturing the positron from the positronium atom, with the electron taking away the excess energy. In this case, the rate of formation of antihydrogen using ground-state positronium would be different compared to using excited-state positronium. This is because of two primary reasons:

- **Positronium Lifetime:** In ground state positronium, the spins of the electron and positron can either be antiparallel (1S_0 singlet state) or parallel (3S_1 triplet state). The singlet and triplet states have a mean lifetime of 125 ps and 142 ns respectively [73]. Excited-state positronium is metastable however, requiring decay to the ground state before electron-positron annihilation takes place. This reduces the rate of depletion of the positronium available for antihydrogen formation, resulting in an overall increased antihydrogen formation rate.
- **Positronium Binding Energy:** When an antiproton displaces an electron from a positronium atom, the resultant antihydrogen typically inherits the binding energy of the positronium. Since it is easier to dislodge loosely bound electrons than tightly bound ones, it follows that antiprotons interacting with excited-state positronium have a much larger antihydrogen-formation cross-section than they would have with

ground-state positronium.

To date, only three-body recombination [36, 37, 65, 68, 69], and antihydrogen synthesis using excited-state positronium [74] have been demonstrated experimentally. When antihydrogen has been made by mixing positron and antiproton plasmas, the high formation rates have been inconsistent with radiative recombination [66, 75, 76]. Attempts to stimulate radiative recombination using lasers have also been unsuccessful [77]. Antihydrogen production with ground-state positronium has not been tried since significant enhancement of antihydrogen production is expected when using excited-state positronium due to the reasons listed earlier.

While [74] has been the only successful attempt at creating antihydrogen atoms using Ps^* , there is considerable promise in using this method since it avoids the heating of positron plasmas due to space-charge or thermalization with a warmer trap environment. By mixing cold antiprotons with excited-state positronium, much colder antihydrogen atoms may be created.

5.2 Dependence on Antiproton Temperature

In both radiative recombination and three-body recombination, the relative-velocity between the positron and antiproton plays a critical role in determining the interaction cross-section, and consequently the rate of antihydrogen formation. Assuming thermal equilibrium between the two species, the relative-velocity is dominated by the positron velocity, since

$$v_{\bar{p}} = \sqrt{\frac{m_e}{m_p}} v_{e^+} \approx 0.02 v_{e^+} \quad (5.22)$$

It is important therefore to ascertain whether antiprotons injected into the positron plasma achieve thermal equilibrium or not prior to antihydrogen formation.

The first experiment to produce cold antihydrogen at ATHENA was carried out using the Simple Mixing technique where antiprotons were shot into a cold (15 K) positron plasma with ~ 15 eV of axial energy [65]. The experiment was then repeated with heated positrons at 175 K, 500 K, and 3500 K [66]. It was observed that the spatial distribution of antihydrogen annihilations was independent of the positron temperature, indicating that the antihydrogen formed had velocity distributions not determined by the positron velocity. Subsequently, a simple analytical model was used to set a lower bound on the axial temperature of the antiprotons at 150 K, concluding that the antiprotons had not thermalized with the positrons prior to antihydrogen formation.

With the mixing parameters used in the study, antihydrogen formation started before antiprotons equilibrated with the positrons, resulting in anti-atoms with temperatures in the hundreds of kelvin. These temperatures are not favorable for trapping, considering that the trap depths are on the order of ~ 1 K. This necessitates using colder antiprotons. In fact, using Eq. 5.22, one can estimate that antiprotons on the order of ~ 2 eV have velocities comparable to ~ 15 K positrons. Therefore, using colder antiprotons and a better mixing technique (e.g. Autoresonant Mixing), one could expect antiprotons to thermalize before antihydrogen atoms begin to form, and result in more trappable anti-atoms. In this situation, both antihydrogen formation and trapping rates would be greatly influenced by the temperature of the positrons, making it imperative for us to understand how these rates scale with the positron parameters.

5.3 Dependence on Positron Temperature

As we have seen earlier in this chapter, the rate of antihydrogen production for both radiative recombination (Eq. 5.11) and three-body recombination (Eq. 5.16) is influenced by the temperature of the positron plasma used for mixing. In particular, colder positrons are expected to result in a higher antihydrogen formation rate, albeit the impact is much higher in the case

of three-body recombination. In an antihydrogen experiment with a given mixing regime, one could study the positron temperature dependence on antihydrogen production and determine whether formation is dominated by RR, TBR, or a mixture of both. Such an understanding would greatly assist us in optimizing the desired positron plasma parameters to create trappable antihydrogen.

The ATHENA experiment which produced the first cold antihydrogen atoms [65] was also the first experiment where one could try to determine whether formation was mainly due to RR or TBR. Antihydrogen was produced by mixing $N_{\bar{p}} = 10^4$ antiprotons with $N_{e^+} = 7 \times 10^7$ positrons, with a positron plasma density of $1.7 \times 10^{14} \text{ m}^{-3}$. Assuming positrons at 15 K, this would give $N_{\bar{p}}\Gamma_{\text{RR}} \approx 30 \text{ Hz}$ and $N_{\bar{p}}\Gamma_{\text{TBR}} \approx 1.4 \text{ MHz}$. The experimentally measured rate at the onset of mixing was $432 \pm 44 \text{ Hz}$ [75], an order of magnitude higher than what was predicted by RR alone, and several orders of magnitude below the calculated rate for TBR.

Taking into account the data points at 175 K and 500 K in this study, it is clear that colder positrons do in fact lead to higher antihydrogen formation rates as expected. Additionally, the high formation rates at 175 K and 500 K are strong indicators for TBR as the dominant process. However, the data is inconsistent with any simple scaling law in the form of aT^{-b} . It was conjectured that the short transit time of the antiprotons in the positron plasma is responsible for this behavior, and simulations [71] show rough agreement with this suppressed three-body recombination process.

Before their dissolution in 2005, the ATHENA collaboration made another attempt at extracting the positron temperature scaling, this time by observing how the antihydrogen production rate evolved as heated positrons cooled down during mixing [76]. Initially, antihydrogen formation was suppressed using RF-heating of the positrons. Once the heating was turned off, the positrons cooled by emitting cyclotron radiation in the 3 T magnetic field. The increase in antihydrogen production during these few seconds of cooling was measured for various levels of initial heating. From the resulting data,

a temperature scaling of $T^{-1.1 \pm 0.5}$ was extracted for temperatures between 1.1×10^3 K to 1.7×10^4 K. As before, the high formation rates were indicative of TBR, but the suppressed scaling again pointed towards a non-steady-state formation mechanism between the positrons and antiprotons.

A similar study of the positron temperature scaling was carried out by the ALPHA collaboration [78]. A pulsed radio-frequency source was used to heat positron plasmas which were then mixed with antiprotons for 1 second. By adjusting the amplitude of the RF heating, positron temperatures between 40 K and 1000 K could be achieved. The rate of antihydrogen formation for several choices of positron temperatures was measured. The result was a temperature scaling of $T^{-0.5 \pm 0.1}$, consistent with the observations of the ATHENA experiment.

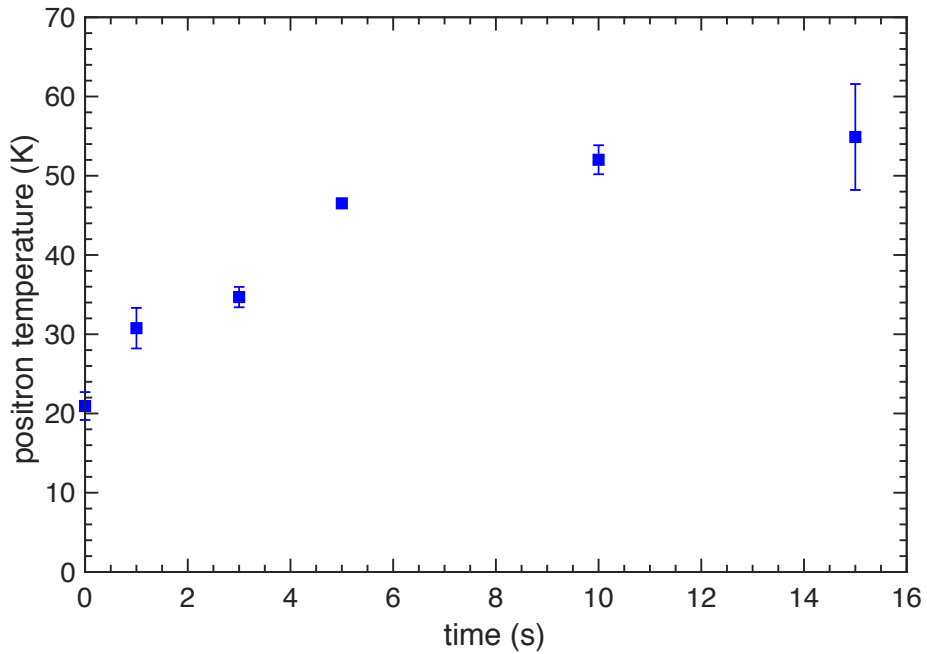
With the development of evaporative cooling for charged particles [79], antiproton plasmas could be cooled from hundreds of kelvin to tens of kelvin, making them more amenable for the production of cold antihydrogen. Alternatively, by using autoresonant mixing (see Section 4.2.2), antiprotons could be injected into a positron plasma without significantly increasing their temperature [67]. The latter technique was used to first create trapped antihydrogen in ALPHA [44], and a combination of both techniques later helped in the production and confinement of antihydrogen for up to a thousand seconds [68]. In both of these instances, and subsequent antihydrogen production via both autoresonant mixing [36] and slow merge [37], the formation rates are consistent with antiprotons having thermalized with the positron plasma prior to antihydrogen formation [80].

5.4 Recent Positron Temperature Studies at ALPHA

Positron plasmas in the Mixing Trap at ALPHA have a peculiar property that they do not equilibrate with the ambient trap temperature: while the

electrodes reach temperatures as low as 6 K, the positron plasma rarely goes below 15K. Furthermore, once the plasma is evaporatively cooled, it subsequently quickly heats up to several tens of kelvin in a matter of a few seconds (an example is shown in Figure 5.1). Heating is similarly observed with electrons, indicative of a charge-neutral source of lepton heating in the trap environment. This is one of the main reasons why the antiproton cloud is mixed with the positron plasma immediately after the latter is evaporatively-cooled.

Figure 5.1: A typical plasma heating curve in the ALPHA Mixing Trap. The plasma consists of 1.5×10^6 positrons at a density of $\sim 6 \times 10^7 \text{ cm}^{-3}$ and a radius of 0.66 mm. The error bars are one standard deviation for three measurements at each point.



In 2016, ALPHA carried out hundreds of experiments in which antihydrogen atoms were produced with different plasma parameters and changing trap environments. In between experiments, positron temperatures and heating rates were periodically measured. The temperature of the remaining positrons after mixing was also measured. Hot positrons and/or high heating rates were often observed to result in diminished production and low

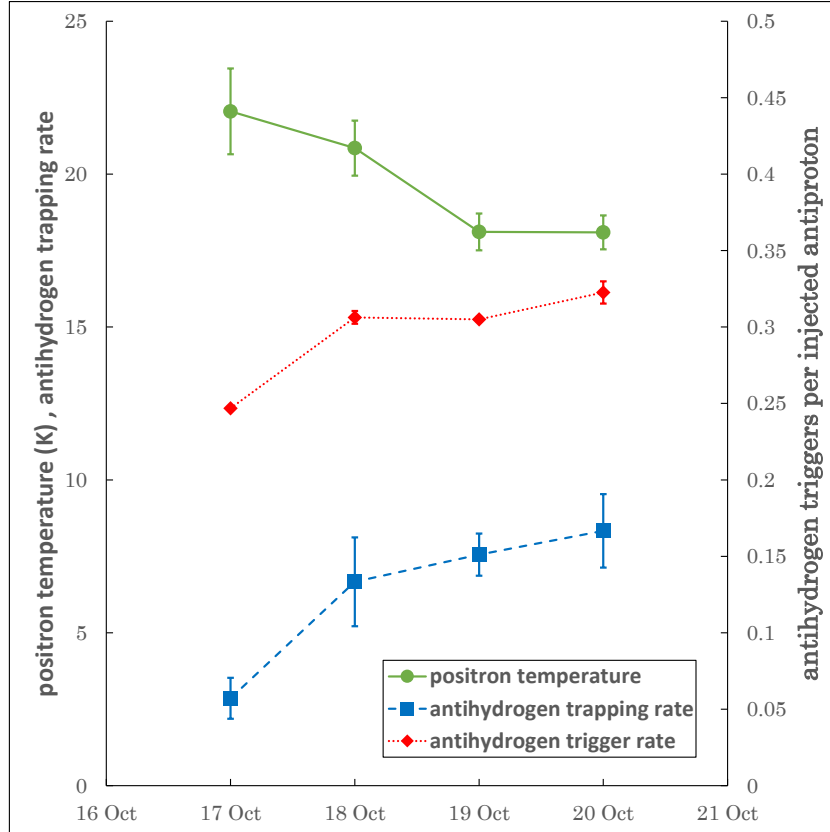
trapping rates.

One such example is shown in Figure 5.2, where antihydrogen mixing triggers, trapping rate, and positron temperature are plotted for several days after a “cool down” of the trap (in a cool down, the internal components of the apparatus are allowed to thermalize with a bath of liquid helium at 4.2 K). The cool down took place on the 16th of October, but not all of the internal components had thermalized immediately. Antihydrogen is made by mixing a plasma of 1.6×10^6 positrons with a radius of 0.66 mm with a cloud of $\sim 80,000$ to $100,000$ antiprotons and a radius of ~ 0.9 mm. During the four-day period following the cool down, a steady increase in the antihydrogen production and trapping rate is observed, in parallel with a steady decline in the positron temperature. While this presents a qualitative relationship between antihydrogen trapping rate and positron plasma temperature, the measurement uncertainty is too large to obtain any meaningful measure of how trapping quantitatively scales with temperature.

A somewhat different study offers a better estimate of this relationship. Recently, it has been observed at ALPHA that exposing the Mixing Trap to ultraviolet radiation increases the nominal lepton heating rate. The purple curve in Figure 5.3a shows how the positron heating rate varies with time as the trap is exposed to UV radiation. The heating rate is calculated by looking at the positron temperatures 0 s, 1 s, and 3 s after EVC, and then obtaining the gradient of the linear fit as the nominal heating rate. Initially, before the laser is turned on, the nominal heating rate is approximately 5 Ks^{-1} . With the introduction of the laser into the trap, the heating rate increases monotonically. After about 200 minutes of exposure, the heating rate has approximately increased by a factor of 4. After another 200 minutes, the heating rate is an order of magnitude higher.

During this study with UV radiation, antihydrogen production was carried out after each 500 seconds of laser exposure. Figure 5.3a also shows the evolution of the antihydrogen trigger rate with time. The relationship between the trigger rate and laser exposure appears to be very linear, suggesting a

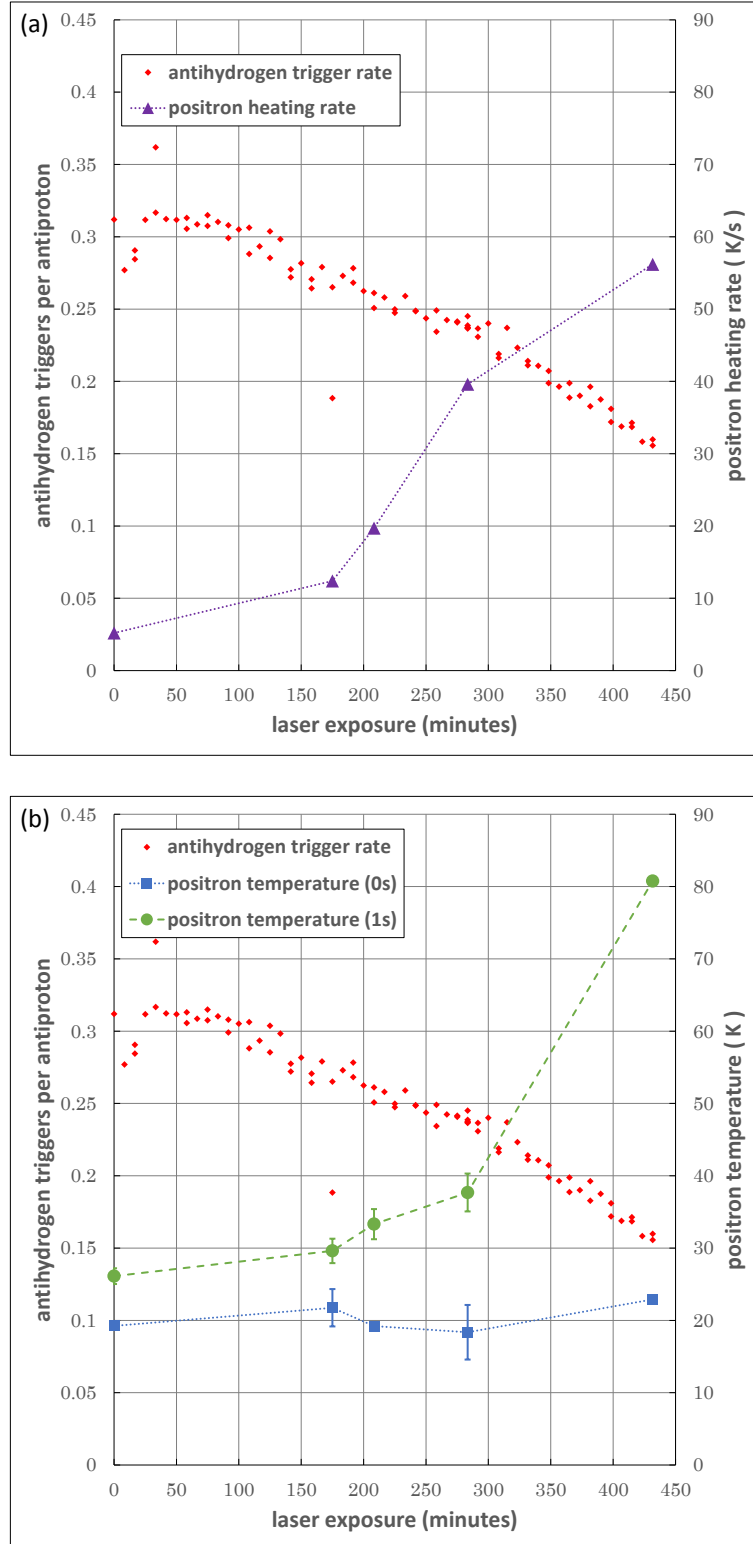
Figure 5.2: Plot showing the evolution of positron temperature (green), mixing triggers (red) and trapping rate (blue) for a 4-day period after a cool down. As the apparatus cools down, the positron temperatures decrease while the mixing triggers and trapping rate increase in tandem. The error bars are standard errors of the mean. The error bars on the trapping rate are from Poisson counting statistics. Data is from October 2016 (Series 560).



uniform mechanism that causes this. At present, this mechanism is not fully understood and is an active area of research within ALPHA.

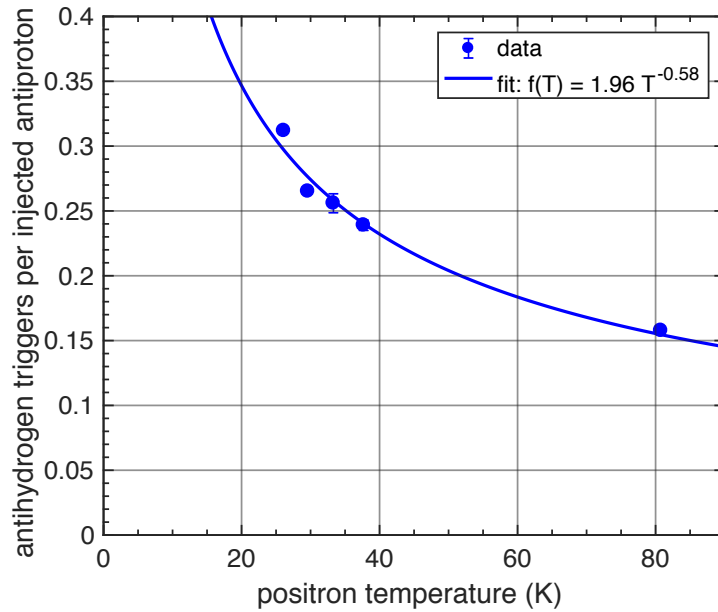
In Figure 5.3b, the positron temperatures immediately after EVC (blue) and 1 s after EVC (green) are plotted. The positron temperature right after evaporative cooling does not seem to be affected by laser exposure, with the temperature remaining constant around 20 K. It is presently unclear whether the technique we use to measure the positron temperature is reliable for measurements immediately after EVC, since it takes some fractions of a second

Figure 5.3: (a) Variation of positron heating rate (purple) and antihydrogen production rate (red) with exposure to ultraviolet radiation. (b) Positron temperature 0 s after EVC (blue) and 1 s after EVC (green) for various laser exposure times. The error bars, where given, are standard errors of the mean. Data is from October 2016 (Series 566-568).



for the positron space-charge to equilibrate in the post-EVC potential well. The stability of the temperature immediately after EVC could therefore be an artifact of the measurement technique. Alternatively, assuming that the technique is indeed reliable for the given plasma parameters, this suggests that the reduction in the rate of antihydrogen formation stems not from a change in initial positron temperature but is due to heating that takes place during mixing. This is reinforced by the positron temperature 1s after EVC, which shows a marked increase as a result of exposure to ultraviolet radiation.

Figure 5.4: A plot of the antihydrogen trigger rate variation with positron plasma temperature. A fitted curve of the form aT^{-b} is also shown. The exponent $b = 0.58 \pm 0.09$, is a measure of the temperature scaling.



Using the positron temperature data in Figure 5.3b and the associated antihydrogen trigger rates, a relationship between trigger rates and positron temperature is plotted in Figure 5.4. A curve of the form aT^{-b} is fitted to the data, resulting in a temperature scaling of $T^{-0.58 \pm 0.09}$, similar to the results of both the ATHENA and ALPHA measurements reported earlier in Section 5.3.

In addition to understanding how positrons impact antihydrogen formation,

it is perhaps even more important to determine how the overall antihydrogen trapping rate scales with temperature. In this study, the variation in antihydrogen trapping was within the noise level of the measurement, and therefore a scaling could not be determined.

In late October 2016, adiabatic cooling of positrons was added to the antihydrogen production sequences. This allowed the positron temperatures at mixing to be as low 10 K compared to the 20K average shown in Figure 5.3a. These low temperatures resulted in increased antihydrogen trapping rates, with some production sequences averaging upwards of 10 trapped anti-atoms per attempt. With this improved rate of trapping, the effect of positron temperature was reinvestigated.

Figure 5.5 shows the results of a similar study which was carried out at higher trapping rates. As before, exposure to UV radiation results in an increase in positron heating rate and a decrease in antihydrogen formation. Similarly, the temperature right after EVC remains unaffected by the exposure. This time however, the minimum average positron temperature is approximately 15 K, 5 K lower compared to the earlier study when adiabatic cooling of positrons was not implemented.

Figure 5.6 shows the antihydrogen trapping data for the same study. Despite the large error bars on the trapping rates, the overall trend of diminished trapping with laser exposure is evident. The data in Figure 5.6b is used to extract a rough scaling of antihydrogen trapping with positron temperature. A plot for this is shown in Figure 5.7, which gives a temperature scaling of $T^{-1.6 \pm 0.9}$.

This is the first reported attempt of quantifying the scaling of antihydrogen trapping rate with the temperature of the positron plasma. The analysis presented here was only made possible as a result of the significantly high trapping rates recently achieved at ALPHA. If this scaling holds true at lower temperatures, a factor of 4 reduction in positron plasma temperature

Figure 5.5: (a) Variation of positron heating rate (purple) and antihydrogen production rate (red) against the duration of laser exposure. (b) Positron temperature 0 s after EVC (blue) and 1 s after EVC (green) at different levels of UV exposure. The positron temperatures immediately after EVC are ~ 15 K. Data is from November 2016 (Series 575).

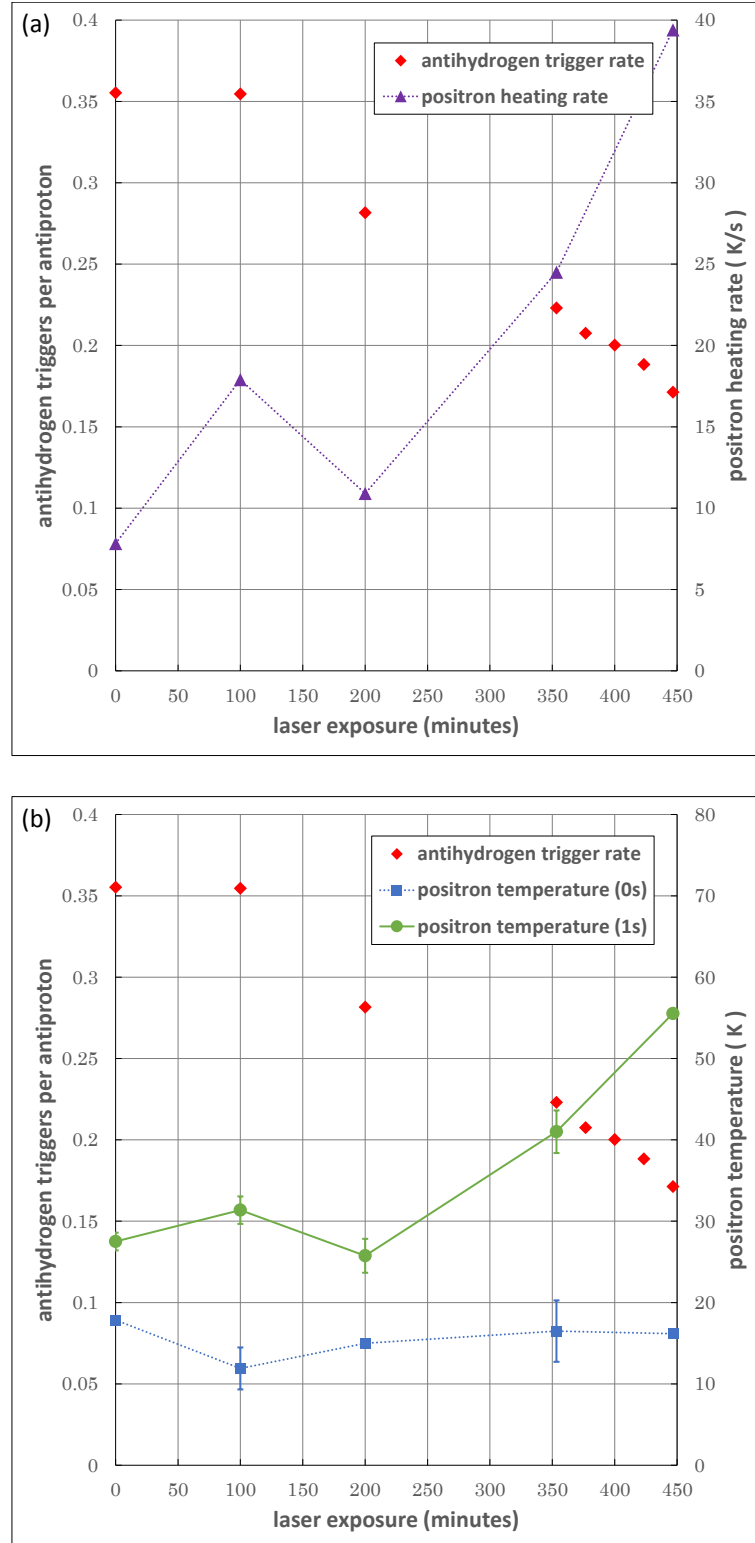


Figure 5.6: (a) Variation of positron heating rate (purple) and antihydrogen trapping rate (brown) with exposure to ultraviolet radiation. The trapping rate also decreases roughly linearly with laser exposure. The error bars on the trapping rate are due to Poisson counting statistics. (b) Positron temperature 0 s after EVC (blue) and 1 s after EVC (green) at various exposure times.

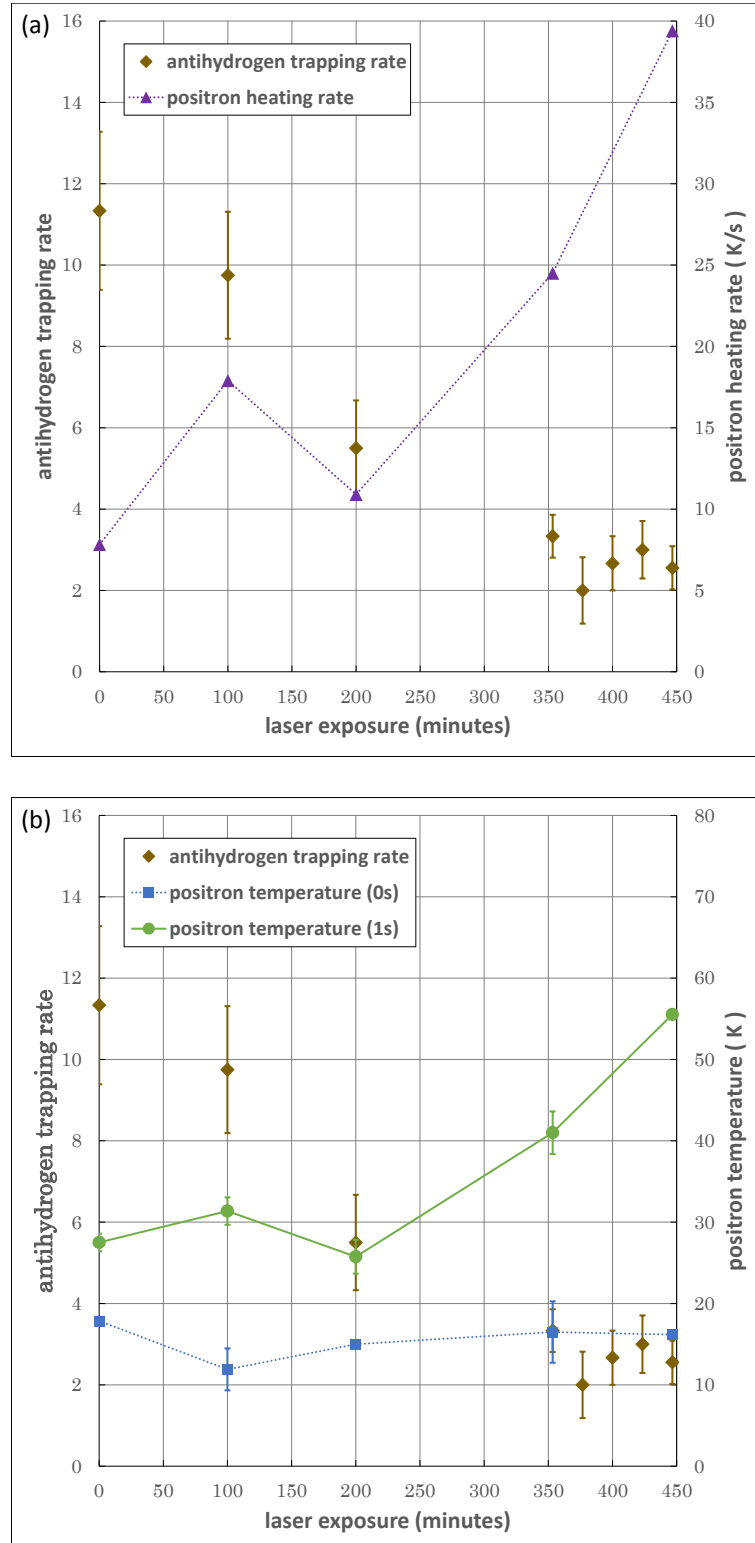
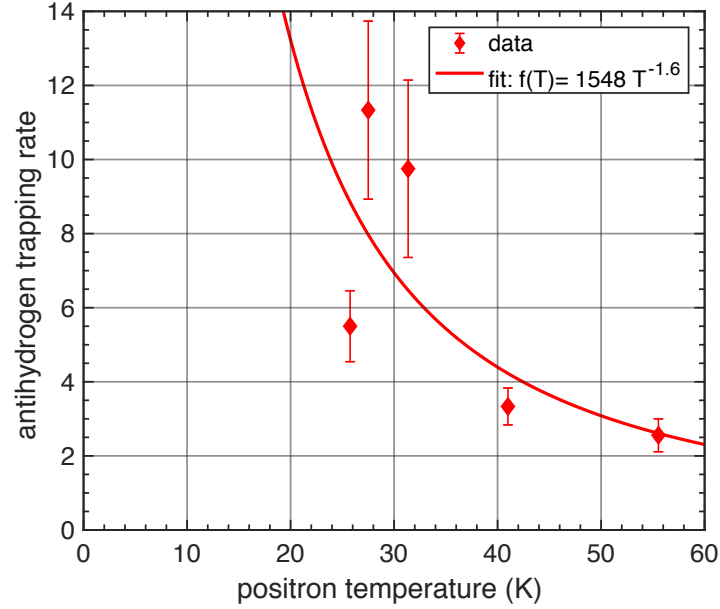


Figure 5.7: A plot of antihydrogen trapping rate variation vs positron plasma temperature. A fitted curve is also shown. The exponent of the fit, -1.6 ± 0.9 , provides a rough temperature scaling of antihydrogen trapping with positron temperature.



could result in a 5-fold to 25-fold increase in trapping rate! Such a dramatic improvement in trapping would allow measurements on antihydrogen to be carried out with much more precision and in significantly less time. This provides the main motivation for the rest of this thesis, where I discuss the implementation of a technique that promises to provide us with significantly colder positrons for antihydrogen production at ALPHA.

Chapter 6

Sympathetic Cooling of Positrons

The process in which particles of one type are used to cool particles of another type is known as sympathetic cooling. In an electromagnetic trap, sympathetic cooling of one charged species (recipient species) can be accomplished by introducing another charged species (intermediary species) that is actively being cooled. As the two types of particles interact via Coulomb collisions, the excess kinetic energy of the recipient particles is transferred to the intermediary particles, which are in turn cooled by the cooling mechanism. The cooling power is effectively transmitted to the recipient as it thermalizes with the intermediary.

The electron-cooling of antiprotons described in Section 2.4.4 is an example of sympathetic cooling in a Penning-Malmberg trap. Here, the electrons behave as the intermediary, cooling the antiprotons via elastic Coulomb collisions while self-cooling through the emission of cyclotron radiation.

6.1 Past Experiments

Sympathetic cooling of ions using a different ion species is a routinely-used cooling technique in Penning traps [81–84]. By laser-cooling the intermedi-

ary ions to sub-kelvin temperatures, the recipient ions can also be cooled to extremely low temperatures. Since laser-cooling requires the ions to have a suitable set of energy levels, only a few types of ions meet this criterion. Among positive ions, Be^+ , Mg^+ , and Ca^+ have successfully been used for sympathetically cooling atomic and molecular ions. To date, no negative ions have been laser-cooled [85]. Efforts are underway however to develop schemes to laser-cool atomic [85, 86] and molecular [87, 88] anions with a view to utilize them for the sympathetic-cooling of other negatively-charged particles.

The first demonstration of sympathetically laser-cooled ions was implemented in 1986 at NIST in Colorado, USA [81]. The experiment used laser-cooled Be^+ ions at $T \approx 10$ mK to cool Hg^+ ions to temperatures < 1 K. Shortly thereafter, the possibility of producing ultracold positrons via sympathetic laser-cooling was proposed [89]. It was foreseen that that by using laser-cooled beryllium ions in a Penning trap, positron plasmas with densities $n > 10^{10} \text{ cm}^{-3}$ and temperatures $T < 10$ mK could potentially be achieved.

At the turn of the century, detailed simulations of sympathetically laser-cooled positrons in a cylindrical Penning trap further bolstered the viability of producing cold, dense positrons with laser-cooled beryllium ions [90]. In experimental investigations that followed [91, 92], Jelenkovic *et al.* reported the first demonstration of sympathetically laser-cooled positrons. The experiment used $\sim 10^5$ Be^+ ions that were laser-cooled to sub-kelvin temperatures (although the beryllium ion temperature was not measured *in situ*, the technique used had previously resulted in recorded ion temperatures of < 100 mK). The resulting sympathetic cooling of positrons resulted in a cold positron plasma on-axis, centrifugally separated from the beryllium ions that had collected at the outer edges of the plasma. The positron plasma had a density of $n \approx 4 \times 10^9 \text{ cm}^{-3}$, and temperatures $T_{\parallel} \approx 5$ K along the axis and $T_{\perp} \approx 10$ K perpendicular to the axis.

The observed centrifugal separation indicated that the positrons had thermalized with the beryllium ions. This would imply that both positrons and

beryllium ions had similar densities and rotational frequencies, consistent with simulation predictions made in Reference [90]. With regards to temperature however, the naive expectation of positron temperatures on the order of ~ 100 mK were not met. This may be because Coulomb collisions in beryllium-positron mixed plasmas are three orders of magnitude weaker than in ion-ion plasmas (due to the low positron mass), resulting in a much reduced average cooling rate. Additionally, the strong 6 T axial magnetic field alters the positron trajectories to closely follow magnetic field lines, a phenomenon known as plasma magnetization. In a magnetized plasma, the cyclotron frequency is much larger than the average rate of collisions between particles. Magnetization of the cold positron plasma in the 6 T field can therefore greatly reduce collisional relaxation rate of the parallel and perpendicular temperatures [53]. Finally, Coulomb interactions are significantly suppressed once centrifugal separation takes place, resulting in the positrons equilibrating at a higher temperature when the external heating rate is at par with the suppressed cooling rate.

6.2 Simulations in the ALPHA Apparatus

To date, the NIST experiment [91, 92] is the only instance where sympathetic laser-cooling of positrons has been successfully performed. Given the significant benefits of having cold positrons for antihydrogen formation and trapping (see Chapter 5), there is tremendous interest in the ALPHA experiment for developing sympathetically laser-cooled positrons. Compared to Jelenkovic *et al.*, there are significant differences in the experimental conditions and trapping apparatus at ALPHA. In particular, the manner in which antihydrogen is created at ALPHA presents two fundamental questions:

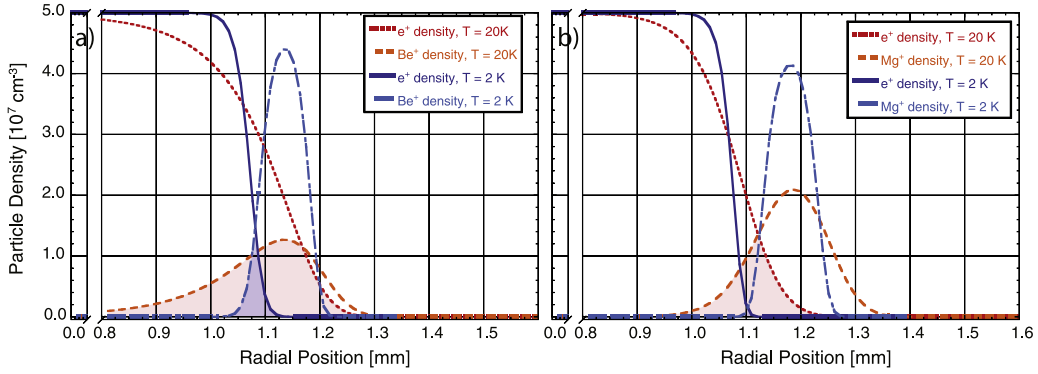
1. Is it possible to sympathetically laser-cool ~ 2 million positrons in the ALPHA Penning-Malmberg trap?
2. How does the presence of laser-cooled ions impact the antihydrogen trapping rate and lifetime?

To answer these questions, Madsen *et al.* [93] performed a study to ascertain the feasibility of sympathetically cooling positrons in the ALPHA trap environment. In particular, the viability of both Be^+ and Mg^+ was investigated. Both ions have a convenient $3s^2S_{1/2} \rightarrow 3p^2P_{3/2}$ cooling transition at 313 nm and 279 nm and Doppler-cooling limits of 0.47 mK and 1.0 mK for Be^+ and Mg^+ respectively. And with the availability of commercial lasers with powers in excess of 100 mW at the two specified wavelengths [94], cooling the ions to less than ~ 100 mK is not an impediment.

The study used a numerical code [95] previously developed to calculate the self-consistent density and temperature distributions of a mixed plasma by relaxation from a given set of initial distributions. As predicted in Reference [96] and observed by Jelenkovic *et al.* in [91, 92], a mixed ion-positron plasma in a Penning trap reaches equilibrium with a partial or total centrifugal separation of the two species. The degree of separation dictates the rate of inter-species collisions, and consequently determines the overall cooling ability of the laser-cooled ion species. Figure 6.1 shows the calculated radial density profile of a mixed ion-positron plasma at equilibrium, at two different temperatures. For Be^+ ions, there is significant ion-positron radial overlap at 20 K, but a much reduced overlap at 2 K, indicating a greater degree of separation at colder temperatures. Similarly, for Mg^+ ions, centrifugal separation is already present at 20 K, and the two species have completely separated at 2 K. Given the substantially larger separation of Mg^+ relative to Be^+ , one would expect a lower sympathetic-cooling power of the former due to a reduced rate of collisions. In this situation, using beryllium ions would be the better choice to sympathetically laser-cool positrons.

Since the extent of radial overlap of the two species influences the overall sympathetic cooling rate, and lower temperatures result in lower radial overlaps, a quantity of interest is the equilibrium positron temperature reached when colder and colder ions are used. Figure 6.2 shows how the equilibrium positron temperature varies with the ion temperature used for cooling. At higher temperatures, the positron temperature closely follows the ion temperature, with no significant difference between using only 50,000 ions

Figure 6.1: (a) Radial density profile for e^+ and Be^+ at 20 K (red) and 2 K (blue) at equilibrium. (b) Radial density profile for e^+ and Mg^+ at 20 K (red) and 2 K (blue) at equilibrium. The shaded area indicates the region where both species have non-zero density. The calculated profiles are for 2×10^6 positrons and 10^5 ions. From [93]

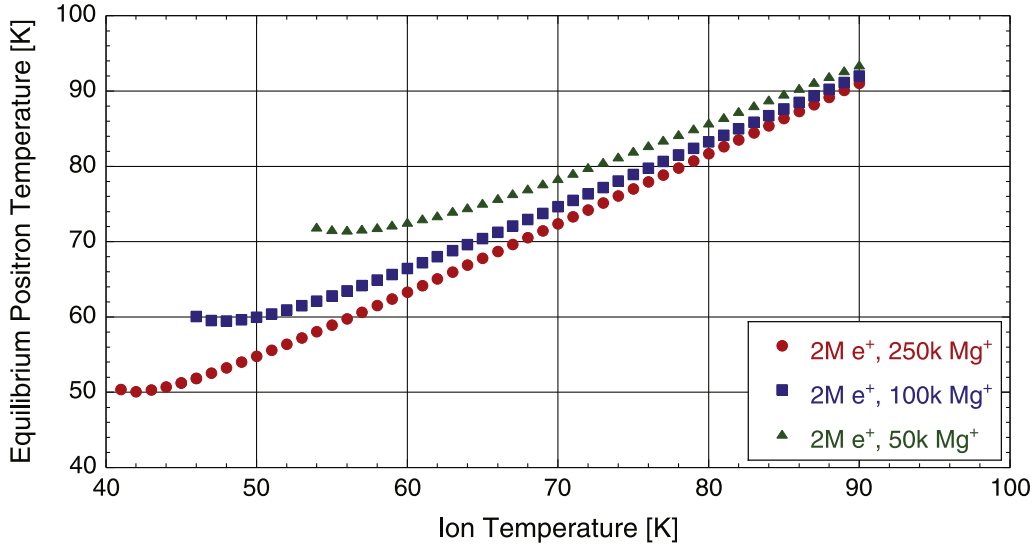


or 250,000 ions. This follows from the understanding that at higher temperatures, collision rates are high enough that cooling power is effectively transmitted to the positrons regardless of ion quantities.

At low ion temperatures however, ion quantities begin to matter. For any given ion temperature between 60 K and 80 K, using more cooling ions results in a lower equilibrium positron temperature. This is because higher ion quantities result in higher ion-positron collision rates, and therefore better cooling. As even colder ions are used, the equilibrium positron temperature reaches a minimum. This is because as ions get colder, centrifugal separation results in a diminished ion cooling power until a minimum positron equilibrium temperature is reached. If even colder ions are used, the cooling power is too low to overcome the external positron heating rate, and the positron equilibrium temperatures would start to rise beyond this minimum. It is worth noting that the minimum positron equilibrium temperature is lower when higher ion quantities are used. Using larger ion numbers may therefore compensate for the reduced cooling due to centrifugal separation.

Figure 6.3 shows how the minimum positron equilibrium temperature varies as a function of ion quantity for both magnesium and beryllium ions. As

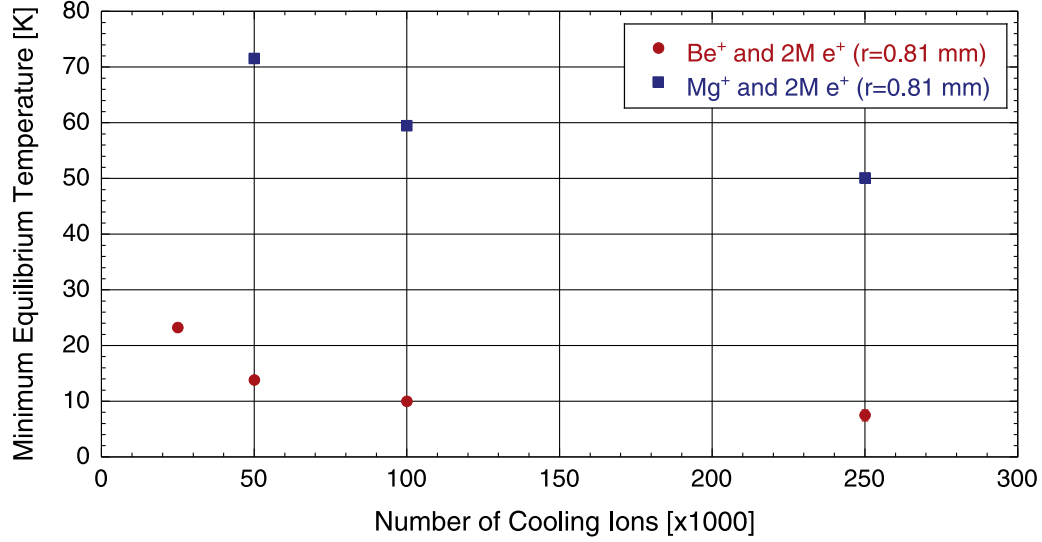
Figure 6.2: Equilibrium temperature of 2×10^6 positrons as a function of the Mg^+ ion temperature for three different ion quantities. The calculation assumes an ambient temperature of 7 K, an external positron heating rate of 26 K s^{-1} , and a positron plasma radius of 0.81 mm. From [93]



observed previously, using higher ion quantities results in a lower minimum temperature. However, for higher ion quantities, the trend begins to saturate, and one might not expect any significant benefit of using more than 10^6 ions for the conditions used in this simulation. This is consistent with the observation that, beyond the saturation point, adding more ions only results in their accumulation on the outside layers of the centrifugally separated ion plasma which offers no improvement to the ion-positron collision rate, and therefore no additional cooling power.

Figure 6.3 does however suggest that a factor of 5 reduction in minimum positron temperature can be achieved simply by using Be^+ ions instead of Mg^+ ions in the sympathetic laser-cooling process. This is explained by the much lower extent of centrifugal separation of Be^+ compared with Mg^+ described earlier in Figure 6.1. Much lower positron temperatures can therefore be obtained when using beryllium, making it the species of choice for the sympathetic cooling of positrons.

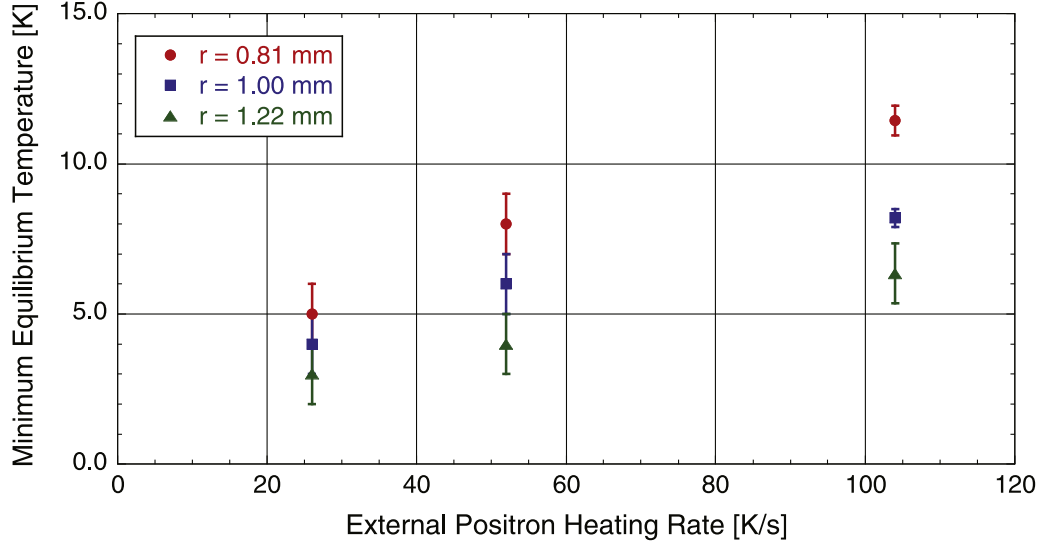
Figure 6.3: Minimum equilibrium temperature of 2×10^6 positrons for several different ion quantities. The results are plotted for both Be^+ (red) and Mg^+ (blue). From [93]



In Section 5.3, we saw how the ambient lepton heating rate is affected by UV radiation in the Mixing Trap. We are therefore interested in understanding how the minimum temperature of sympathetically laser-cooled positrons is affected by the external positron heating rates. Figure 6.4 shows the minimum positron equilibrium temperature of 2×10^6 positrons for different levels of external heating. When heating rates are low (26 K s^{-1}), positron temperatures $< 5 \text{ K}$ are attainable. For a high positron heating rate however (104 K s^{-1}), the minimum equilibrium positron temperature is higher and can be anywhere from 6 K to 12 K , depending on the radial size of the positron plasma, with larger plasmas resulting in lower temperatures. This is because larger positron plasmas share a larger surface area of overlap with the cooling ions, and therefore provide a higher collision rate and consequent cooling power.

The question about the effect of beryllium ions on the antiproton and antihydrogen lifetime is an important one since premature annihilation of antiprotons and unwanted loss of antihydrogen atoms need to be avoided. In the case of a mixed ion-positron plasma in equilibrium at 5 K with 10^5 beryl-

Figure 6.4: Minimum temperature of 2×10^6 positrons sympathetically cooled with 250,000 Be^+ ions as a function of external positron heating rate. From [93]



lithium ions and 2×10^6 positrons, the ion density is on the order of 10^2 cm^{-3} along the trap axis (where the positrons are located) and $\sim 10^7 \text{ cm}^{-3}$ at the band where the ions have centrifugally separated [93]. These densities correspond to an antiproton-ion collision rate of $\sim 1.5 \times 10^{-4} \text{ s}^{-1}$ and 15 s^{-1} respectively. In the case of autoresonant injection of antiprotons into the positron cloud, almost all of the antiprotons are injected along the trap axis where the collision rate is negligible. In the case of smerge, where the positron and antiproton clouds are of comparable radial size, a small fraction of antiprotons may be exposed to ions in the high density region and annihilate. Simulations indicate that for a positron plasma at 10 K and density $> 1 \times 10^{14} \text{ m}^{-3}$, antiprotons injected at higher radii produce proportionally fewer trappable antihydrogen atoms [97]. This effect is expected to be more pronounced at lower temperatures. The loss of these higher-radii antiprotons would therefore not be a serious detriment to the overall antihydrogen trapping rate.

The antihydrogen-ion collision rate calculated by Madsen *et al.* in the high ion-density region is on the order of 0.1 s^{-1} [93]. Antihydrogen atoms once created spend most of their life near the trap walls where ions are absent.

Furthermore, since antihydrogen formation only takes place for 1.5 s after which the ions are cleared out, this mechanism of antihydrogen loss can effectively be neglected.

In summary, laser-cooled beryllium ions present a compelling case to sympathetically laser-cool positrons in the ALPHA Mixing Trap. Since typical mixing cycles employ 2×10^6 positrons for antihydrogen formation, a few 10^5 beryllium ions would be sufficient to transmit cooling to the positron cloud and achieve temperatures of ~ 5 K. Assuming the positron temperature scaling of antihydrogen trapping (Figure 5.7) holds true, this four-fold drop in positron plasma temperature should result in a factor of ~ 12 improvement in trapping rate. This provides us with the motivation to develop a low-energy beryllium ion source that is compatible with the ALPHA apparatus and can reliably deliver $\sim 10^5$ laser-cooled ions.

Chapter 7

Simulations of 1-D Beryllium Laser Cooling in ALPHA

In order to sympathetically laser-cool positrons in ALPHA, the beryllium ions would need to be loaded into the center of the Mixing Trap where one of the laser pathways crosses the trap axis (see Figure 2.7). Using a single 313 nm continuous-wave laser, the ions would then be cooled in a process known as Doppler cooling.

The method of Doppler laser cooling utilizes the momentum of photons to slow down and thereby cool the ions. An ion illuminated by laser light of wavelength λ can absorb a photon to reach an excited state, increasing its momentum by $\frac{h}{\lambda}$ in the incident photon direction. The excited ion then re-emits the photon in a random direction, changing its momentum by $-\frac{h}{\lambda}$ in that direction. Since the photon absorptions are directional but the re-emissions are approximately isotropic, the net effect after N scatterings is a change in the ion's momentum by $\frac{Nh}{\lambda}$ in the direction of the incident photons. Depending on the electronic structure and energy levels of the ion species, this scattering mechanism is resonant at certain wavelengths of light. For a laser red-detuned below the resonance frequency ω_0 , ions moving towards the laser observe Doppler-shifted photons closer to resonance, scattering them at a much higher rate than ions moving away from the laser. For a closed two-level system described above with ions cycling between a ground state and

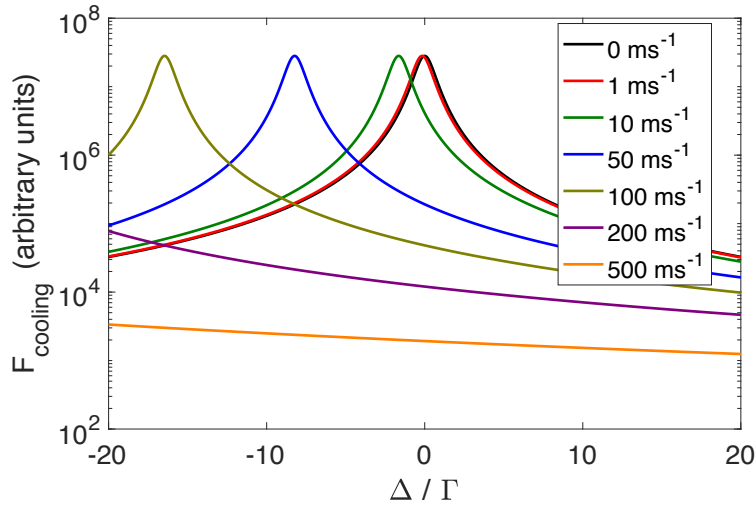
an excited state, the photon scattering rate R_{scatt} is given by

$$R_{\text{scatt}} = \frac{\Gamma}{2} \frac{I/I_{\text{sat}}}{1 + I/I_{\text{sat}} + 4\delta^2/\Gamma^2} \quad (7.1)$$

where Γ is the natural linewidth of the excited state, I is the laser intensity, I_{sat} is the saturation intensity and $\delta(B) = \omega - \omega_0(B) + kv$ is the frequency detuning from resonance, taking into account the Doppler shift kv . The laser-cooling force is then simply given by

$$F_{\text{cooling}} = \text{photon momentum} \times \text{scattering rate} = \hbar k \times R_{\text{scatt}} \quad (7.2)$$

Figure 7.1: Laser-cooling force vs. laser detuning for ions at various velocities and $I/I_{\text{sat}} = 0.86$. For positive velocities indicated, the ions are moving towards the laser.



For a fixed resonance frequency ω_0 , we define the intentional detuning of the laser Δ away from resonance as

$$\Delta = \omega - \omega_0 \quad (7.3)$$

where ω is the laser frequency.

Figure 7.1 shows the laser-cooling force as a function of laser detuning for ions at different velocities. For stationary ions (0 ms^{-1}), since there is no

Doppler shift, the cooling force is highest for laser frequency on resonance with the transition frequency (i.e. $\Delta = 0$). For slightly energetic ions ($1 \text{ ms}^{-1} - 50 \text{ ms}^{-1}$), the cooling peak is shifted to the left, indicating that cooling force is highest when the laser frequency is slightly red-detuned. For ions with sufficiently high energy ($> 100 \text{ ms}^{-1}$), the frequency of peak cooling is far off-resonance ($\Delta < -10\Gamma$) from the transition frequency. In this regime, photon scattering rates close to the transition frequency are more than two orders of magnitude lower, consequently a much smaller cooling force is seen for small detunings.

In Eq. 7.1, the scattering rate is implicitly dependent on the magnetic field via the magnetic dependence of the resonant transition frequency $\omega_0(B)$. Therefore, a comprehensive understanding of the beryllium energy levels is essential for laser-cooling of Be^+ in the ALPHA apparatus.

7.1 $^9\text{Be}^+$ Energy Levels

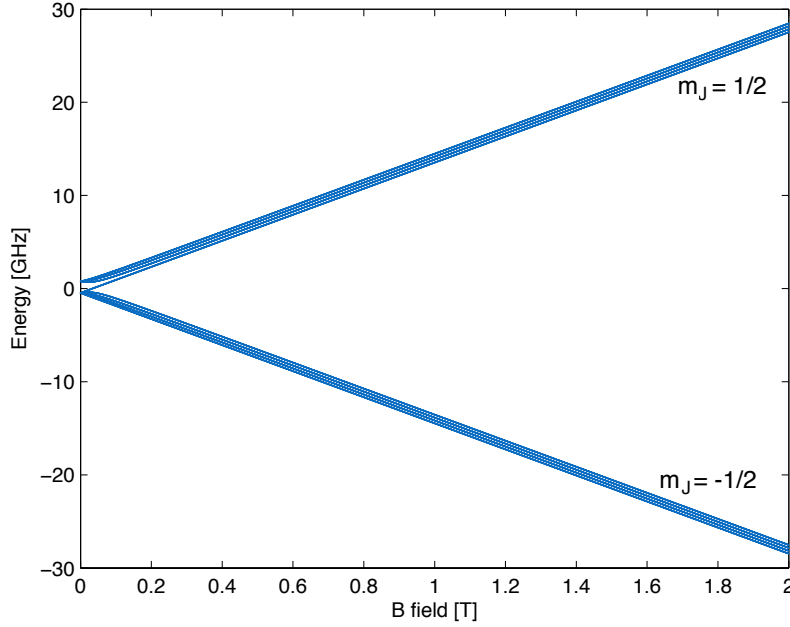
A monoisotopic element, beryllium has only one stable isotope ^9Be with four protons and five neutrons in the nucleus. A singly ionized $^9\text{Be}^+$ atom in the ground state has an electronic configuration of $1s^2 2s^1$. Interaction between the nuclear spin $I = \frac{3}{2}$ and the angular momentum of the valence electron $J = \frac{1}{2}$ leads to a hyperfine splitting of the energy levels. In the presence of an external magnetic field, the effective Hamiltonian is given by

$$H = hA\mathbf{J} \cdot \mathbf{I} + \mu_B g_F \mathbf{F} \cdot \mathbf{B} \quad (7.4)$$

where the first term is the magnetic dipole hyperfine interaction, and the second term is the Zeeman interaction. The Hamiltonian can be solved numerically to obtain the ground state energy eigenvalues at different magnetic fields.

In the Mixing Trap of the ALPHA apparatus, the magnetic field generally varies between 1 T and 2 T. Figure 7.2 shows the energy levels of the $2S_{1/2}$

Figure 7.2: $2S_{1/2}$ energy levels at high magnetic fields. The upper manifold corresponds to $m_J = \frac{1}{2}$ and the lower manifold corresponds to $m_J = -\frac{1}{2}$. At a magnetic field of 1 T, the spacing between adjacent m_I sub-levels is ~ 300 MHz.

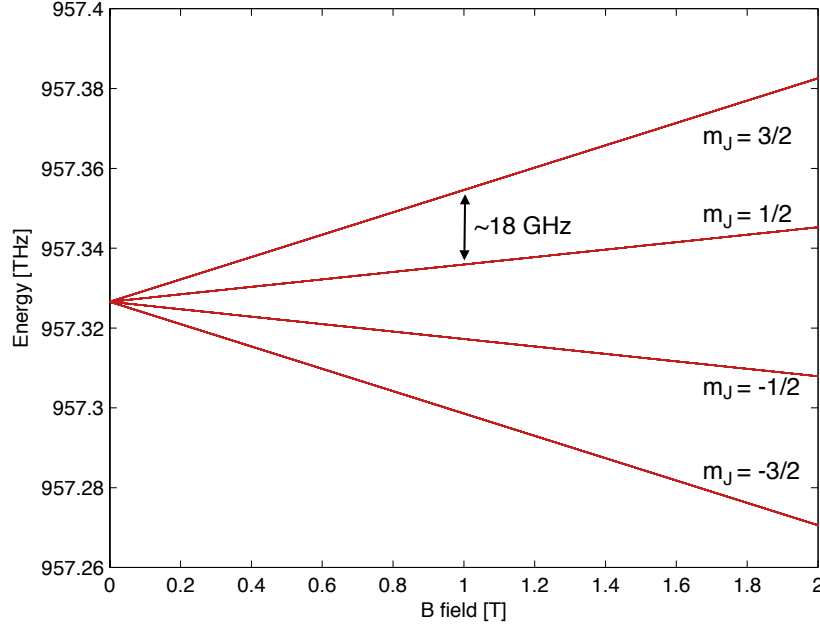


state at these high magnetic fields. The energy levels are split into two manifolds, $m_J = \frac{1}{2}$ and $m_J = -\frac{1}{2}$, with four m_I sub-levels in each manifold.

Laser cooling of ${}^9\text{Be}^+$ involves the $2S_{1/2}$ and $2P_{3/2}$ energy levels. Figure 7.3 shows the energy levels of the $2P_{3/2}$ state at fields up to 2 T. As with the $2S_{1/2}$ state, the energy levels are grouped according to their m_J values.

Figure 7.4 shows a simplified diagram for the beryllium energy levels relevant to laser cooling. The $2S_{1/2}$ ground state energy splitting between the two manifolds is small, so we can assume approximately equal initial populations in all sub-levels of the ground state. Laser cooling is typically performed by exciting the $2S_{1/2} |m_J = \frac{1}{2}, m_I = \frac{3}{2}\rangle$ sub-level of the ground state to the $2P_{3/2} |m_J = \frac{3}{2}, m_I = \frac{3}{2}\rangle$ excited state. Repump light is needed to drive the ions in the $2S_{1/2} |m_J = -\frac{1}{2}\rangle$ manifold to the $2P_{3/2} |m_J = \frac{1}{2}\rangle$ manifold, which may decay to either of the ground state hyperfine manifolds. By using σ^+ polarized light, after several transitions, the ions are

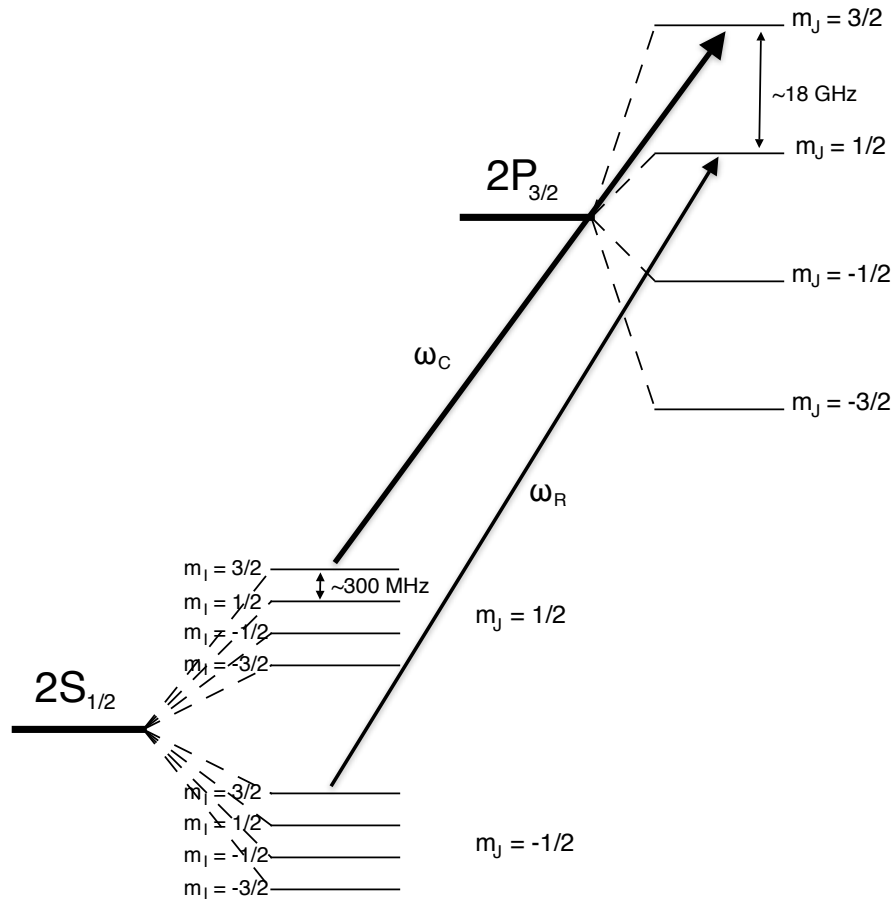
Figure 7.3: $2P_{3/2}$ energy levels at high magnetic fields. From upper to lower, the manifolds correspond to $m_J = \frac{3}{2}$, $m_J = \frac{1}{2}$, $m_J = -\frac{1}{2}$ and $m_J = -\frac{3}{2}$. At a magnetic field of 1 T, the spacing between adjacent m_J sub-levels is ~ 1.5 MHz.



pumped to the $2S_{1/2} |\frac{1}{2} \frac{3}{2}\rangle \rightarrow 2P_{3/2} |\frac{3}{2} \frac{3}{2}\rangle$ transition. Decay from this state to the $2S_{1/2} |m_J = -\frac{1}{2}\rangle$ manifold is forbidden, effectively confining any ions in this state to the closed $2S_{1/2} |\frac{1}{2} \frac{3}{2}\rangle \rightarrow 2P_{3/2} |\frac{3}{2} \frac{3}{2}\rangle \rightarrow 2S_{1/2} |\frac{1}{2} \frac{3}{2}\rangle$ cooling cycle.

At a magnetic field of 1 T, the difference between the cooling and repump transitions is $\Delta\omega = \omega_C - \omega_R \simeq 8$ GHz. The cooling laser can therefore off-resonantly repump the ions in the $2S_{1/2} |m_J = -\frac{1}{2}\rangle$ manifold with a time constant of a few seconds [98]. In contrast, the separation between the $2P_{3/2}$ m_J manifolds at 1 T is ~ 18 GHz, significantly larger than $\Delta\omega$. As a consequence, the repump rate is higher than the rate of ions decaying to the $2S_{1/2} |m_J = -\frac{1}{2}\rangle$ manifold [99]. Therefore, by appropriately red-detuning the cooling laser, it is possible to maintain a large fraction of the population in a cooling transition cycle.

Figure 7.4: Simplified energy level diagram of ${}^9\text{Be}^+$ at a magnetic field of 1 T. The difference between the cooling transition ω_C and the repump transition ω_R is ~ 8 GHz.



7.2 Simulations of ${}^9\text{Be}^+$ Laser-Cooling

In this section, we outline the atomic parameters of the beryllium ion, the electromagnetic parameters of the ALPHA Mixing Trap and the parameters of the cooling laser used to simulate the cooling of ${}^9\text{Be}^+$ ion in ALPHA. In Section 7.2.2, we provide preliminary results of our simulations and suggest possible explanations for the cooling behaviors observed.

As discussed earlier in Section 7.1, it is possible to optically pump all ground state ${}^9\text{Be}^+$ ions to the $2S_{1/2} \left| \frac{1}{2} \frac{3}{2} \right\rangle \rightarrow 2P_{3/2} \left| \frac{3}{2} \frac{3}{2} \right\rangle \rightarrow 2S_{1/2} \left| \frac{1}{2} \frac{3}{2} \right\rangle$ cooling cycle. Therefore, we only need to simulate a simplified two-level laser system with the corresponding energy separation, which depends on the local magnetic field. Laser cooling of ${}^9\text{Be}^+$ can be carried out using a 313 nm laser with σ^+ polarization.

7.2.1 Numerical Method

Laser cooling of ${}^9\text{Be}^+$ ions in the ALPHA mixing trap was simulated using MATLAB. The simulation takes into account the initial position and velocity of the simulated particles and individually calculates the trajectory of each particle during the laser-cooling process. Inter-particle interactions as well as gravity are neglected in these simulations in order to save on computation time.

As mentioned previously in Chapter 2, charged particles experience a Lorentz force in the presence of external electric and magnetic fields. Using Equation 2.1, we can derive the three-dimensional equations of motion in cartesian coordinates as follows

$$\begin{aligned}\ddot{x} &= \frac{q}{m} (E_x + v_y B_z - v_z B_y) \\ \ddot{y} &= \frac{q}{m} (E_y - v_x B_z + v_z B_x) \\ \ddot{z} &= \frac{q}{m} (E_z + v_x B_y - v_y B_x)\end{aligned}\tag{7.5}$$

For our simulations, we used MATLAB's built-in ODE45 solver to solve these differential equations of particle motion.

Trap Geometry and Fields

Figure 2.7 shows the ALPHA Penning-Malmberg trap. The central Mixing Trap consists of 13 thin electrodes of internal bore diameter ~ 44.5 mm. The total trap length is ~ 280 mm. Standard configuration of electrode potentials (Table 7.1, Figure 7.5a) for the nested-potential-well (holding both positrons and antiprotons simultaneously prior to positron EVC) result in the generation of axial and radial electric fields in the Mixing Trap (Figure 7.5b).

For Be^+ ions trapped on-axis in a 1 T magnetic field with $T \sim 40$ K, the cyclotron radius $r_c = \frac{mv_\perp}{qB}$ is of the order $10\mu\text{m}$, so the ions do not experience any appreciable radial electric fields. In our simulations, we only consider the axial electric field and ignore the radial electric field which is effectively zero within $\pm 10\mu\text{m}$ of the trap axis.

Table 7.1: Mixing Trap nested-potential-well electrode potentials.

Electrode	E10	E11	E12	E13	E14	E15	E16	E17	E18
Potential (V)	0	0	-13.5	8	-10	8	-13.5	0	0

The two mirror coils at each end of the Mixing Trap, an external 1 T solenoid and the octupole produce the minimum fields necessary to form the magnetic trap. The on-axis axial and transverse magnetic fields are shown in Figure 7.6.

For simulations discussed in this report, the transverse magnetic fields were ignored in order to reduce computation time. One can estimate the time scale of magnetic mixing between the axial and transverse motions by defining a parameter ν that characterizes the rate of relaxation between the two

Figure 7.5: (a) Mixing trap applied (red) and calculated on-axis (blue) electrode potentials. The Be^+ ions sit in the central potential well (from $z = -18$ mm to $z = 18$ mm). (b) Mixing trap on-axis axial (blue) and radial (magenta) electric fields.

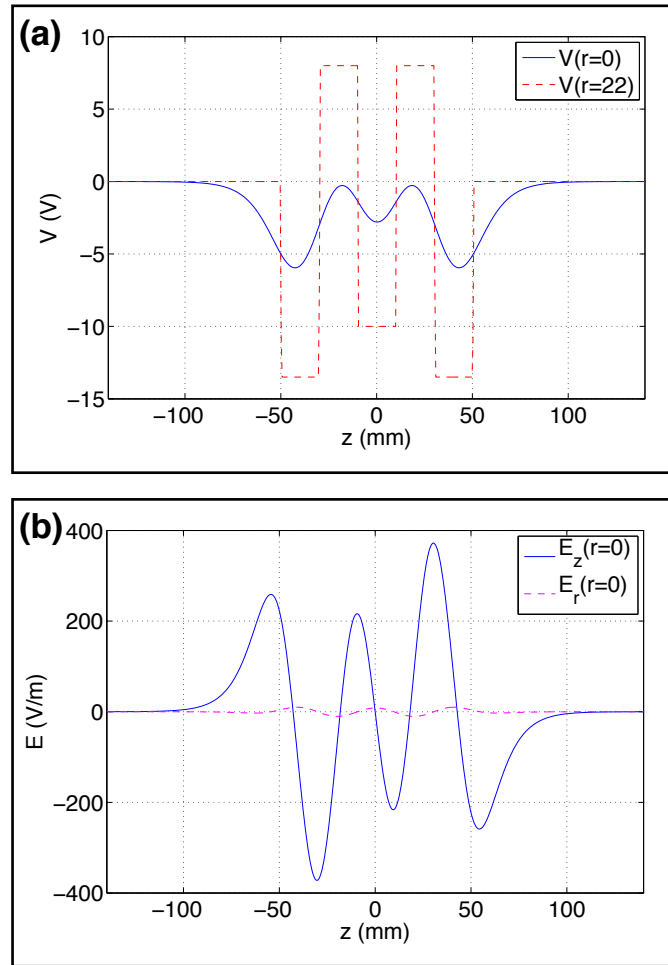
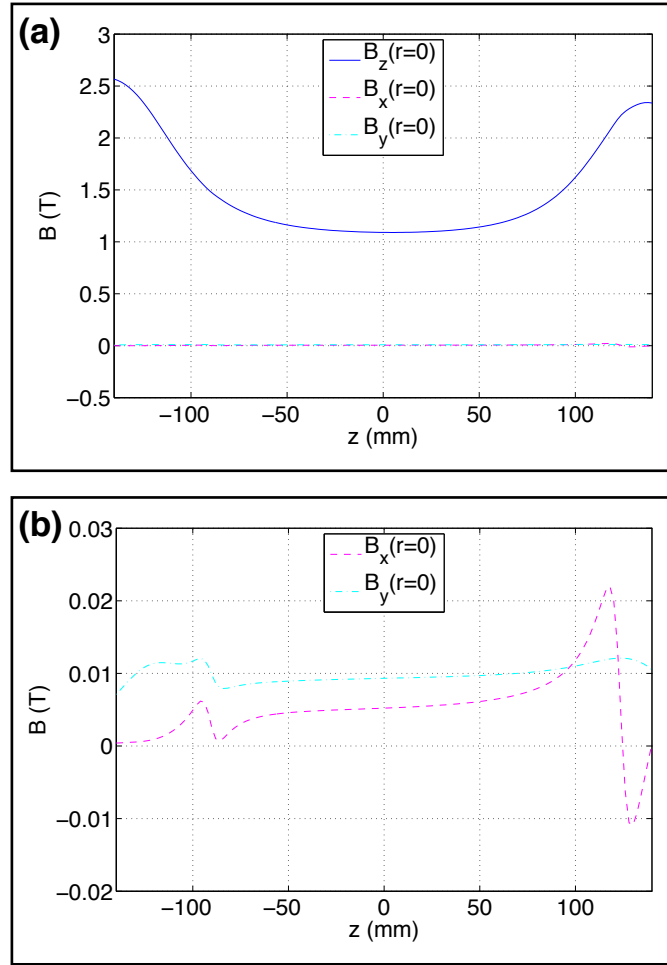


Figure 7.6: (a) On-axis axial (blue) and transverse magnetic fields. In the central potential well from Figure 7.5a, the magnetic field varies by ~ 0.01 T. (b) Zoomed-in transverse magnetic fields only.



motions. Using Newton's law of cooling, we have

$$\begin{aligned}\frac{dT_{\perp}}{dt} &= \nu (T_{\parallel} - T_{\perp}) \\ \frac{dT_{\parallel}}{dt} &= -2\nu (T_{\parallel} - T_{\perp})\end{aligned}\tag{7.6}$$

In a strongly magnetized plasma, Reference [100] gives

$$\nu \sim nb^2\bar{v} \left(\frac{r_c}{b}\right)^{1/5} \exp \left[-2.34 \left(\frac{r_c}{b}\right)^{-2/5}\right]\tag{7.7}$$

where n is the plasma density, $b = e^2/k_B T$ is the classical distance of closest approach, $\bar{v} = \sqrt{k_B T/m}$ is the thermal velocity, and $r_c = mv_{\perp}/qB$ is the cyclotron radius.

Since $r_c/b \ll 1$ in a magnetized plasma, the exponential term dominates, and magnetic mixing decreases rapidly with decreasing temperature or increasing magnetic field. Consequently, results presented here ignore the slow ($1/\nu \sim 1$ s) magnetic mixing of transverse and axial motions and effectively simulate one-dimensional laser cooling in the axial direction only. Therefore, only axial laser-cooling results are presented.

In the central potential well where the Be^+ ions are confined, the axial magnetic field strength varies by ~ 0.01 T (see Figure 7.6). This would result in a variable photon scattering rate given by Equation 7.1 due to the variation of the resonance $\omega_0(B)$ relative to the laser frequency ω . This effect has explicitly been taken into consideration for these simulations by incorporating the full axial field map. In Section 7.2.2, we will see that the trapped ions only go as far as ± 0.44 mm from the trap center where the axial magnetic field and likewise $\omega_0(B)$ are effectively constant. In the future, a custom integration algorithm would need to be developed in order to incorporate the full radial magnetic field map in the simulations without being computationally expensive.

Light Scattering

For the $2S_{1/2}$ to $2P_{3/2}$ laser cooling in ${}^9\text{Be}^+$, the natural linewidth is $\Gamma = 2\pi \times 19.4 \text{ MHz} \simeq 122 \text{ MHz}$. σ^+ light at $\sim 313 \text{ nm}$ is used to drive this cooling transition. For photons of wavelength $\lambda = 313 \text{ nm}$, the recoil velocity of ${}^9\text{Be}^+$ for absorption or emission is

$$v_{\text{recoil}} = \frac{h}{m\lambda} = 0.1414 \text{ ms}^{-1} \quad (7.8)$$

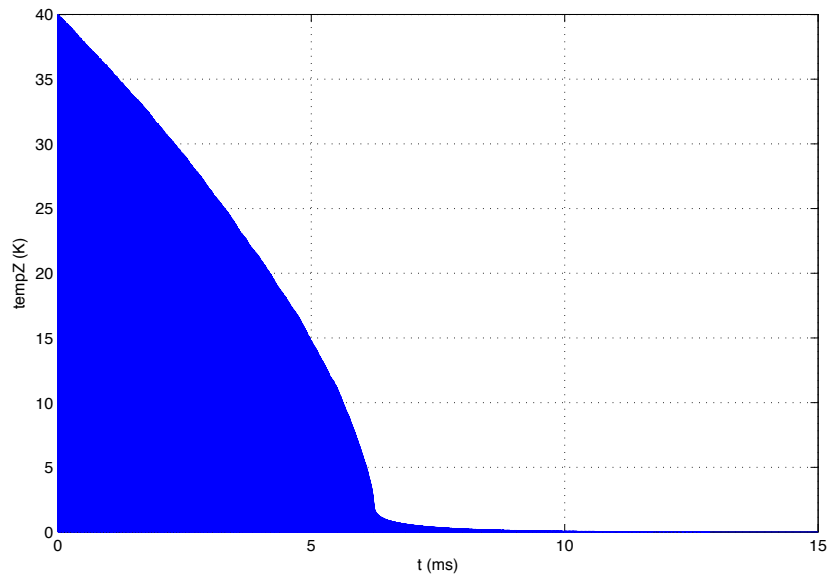
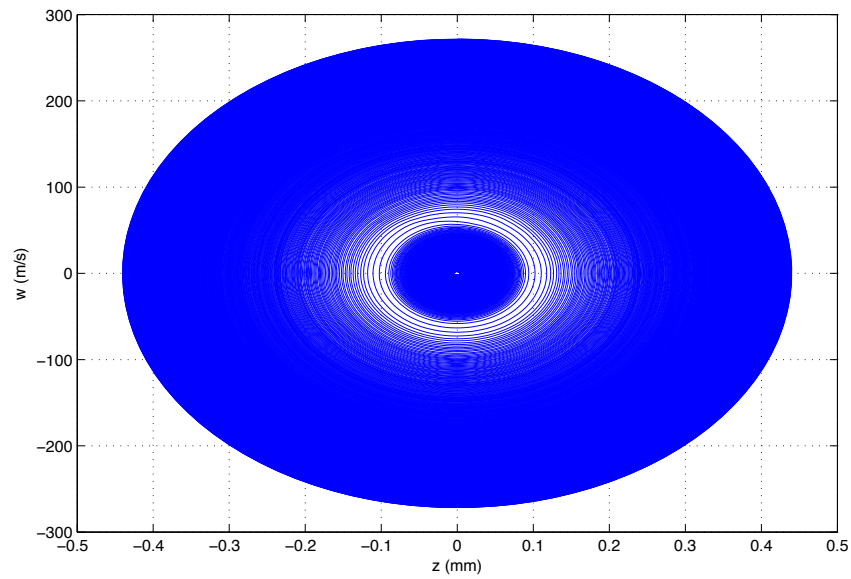
For our simulations, we assumed a laser output power of 5 mW with a beam waist of 3 mm inside the trap. This corresponds to an intensity close to the saturation intensity, with $I/I_{\text{sat}} = 0.86$.

7.2.2 Simulation Results

As a starting point, we simulate the laser cooling of a single ${}^9\text{Be}^+$ ion with an axial energy of 40 K (in temperature units) at the center of the ALPHA mixing trap. The choice of this temperature is motivated by the fact that we can easily cool trapped beryllium ions down to tens of kelvin by mixing them with the 20 K positron plasmas normally prepared at ALPHA. The positrons radiate energy via cyclotron emission and sympathetically cool the beryllium ions via Coulomb collisions. Once the mixed positron-beryllium plasma has thermalized, laser-cooling of beryllium can commence. For our simulations, we conservatively assume an equilibrium temperature of 40 K when laser-cooling begins.

Figure 7.7 shows a graph of the axial energy of a single Be^+ ion as a function of time, being cooled with an on-axis 313 nm laser red-detuned by 10Γ . The initial cooling rate is very high, with the ion energy falling from 40 K to 5 K in just over 6 ms. Beyond a critical turning point around 7 ms, the cooling rate drops drastically: it takes another 7 ms for the energy to fall below 5 mK.

Eq. 7.1 and Figure 7.1 can be used to explain the cooling behavior observed here. At 40 K ($\sim 272 \text{ ms}^{-1}$), there is a high photon scattering rate and

Figure 7.7: Plot of axial energy vs. time.**Figure 7.8:** Axial z vs v_z phase-space plot. The ion starts at $z = 0$ and $v_z = 272$ ms^{-1} and then spirals inwards in phase-space as it cools down.

a moderate cooling force at 10Γ red-detuning, working to reduce the axial energy of the ion. As the ion cools, F_{cooling} increases until the ion reaches 2 K ($\sim 61 \text{ ms}^{-1}$), where the cooling rate is at its peak. For energies lower than 2 K (v_z less than $\sim 60 \text{ ms}^{-1}$), the cooling force at 10Γ decreases, thus we observe a marked drop in the ion cooling rate. At this moment, the laser detuning is too large and off-resonant from the cooling transition frequency and the photon scattering rate is low as a result. Lowering the detuning at this point would subsequently increase the cooling rate.

Figure 7.8 is a phase-space plot of the axial displacement z and velocity w for the same ion laser-cooled from an energy of 40 K. With an initial displacement range of $\pm 0.44 \text{ mm}$ and velocity range of $\pm 272 \text{ ms}^{-1}$, the ion is cooled down and slowly spirals into the center of the phase-space plot.

Figure 7.9: Plot of cooling times vs. initial laser detuning for a ${}^9\text{Be}^+$ ion at 40 K.

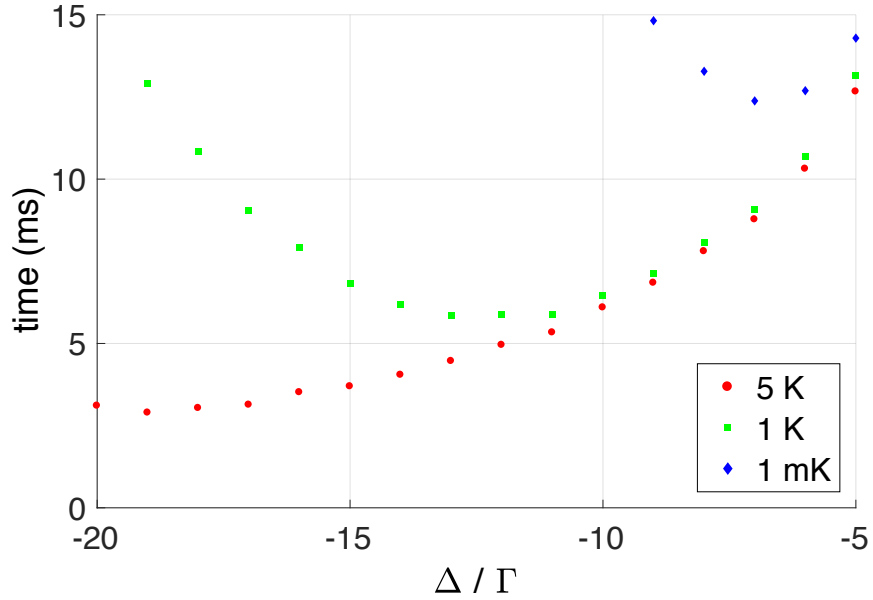
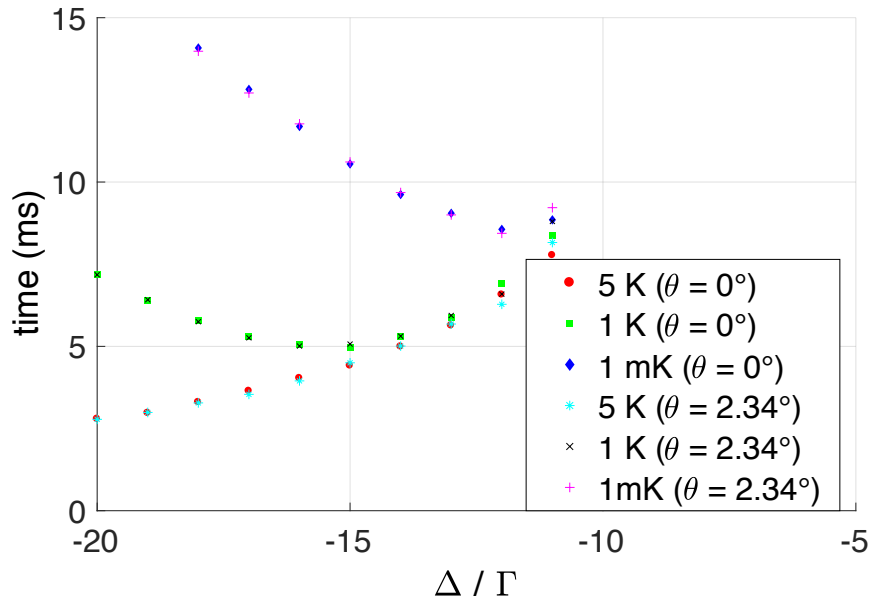


Figure 7.9 is a plot of ion cooling times versus detuning of the cooling laser. For large detuning, it takes less than 5 ms for the ion to be cooled to 5 K. However, it takes an extremely long time to reach energies of a few mK.

As the detuning is decreased, the initial cooling rate is only slightly lower, but sub-kelvin energies are reached in less than 10 ms. For laser detuning between 9Γ and 5Γ , sub-milliKelvin energies are reached in less than 15 ms. For reference, the 313 nm Doppler laser cooling temperature limit for ${}^9\text{Be}^+$ ions is 0.5 mK.

Figure 7.10: Plot of cooling times vs. initial laser detuning including laser chirping at 100 GHz/s.

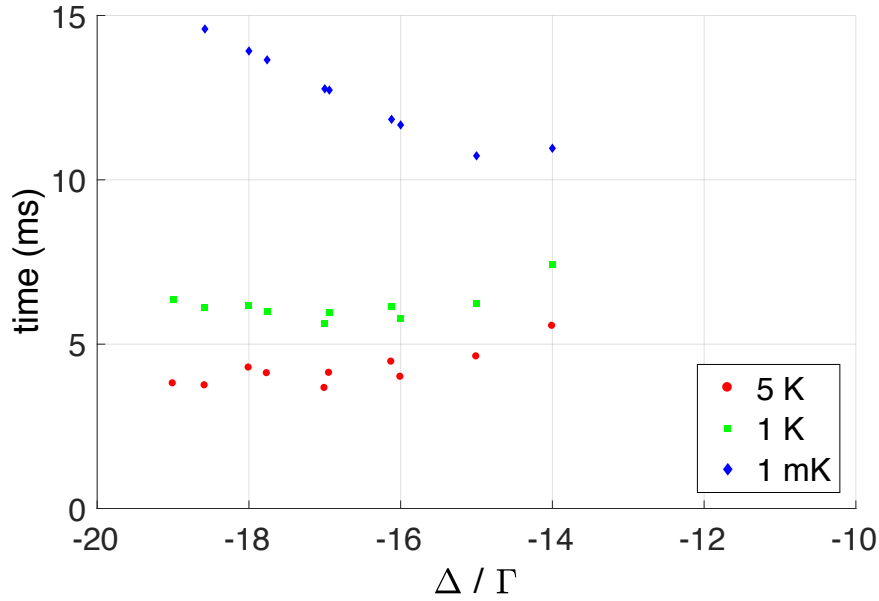


In order to speed up the cooling process, chirping of the laser would be required. From Figure 7.9, we estimate an ideal chirping rate of around 100 GHz/s that is still within the technical capabilities of our laser system. We simulate the cooling of a 40 K ion, this time with chirping at the aforementioned rate. Figure 7.10 shows the results. An initial detuning of less than 15Γ with chirping is sufficient to reach sub-milliKelvin energies within 10 ms.

The ALPHA apparatus is currently designed to deliver the cooling laser at an angle of $\theta = 2.34^\circ$ with respect to the trap axis. Therefore, we also simulate ion cooling with the laser being delivered at this angle. The results are plotted in Figure 7.10. The cooling times are nearly identical to those

obtained with the laser delivered on-axis. Small angular variations in the laser delivery system should therefore not have any significant impact on our laser cooling abilities.

Figure 7.11: Plot of average cooling times vs. initial laser detuning for a distribution of ${}^9\text{Be}^+$ ions centered at 40 K.



Next, we simulate the laser cooling of an ensemble of 100 ${}^9\text{Be}^+$ ions, with a Gaussian temperature distribution centered at 40 K. While the simulations ignore inter-particle interactions, they are useful to estimate the average laser-cooling behavior of the ensemble from a given distribution. Figure 7.11 is a graph of average cooling times for this ensemble against initial laser detuning. For detuning below 14Γ , the average ensemble temperature remains above 5K at all times. The region between 14Γ and 16Γ appears to be the ideal region for initial red-detuning of the laser with a chirping rate of 100 GHz/s. Beyond 16Γ , it takes longer and longer for the ensemble to reach sub-milliKelvin temperatures.

For the sympathetic laser-cooling of positrons with Be^+ ions, it would be preferable to have ions not colder than 1 K in order to avoid the prema-

ture centrifugal separation of the two charged species. We would therefore ideally have the cooling laser far red-detuned from resonance. In addition to preventing the ions from cooling too quickly, this would have the added benefit of providing enough time for collisional temperature relaxation and thermalization between the axial and transverse degrees of freedom, thereby effectively cooling the ions in 3-D.

While the weak radial electric and magnetic fields have been ignored in this study, their presence is expected to speed up the mixing of particle axial and radial motion, thereby enhancing the indirect laser-cooling rate in the transverse direction. In contrast, since the cooling laser comes in almost parallel to the axis of the trap, there is going to be very little direct laser cooling in the transverse direction compared to the axial direction. For this reason we neglect the Doppler cooling due to cyclotron motion in our study. Further discussion on the laser-cooling of ions using a single axial laser can be found in References [\[101\]](#) and [\[102\]](#).

The preliminary 1-D simulations results presented here support the feasibility of laser-cooling beryllium ions in the ALPHA Mixing Trap. Once a suitable beryllium ion source is ready, a more comprehensive investigation will be undertaken that takes into account the full electromagnetic geometry of the ALPHA apparatus and studies the effects of laser-cooling in all three dimensions.

Chapter 8

Ion Production Using Pulsed Laser Ablation

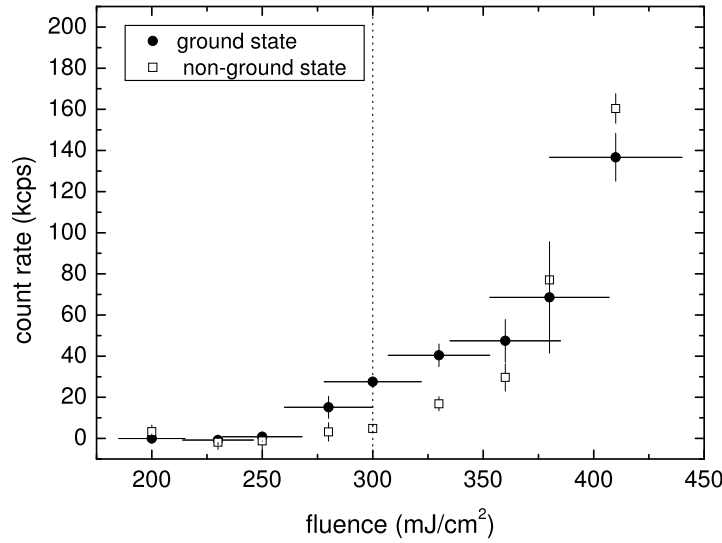
The removal of material from a surface by high-intensity short laser pulses is known as pulsed laser ablation (PLA). A wide variety of microscopic mechanisms compete during the ablation process: material may desorb due to the evaporation of the target surface from heating; alternatively, the laser pulse may induce mechanical stress and defects, resulting in the release of material; if the photon energy is sufficiently high, direct bond-breaking may also occur resulting in material removal from the target. A detailed treatment of ablation mechanisms based on several different models can be found in Reference [\[103\]](#).

The process of laser ablation can lead to the creation of atoms, ions, molecules, and even large clusters and fragments [\[103\]](#). In this chapter, I will give a brief review of the present understanding of the ablation of metals, and the potential for using this technique in the production of laser-ablated Be^+ ions. I will briefly describe the constraints imposed by the ALPHA experiment for the design of a compatible ion source.

8.1 Laser Ablation of Metals: A Brief Review

The ablation of metal targets with nanosecond laser pulses has been extensively studied during the past three decades [104–116]. While an exhaustive study of any metallic element has never been carried out, most of the investigations involve two or more metals which can be directly compared. By collating the results from several experiments, a coherent picture emerges that describes how the production of metallic ions is affected by different parameters in the ablation process. This understanding in turn informs our choice of ablation parameters in our attempt to create low-energy beryllium ions for the sympathetic cooling of positrons in ALPHA.

Figure 8.1: Production rate of ground state and excited state (or ionized) atoms of calcium as a function of laser fluence. To the right of the dashed line, a significant fraction of the particles are ions. From [113].



Investigations conducted so far reveal that for the ablation of metals with nanosecond laser pulses, the properties of the ablated plume are correlated with a parameter known as the “fluence”, i.e. the laser pulse energy per unit area. The first observation is that below a certain fluence, no ablation takes place. Figure 8.1 plots the production of atoms (labeled “ground state”) and ions (labeled “non-ground state”) as a result of ablating a calcium target with 1064 nm wavelength laser pulses (10 ns pulse length). At fluences below 0.25

Jcm^{-2} , almost no atoms or ions are produced. At 0.3 Jcm^{-2} and above, the production of both atoms and ions is observed. This suggests the existence of an “ablation threshold” such that atoms and ions are only produced in significant quantities if fluences above threshold are employed in the ablation process.

A second thing to note in Figure 8.1 is that the ion fraction increases with increasing fluence. At a fluence of 0.3 Jcm^{-2} , the ablated material is dominated by neutral atoms. Near 0.4 Jcm^{-2} however, the number of ions detected is almost equal to the number of atoms in this study. Therefore, for a given amount of ablated material, the ion fraction can be tuned by modifying the laser pulse energy and/or spot size.

Figure 8.2: Ion yield as a function of laser fluence for ablation of a copper target. Adapted from [107].

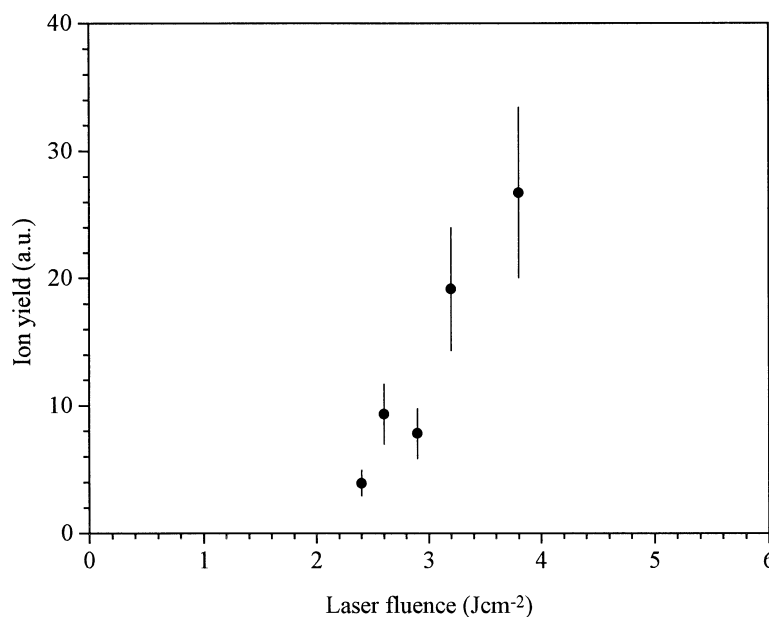
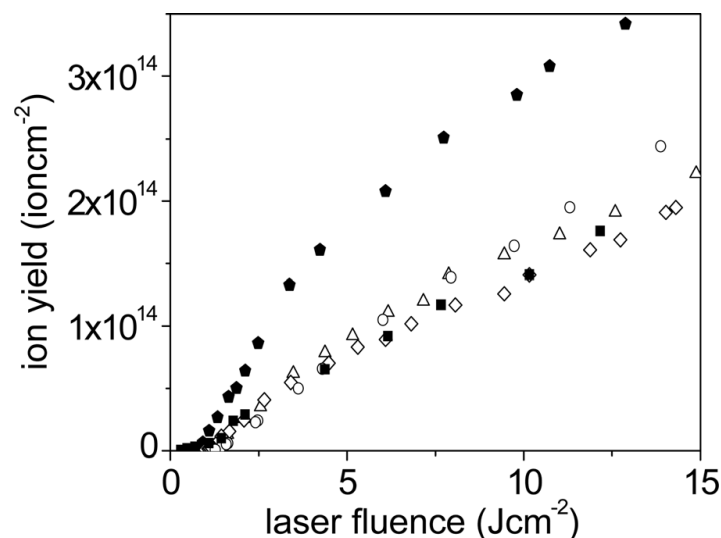


Figure 8.2 shows the ion yield with respect to fluence for a different experiment in which 351 nm wavelength laser pulses (16 ns pulse length) were used to ablate a copper target. As before, the increase in laser fluence results in an increase in ion production, with the relationship being approximately linear within experimental error. By extrapolating the linear fit backwards,

one can approximate the ablation threshold as the fluence value where the linear fit crosses the horizontal axis. This is in fact the standard method with which ablation thresholds are obtained experimentally.

Figure 8.3: Ion yield as a function of laser fluence for ablation of several different metal targets, including Al (■), Cu (○), Ag (◇), Au (△), and Bi (closed pentagon). From [112].

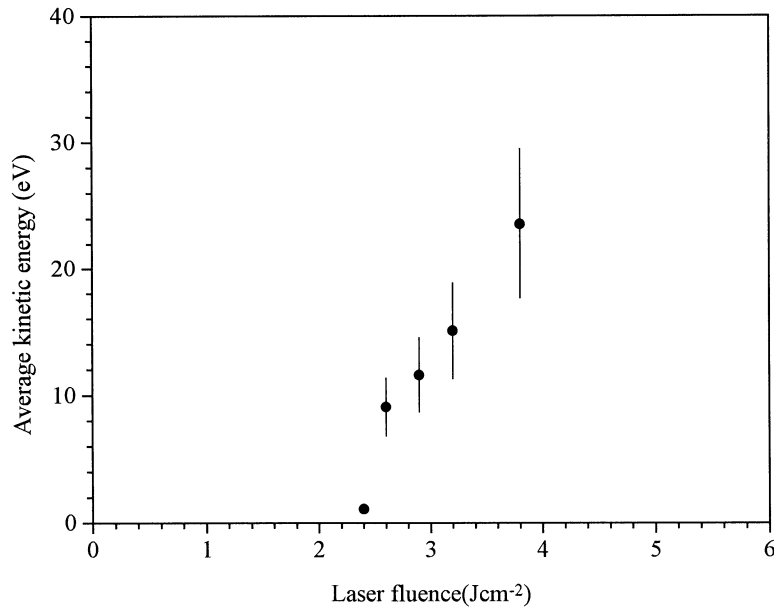


For laser ablation of metal targets using nanosecond pulses, it has repeatedly been observed that most metals have very similar ablation thresholds and ion yields. Figure 8.3 shows the ion yields of five different metals: aluminum, copper, silver, gold, and bismuth. In this study, the ablation targets were impinged with 193 nm wavelength laser pulses (20 – 25 ns pulse length). All five metal species were observed to have ablation thresholds between 0.5 J cm^{-2} and 1.5 J cm^{-2} . Furthermore, apart from bismuth, the ion yields of all of the metals under investigation increase by almost the same amount with increasing fluence.

The deviation of bismuth can be attributed to the significantly lower melting point it has (545 K) compared to the other metals under investigation (933 K, 1358 K, 1235 K, and 1337 K for Al, Cu, Ag, and Au respectively). This is in line with the expectation that ablation dominated by thermal ef-

fects (melting and evaporation due to heating) will result in lower ablation thresholds and higher ions yields for metals with lower melting points. If this model also applies to beryllium (melting point: 1560 K), we can expect it to have ion yields much lower than bismuth and aluminum for ablation at the same fluence.

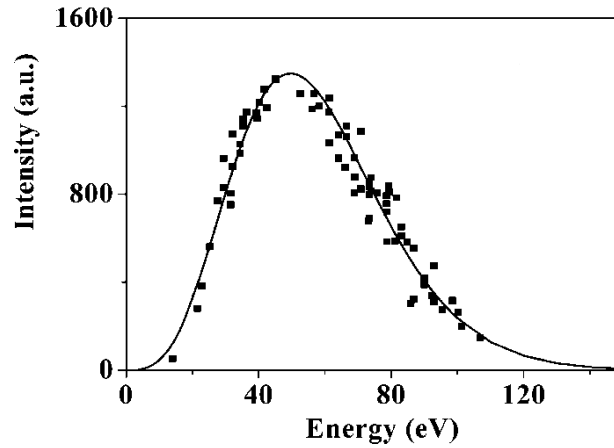
Figure 8.4: Average ion kinetic energy as a function of laser fluence for ablation of a copper target. Adapted from [107].



For the trapping and laser-cooling of beryllium ions, an important consideration is the kinetic energy distribution of the ablated ions. If the ions have a high average kinetic energy and/or a large energy spread, it would be much more difficult to trap and cool them since stronger electric fields and longer traps would be required. For the same data shown in Figure 8.2, Figure 8.4 shows how the average kinetic energy of ablated ions typically varies with fluence. For ablation just above threshold, the ions have less than 1 eV average kinetic energy. For higher fluences however, a linear increase in the average kinetic energy is observed. Carrying out ablation just above threshold could therefore be used to produce as slow ions as possible. Low fluences unfortunately also result in a low ion fraction. Therefore, for the purpose of creating Be⁺ in ALPHA, an intermediate laser fluence would be needed to

optimize between the desired ion production and ion kinetic energy.

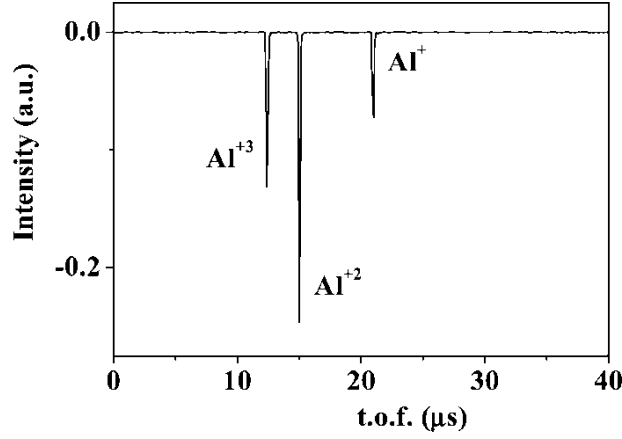
Figure 8.5: Maxwell-Boltzmann kinetic energy distribution of Al^+ at a fluence just above threshold. Adapted from [111].



Since we are interested in trapping low energy metal ions, it is worthwhile knowing not just their average kinetic energy, but the entire energy distribution. Figure 8.5 shows the kinetic energy distribution of the Al^+ ions that are ablated using 532 nm laser pulses (5 ns pulse length) at a fluence of 1.1 Jcm^{-2} , just above the threshold for singly charged ions. The closed squares in the plot are experimental data points. A Maxwell-Boltzmann distribution, shown by the solid line, is an excellent fit to the data, indicating that the ablated plume of ions is in a thermodynamic equilibrium. Using laser fluences just above threshold not only ensures a lower average kinetic energy, but also a higher fraction of ions at the lower end of the energy distribution.

For ion sources where only a particular charge species is of interest, the accompaniment of ions of lower or higher charge needs to be determined as well. When ions are produced using PLA, one way to ascertain whether several different charged species are produced is to look at time-of-flight data. Figure 8.6 shows the time-of-flight spectrum as a result of ablating an aluminum target with 532 nm pulses (5 ns pulse length) at a fluence of 5.3 Jcm^{-2} . Three distinct peaks, corresponding to Al^+ , Al^{2+} , and Al^{3+} are clearly visible. In this particular situation, 5.3 Jcm^{-2} is above the fluence threshold of both

Figure 8.6: Time-of-flight spectrum as a result of ablating an Al target at a fluence of 5.3 Jcm^{-2} . Adapted from [111].

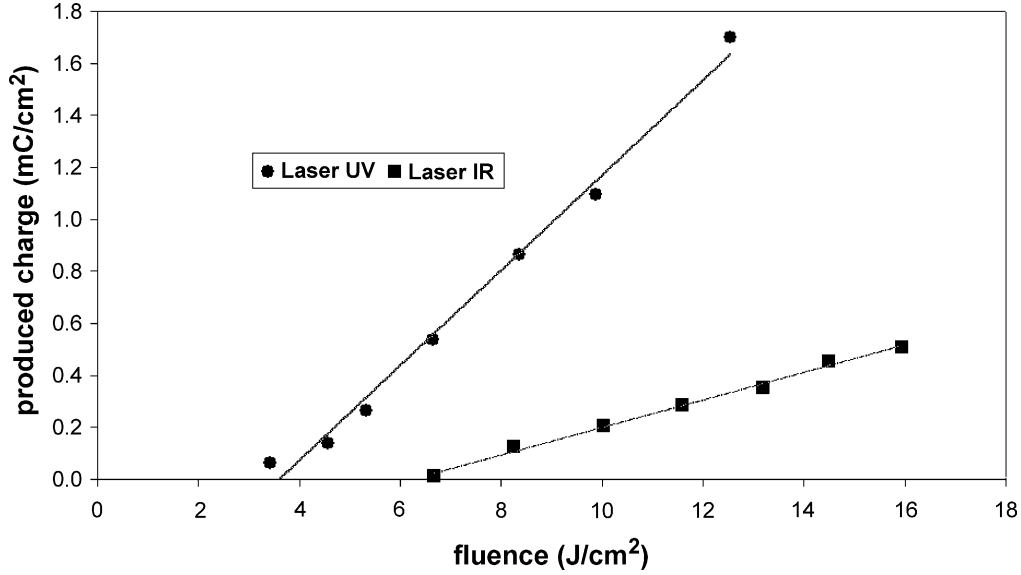


Al^{2+} and Al^{3+} ions. By carrying out ablation and measuring the time-of-flight spectrum for several different fluences, the threshold fluence for the production of Al^+ , Al^{2+} , and Al^{3+} is obtained.

At ALPHA, we intend to laser-cool only singly-charged beryllium ions for the sympathetic cooling of positrons. Therefore, ablation of a beryllium target should ideally take place with fluences above the fluence threshold of Be^+ but below that of Be^{2+} to prevent the production of unwanted Be^{2+} ions.

So far, we have looked at how ablation thresholds, ion yields, and kinetic energy distributions depend on the ablation of metals at different fluences. These three quantities also depend on the wavelength of the laser pulses used. Figure 8.7 is a plot of the ion yield as a result of ablating a copper target at two different laser wavelengths (308 nm wavelength at 9 ns pulse length and 1064 nm wavelength at 20 ns pulse length). For the same laser fluence, the UV laser pulses produce a significantly larger number of ions (by at least a factor of 4) compared with when IR laser pulses are used. Furthermore, the fluence threshold for the production of ions is approximately a factor of 2 smaller in the case of the UV laser. Therefore, in an ion source where a higher total ion fraction is preferable, it would be more pertinent to carry out ablation with shorter wavelength UV lasers rather than longer wavelength

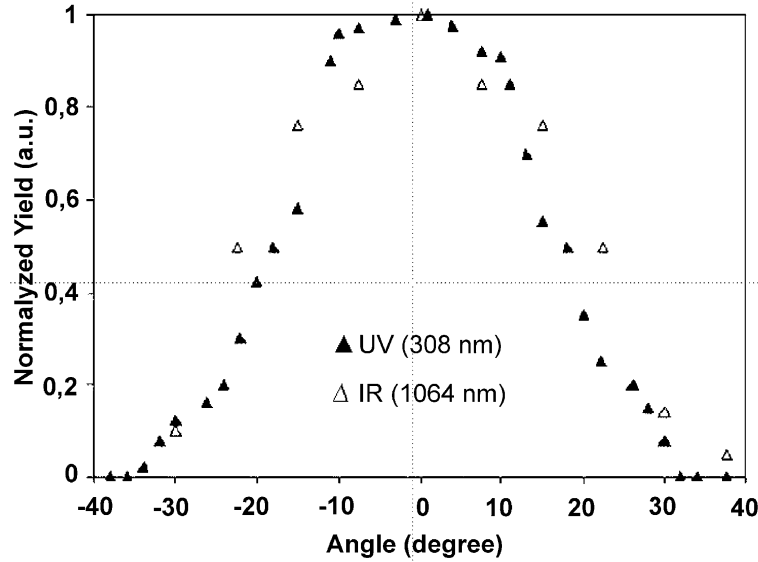
Figure 8.7: Ion yield as a function of laser fluence for ablation of a Cu target at two different wavelengths. From [109].



visible or IR lasers. Using UV lasers would also however mean a lower fluence threshold for doubly and triply charged ions, thereby reducing the range of fluences that may be used for ion production.

In the design of an ion source using PLA, the angular distribution of the ablated species is an important consideration, since maximizing the fraction of trapped, low transverse-energy ions in a Penning-Malmberg trap means having ions with as little radial divergence as possible. Figure 8.8 shows the variation in ion yield with angle from the surface normal when a UV/IR laser is used to ablate a copper target (the UV laser has an impact angle of 70° , while the IR laser has an impact angle of 45°). The ion production is evidently not isotropic, but rather has a maximum normal to the target surface for both laser wavelengths used. At angles of 30° or more, the ion yield is $< 10\%$ relative to the surface normal. Compared to the isotropic ion production from an oven, PLA is better-suited for the production of a collimated ion beam, since most of the ions are already emitted with relatively small solid angles. Given the desire to produce a beam of Be^+ ions to be trapped in the ALPHA Penning-Malmberg trap, this feature of maximal

Figure 8.8: Normalized ion yield as a function of angular distribution due to the ablation of a Cu target at two different wavelengths. From [109].



production normal to the target surface fixes the optimal orientation of the target in the ion source.

Although no comprehensive study on the pulsed laser ablation of any single element exists, the similarity of results for a variety of metals allows us to draw general conclusions with regards to the ablation of metals from nanosecond laser pulses:

- There is a minimum fluence required for the production of ions using pulsed laser ablation. This is known as the ablation threshold or fluence threshold.
- For a given metal target, there is a fluence threshold for singly charged ions, a larger fluence threshold for doubly charged ions, and an even larger one for triply charged ions etc.
- The fluence threshold is lower for ablation carried out with shorter wavelength laser pulses.
- The ion yield increases with increasing fluence.

- The ion yield increases with decreasing wavelength of the laser pulse.
- Ions ablated close to the fluence threshold approximately follow a Maxwell-Boltzmann kinetic energy distribution
- The average ion kinetic energy increases with increasing fluence
- The angular distribution of ions is peaked in the forward direction normal to the target surface

Armed with these general principles about the ablation of metals, we can begin to explain the suitability of producing beryllium ions using this technique, particularly in comparison to the research community norm of using an oven. Eventually, given the spatial constraints of the ALPHA apparatus and the temporal constraints of a PhD degree, we will discuss the design and execution of an experimental program that would result in a beryllium ion source compatible with the ALPHA experiment and suitable for the sympathetic cooling of positrons in the Mixing Trap at ALPHA.

8.2 Ion Source Compatibility with ALPHA

The synthesis of antihydrogen in the ALPHA Mixing Trap takes place in one of the most extreme conditions found in a laboratory setup: the trap electrodes are at cryogenic temperatures (~ 6 K); the trap vacuum where antihydrogen is made is at 10^{-13} mbar or better; the magnetic fields are non-homogenous and vary over two orders of magnitude up to 3 T; custom low-noise electronics are used to provide the necessary electric fields; and finally, there are severe mechanical and spatial limitations that need to be overcome for the apparatus to function as an antihydrogen synthesizing device.

With these serious constraints presented by the ALPHA apparatus, the design of a beryllium ion source using pulsed laser ablation needs to be carefully planned, optimized, and ideally demonstrated in an external setup prior to

integration into ALPHA. In particular, the ion source needs to meet the following criteria:

1. The thermal load of the ion source should be minimal with no detrimental effects on the rest of the apparatus.
2. The process with which ions are produced should not cause pressure spikes in the UHV region of the trap, and should not deteriorate the baseline trap vacuum in the long run (several weeks to months).
3. The ion source should produce a tunable and reproducible number of trappable ions on demand, taking into consideration the electric and magnetic field limitations of the trap.
4. The ion source should be easily integrable into the ALPHA apparatus without any major modifications to the existing setup, and without undermining the operation of existing components.

The remaining chapters in this thesis deal with the design and implementation of a laser-ablated beryllium ion source that has been tested to meet the stringent criteria of the ALPHA trap environment, and is being prepared for integration into the main apparatus in the near term.

Chapter 9

Ablation Testing Apparatus

In this Chapter, I outline the design of an apparatus that has been used to validate the efficacy of an ablation ion source suitable for installation in ALPHA.

9.1 Ablation Laser

The laser used to carry out the ablation measurements is a pulsed Nd:YAG laser (manufacturer: Quantel, model: Brilliant) with an output wavelength of 1064 nm in the infrared region of the electromagnetic spectrum. The laser pulses are 5 ns long with an average pulse energy of 40 mJ and a spot size of 5 mm. The laser operates at a factory-default repetition rate of 20 Hz.

The laser system can be operated in two independent modes, Auto and External. In Auto mode, the laser operates continuously, emitting a pulse every 50 ms based on a clock internal to the laser system. In External mode, the laser pulse is triggered by a 5 V TTL external input. This allows laser pulses to be produced on demand, allowing the emission of a laser pulse to be easily synchronized to the rest of the experiment. There is a down side to operating in External mode however: since the thermal stabilization of the laser is only good to within a few degrees, pulses produced in External mode can have a large deviation from the average energy depending on the temperature of

the laser crystal. This corresponds to a large error bar on any measurements done with the laser in External mode. In Auto mode however, the continuous operation “warms up” the laser to a steady temperature after a couple of hours, resulting in the production of pulses that are within 10% of the average energy. The reproducibility of pulse energy comes at the cost of a $\sim 2 - 3$ hours of warmup time. In all of the results presented in this thesis, the laser was operated in Auto mode, and measurements were only made after laser pulse stability was verified.

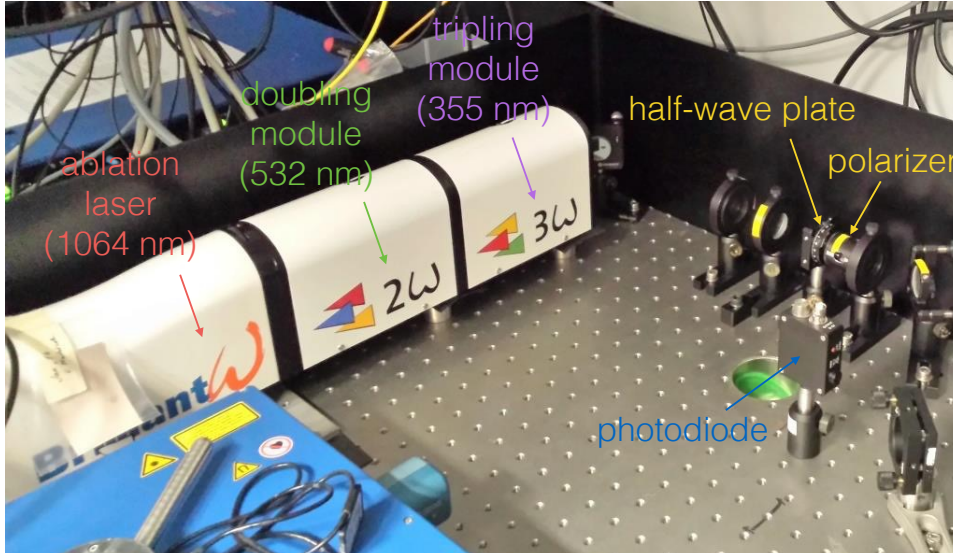
As mentioned in Section 8.1, using a shorter wavelength for ablation typically results in a higher fraction of ions among the ablated particles. Luckily, the Quantel Brilliant line of products includes laser frequency doubling and tripling modules that allow the production of laser pulses at 532 nm and 355 nm respectively from the 1064 fundamental. Generating shorter wavelengths is thus a simple matter of adding the doubling and tripling stages to the main laser, as shown in Figure 9.1. This gives us the ability to perform pulsed laser ablation with a preferred lower wavelength of 355 nm. All of the laser ablation measurements discussed in this thesis were carried out with 355 nm ultraviolet pulses.

9.2 Beam Optics

The frequency tripling stage of the ablation laser outputs ~ 40 mJ pulses at 355nm. As we will discover in the next chapter, less than 0.5 mJ of pulse energy is actually required to carry out the ablation experiments. We must therefore have a mechanism to reduce the energy of the laser beam before it is sent to the experimental area. This attenuation is provided by a simple optical attenuator comprising a linear polarizer and a half-wave plate, as shown in Figure 9.1.

The optical attenuator works on a simple principle. First, the polarized beam from the laser passes through the half-wave plate where its polarization

Figure 9.1: Photograph of the ablation laser situated inside the ALPHA laser lab. The frequency doubling and tripling modules are labeled 2ω and 3ω respectively. Also shown are some of the beam optics, in particular a beam attenuator that reduces the energy of the transmitted laser pulse, and a photodiode that measures the pulse energy after attenuation.



angle is rotated. Rotation of the half-wave plate by ϕ degrees rotates the polarization angle by 2ϕ degrees. Next, the beam passes through the polarizer which only transmits light along its polarization axis. The intensity of light I transmitted by the polarizer is given by

$$I = I_0 \cos^2(\theta) \quad (9.1)$$

where I_0 is the initial intensity of the linearly-polarized beam and θ is the angle between the beam's initial direction of polarization and the axis of the polarizer. Combining the half-wave plate and the polarizer produces a tunable optical attenuator where the transmitted beam intensity is given by

$$I = I_0 \cos^2(2\phi) \quad (9.2)$$

By rotating the half-wave plate between 0° and 45° , continuous attenuation of the beam is achieved from $I = I_0$ to $I = 0$. In practice, the minimum amount of transmitted light is determined by the quality of the polarizer and the half-wave plate as well as the polarization of the initial beam. In the setup shown in Figure 9.1, a minimum of $\sim 50\mu\text{J}$ pulse energy is always

transmitted.

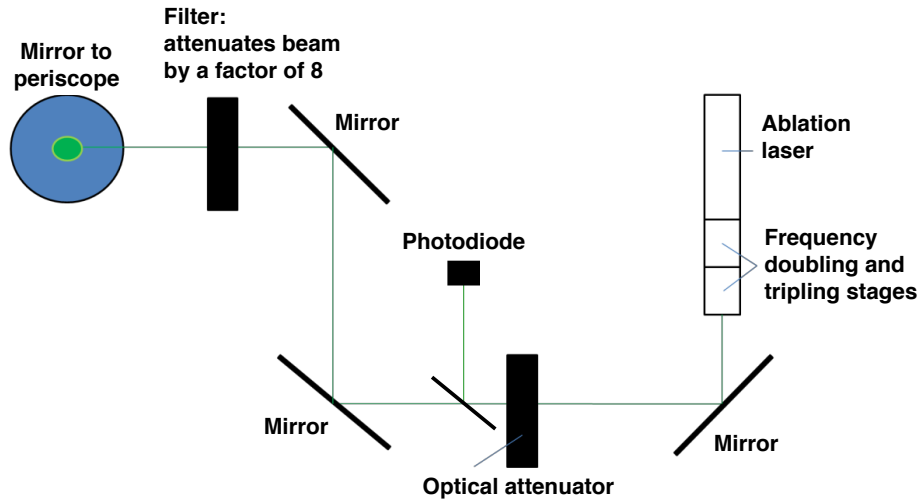
After the laser beam is attenuated, it passes through a pickoff mirror which allows almost 99% of the light to pass through but reflects $\sim 1\%$ towards a calibrated photodiode. This photodiode converts the incoming light into an analog signal (0 V to 5 V). The photodiode response can therefore be used to extract the energy of the pulse used in an ablation experiment.

The ablation laser mentioned above is a Class 4 laser, and therefore, for safety reasons, is installed in a dedicated laser lab that ALPHA uses to house all of its high power lasers. The ablation chamber (see Section 9.4), where all of the ablation experiments take place, is situated next to the wall just outside the laser lab. Transporting the laser pulse from the lab to the experimental area therefore requires a complicated set of beam optics that have to be securely mounted, aligned, and fixed into position in such a way that the laser pulse propagates all the way to the experiment and the beam path remains stable on a day-to-day basis.

Figure 9.2 shows a sketch of the optics deployed inside the laser lab. After passing through the optical attenuator described above, the laser beam is directed via high-reflectivity mirrors to a periscope. If pulse energies $< 50\mu\text{J}$ are needed, an additional filter may be used that attenuates the laser pulse by a constant factor of 8. This filter sits just before the mirror to the periscope.

The periscope is basically a set of mirrors enclosed in a tube that directs the beam path from inside the laser lab, through the wall of the room, to the outside of the laser lab where the experimental bench is located. From the output of the periscope, the laser beam is reflected by another set of high-reflectivity mirrors, through a focusing lens and into the ablation chamber, as shown in Figure 9.3. The mirrors in the experimental area are adjustable in order to correctly align the beam path. Since the 355 nm wavelength is invisible to the human eye, a visible green laser pointer is used as an “alignment laser” in order to optimize the orientation of the mirrors such that the

Figure 9.2: Layout of the beam optics in the laser lab. Light from the ablation laser is reflected by mirrors and passes through a couple of attenuators on its way to the input mirror of the periscope. The periscope then delivers the beam from inside the laser hut to the experimental setup situated next to the wall outside. Components not drawn to scale. Adapted from [117]



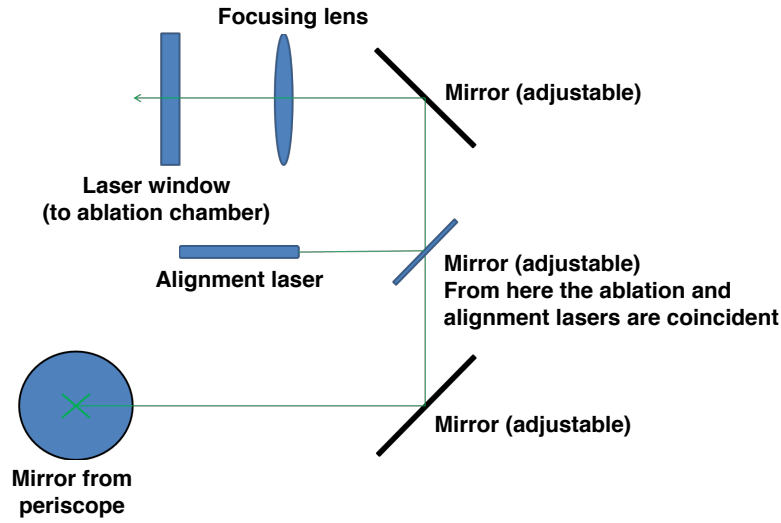
laser beam passes through the center of the window of the ablation chamber. The lens that sits just before the window has a focal length of 25 cm and is used to focus the ablation pulse onto the ablation target in the ion source.

9.3 Ion Source

Figure 9.4 shows a schematic of the ablation ion source which is composed of four key components: the UV mirror, the ablation target, the acceleration plate, and the shield. Focused light coming from the laser bench (bottom of the figure) enters the ion source and is reflected by the mirror onto the ablation target. The orientation of the mirror is such that the reflected light makes an angle of 50° with the target surface.

If the incident laser pulse has sufficient intensity, it will result in the release of atoms and ions from the target surface. The number, kinetic energy and charge of the ablated particles depend on a number of factors. Laser pulse

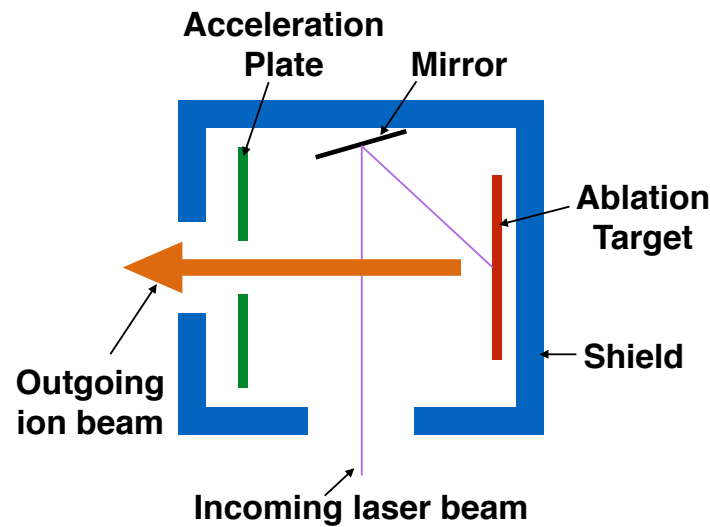
Figure 9.3: Schematic of the beam optics in the experimental area outside the laser lab. The outgoing beam from the periscope is reflected towards the entrance window of the ablation chamber, where it overlaps with the beam from an alignment laser. Adjustable mirrors are used in order to achieve a perfect overlap between the two laser beams. A convex lens is used to focus the beam on the ablation target inside the chamber. Components not drawn to scale. Adapted from [117]



parameters such as wavelength, pulse duration, pulse energy and spot size have a major impact on the ablated particles. For the experiments discussed in this thesis, the laser wavelength is fixed at 355 nm and the pulse duration is 5 ns. The pulse energy is varied however by changing the orientation of the half-wave plate of the optical attenuator described in 9.2, and the spot size is varied by moving the distance of the focusing lens from the laser window.

Ablation target parameters such as choice of material and impurity content also influence the properties of the particles released. Experiments discussed in this thesis use high-purity (> 99%) aluminum and beryllium metal foil as ablation targets. Elemental aluminum has a valency of +3; laser ablation of aluminum can therefore result in the production of Al, Al⁺, Al²⁺, and Al³⁺ ions. Beryllium has a valency of only +2 however, so Be atoms, Be⁺ ions, and Be²⁺ ions are the only possible products from pulsed laser ablation.

Figure 9.4: An illustration of the Ion Source. The components are not drawn to scale. An incoming laser pulse (purple) is reflected by a mirror and hits the ablation target (red). Ions generated are accelerated by the acceleration plate (green), with some of the ions passing through the aperture as a collimated beam (orange arrow). The target, acceleration plate and shield are electrically insulated from each other, and an external power supply can be used to apply a ± 80 V potential bias on each of these components.



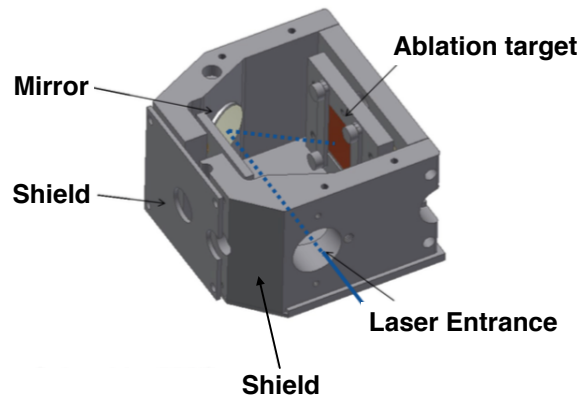
Surface roughness and electrical bias of the target influence the velocity spread of the ions produced. While they were smooth to the touch, the exact surface roughness of the targets used in our experiments was not measured. The targets were electrically connected to a voltage supply however that could provide a potential bias of ± 80 V. By choosing a high positive potential, ions generated could be quickly accelerated away while preventing any electrons from leaving the surface. The choice of potential bias also influences the fraction of singly and doubly charged ions produced.

When the ablation target is positively biased, ions generated quickly fly away from the ablated surface and spread in all directions. While most of the particles simply strike nearby surfaces, a small fraction of particles that are generated perpendicular to the ablation target pass through the aperture in the source and exit as a collimated beam. The collimation is assisted by the acceleration plate shown in Figure 9.4. The acceleration plate is electrically

isolated from both the target as well as the shield. By applying a negative potential on the acceleration plate, ions generated at the target are focused towards the aperture, resulting in a more intense ion beam.

For the design of the ion source, the distance between the target and the acceleration plate (z_{accel}) and the diameter of the aperture (r_{accel}) were optimized by carrying out particle trajectory simulations. The objective was to maximize the fraction of ions that make it through the aperture while minimizing vacuum contamination of the chamber due to accompanying neutral species. The optimal range for z_{accel} was found to be between 10 mm and 50 mm, while r_{accel} anywhere between 3mm and 10mm would work. Consequently, $z_{accel} = 29$ mm and $r_{accel} = 6.5$ mm were chosen for the final design, and the ions source was fabricated with these parameters. Figure 9.5 shows an accurate three-dimensional model of final source. The ablation target, acceleration plate and shield can all be independently biased. The ability to bias the first two will be an important feature for the experiments discussed in this thesis. Unless stated otherwise, the shield was kept at ground potential throughout the measurements.

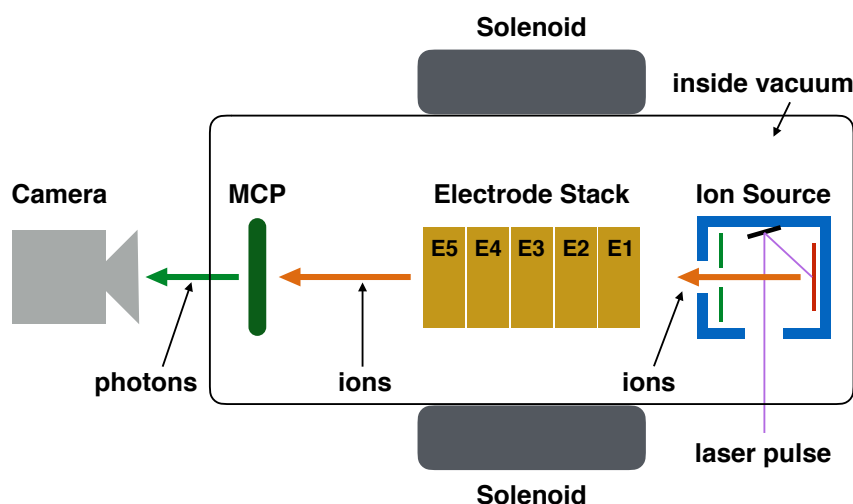
Figure 9.5: An accurate 3D model of the Ion Source, with components drawn to scale. The ablation target (red) could either be made from aluminum or beryllium. The acceleration plate and shield (grey) are made from aluminum. The laser pathway is shown in blue.



9.4 Ablation Chamber

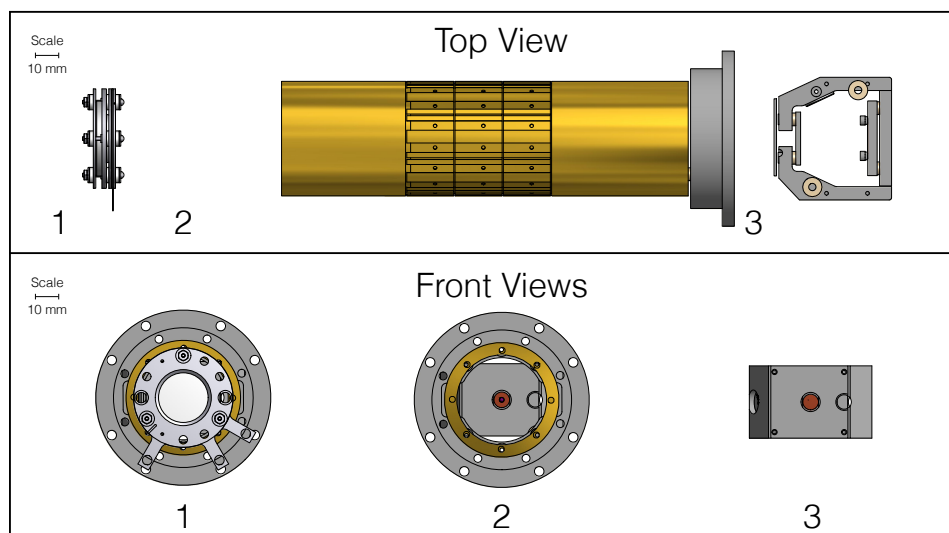
The ion source described in the previous section sits inside a so-called “ablation chamber”, a vacuum chamber where all of the ablation experiments take place. Figure 9.6 shows a schematic of the ablation chamber which is composed of four key components: the ion source, the electrode stack, the solenoid, and the MCP. A 355 nm laser pulse from the ablation laser is used to generate ions as described in Section 9.3. The ions exit the ion source as a collimated beam and travel towards the electrode stack.

Figure 9.6: An illustration of the Ablation Chamber. The components are not drawn to scale. An incoming laser pulse generates ions from the ion source. An electrode stack and a solenoid function as a Penning-Malmberg trap inside the ablation chamber and can be used to block, trap, and/or accelerate the incoming ions. The ions can be dumped to the MCP, which generates photons. A camera situated just outside the vacuum region of the ablation chamber is used to image the photons coming from the MCP.



The electrode stack is composed of five hollow cylindrical electrodes made from gold-plated aluminum, with an inner radius of 22.3 mm and a thickness of 1.5mm (see Figure 9.7 for scaled drawings). Although the choice of electrode material and dimensions is not particularly sensitive, these particular electrodes are used since they closely match the electrodes used in the ALPHA Mixing Trap (see Figure 2.7). Since the ion source is eventually

Figure 9.7: Scaled drawings of the key components of the Ablation Chamber are shown. Front views 1, 2, and 3 show the various apertures in the experimental setup.



intended for installation in the main ALPHA setup, using previously-used ALPHA electrodes reduces another variable when trying to establish the viability of the source in the ablation chamber.

Each electrode is insulated from its neighbor by a set of 1 mm thick ruby balls. The insulation is sufficient as to allow the application of up to ± 400 V on neighboring electrodes without the risk of arcing. Each electrode is connected to an external amplifier that takes an analog input of ± 10 V and outputs ± 400 V. These amplifiers are custom-made in order to have a quick response time ($< 1\mu\text{s}$) and a low background noise level (< 100 mV). The amplifier control lines are integrated into the main computer control system described later, allowing fast application of potentials for each individual electrode.

The solenoid, which provides the axial magnetic field for the Penning-Malmberg trap, is a cylindrical copper-wire electromagnet. The solenoid has a maximum operating temperature of 50 C beyond which effective operation is not guaranteed and damage to the copper wire is possible. For currents

of less than 30 A, the solenoid can be operated continuously without any need for water-cooling, since the copper wire temperature remains below the damage threshold. For currents greater than 30 A however, a water-cooling system is needed to carry heat away from the solenoid. During the course of any experiment, the solenoid temperature is continuously monitored as a safety measure. For results presented in Chapter 10, the maximum operating current in the solenoid is 50 A, and the solenoid temperature is never allowed to exceed 45 C.

9.5 Particle Detectors

A custom-designed MCP manufactured by Photonis is the primary detector used to measure the ion beam generated as a result of pulsed laser ablation. The MCP detector consists of two micro-channel plates and a phosphor screen. The micro-channel plates amplify the signal due to incoming ions by creating a cascade of electrons for each incoming ion. By correctly choosing bias voltages for the MCP front and back, the desired ion-to-electron gain factor can be achieved. The electrons then impact the phosphor screen, which generates 545 nm photons as a result. A CCD camera (manufacturer: PCO, model: Pixelfly) is used to take an integrated image of the photons. The image intensity is therefore a direct measure of the ions arriving at the MCP, and the spatial distribution of the bright pixels is a measure of the radial spread of the ion beam.

While the MCP image provides a direct measure of the transverse distribution and relative intensity of the generated ion beam, it does not provide an absolute measure of the total charge reaching the MCP. To obtain this number, one must operate the front of the MCP as a Faraday cup. The desire to use the MCP both as an amplifier for imaging and a Faraday cup requires the ability to apply a high potential bias in steady-state while at the same time extract the small voltage fluctuation due to the arrival of ions. This is accomplished by means of a low-voltage pickoff circuit that uses a network of resistors and capacitors to separate the high-voltage and low-voltage signals. When charged particles arrive at the MCP in a short time interval, a

momentary surplus of charge creates an additional potential on top of the applied bias. The pickoff circuit decouples this small signal from the applied bias, which can be extracted as a separate output. The magnitude of this signal and the capacitance of the circuit can then be used to calculate the total charge.

The low-voltage pickoff circuit just described can also be used to measure the total number of ions released from the ablation target, which is nominally charged at +40 V. In this case, the ions leaving the target results in a momentary charge deficit instead of a surplus. The pickoff circuit is again able to decouple this ablation signal, providing us with a measure of the amount of charge ablated. The pickoff circuits are therefore an essential component in our ability to detect the number of ions leaving from or arriving at a given surface, and are used extensively for all ablation measurements discussed in this thesis.

9.6 Control System

All of the hardware outlined in this chapter, particularly the ablation laser, the ion source, the electrodes, the solenoid, the MCP, and the camera, must be carefully linked together and controlled in order to reliably perform the ablation experiments. A centralized computer-controlled system, known as the Beryllium Test Sequencer (BTS), is responsible for controlling all of these individual pieces of hardware. By sending and receiving digital and analog signals, the BTS can be used to interact with the ablation apparatus and execute a series of operations with precise amplitude and timing information.

At the lowest level in the control system, the BTS has a simple collection of digital and analog outputs. These outputs define the “state” of the system. A state can last for a pre-defined time, or may continue until an external trigger is received that signals the BTS to move to the next state. Depending on the timing of each state and the choice of input triggers, a collection

of states defines what is known as a “sub-macro”. At the highest level, a user can pre-program a highly customizable “sequence” of sub-macros. By executing the desired sequences, the user effectively controls the parameters and precise timing of an experiment from start to finish, making it possible to make measurements that would not be possible to conduct manually.

At the heart of the BTS hardware is the cRIO-9074 from National Instruments, an embedded controller that features a reconfigurable FPGA module and a 400 MHz real-time processor. The FPGA is programmed to execute pre-defined sequences with < 100 ns jitter. Through the cRIO chassis, the FPGA is connected to independent digital and analog modules. In particular, we use a fast digital input-output module (NI 9402), a slow digital input-output module (NI 9403) and an analog output-only module (NI 9264). The fast and slow digital modules have nominal update rates of 55 ns and $7\mu\text{s}$ respectively, which means that it takes approximately that amount of time for an individual channel to switch from high to low or vice versa. The analog-only module has a latency of $\sim 4\mu\text{s}$.

Most of the instruments requiring a 5V-TTL digital trigger, such as the laser lab safety lamp, the laser shutter, and the MCP camera, are connected to the slow digital module. Only the trapping gate electrode E1 and the data acquisition digital trigger are connected to the fast digital module. Use of the fast digital module is only needed when acquiring time-of-flight data and when attempting to trap ions that have a bunch length of only a few μs , when update rates of less than $1\mu\text{s}$ are required.

The potentials on the electrodes, the current running through the solenoid, and the MCP bias voltages are controlled by the analog output module. The analog output module only provides ± 10 V, so custom-made amplifiers are used in order to allow the individual electrodes to reach potentials as high as ± 400 V. For the solenoid and the MCP, commercial power supplies are used to provide the necessary current and voltage demands.

9.7 Data Acquisition

For every ablation measurement, there are four essential pieces of data that need to be recorded:

1. The fluence of the ablation pulse
2. The number of ions leaving the ablation target
3. The number of ions arriving at the MCP
4. The image of the photons generated by the MCP phosphor screen

The first three pieces of data are recorded on a Tektronix DPO5104 oscilloscope. This fast scope has four independent channels and can record up to 5 GS/s sampling rates. The use of this advanced oscilloscope is important for our measurements since we need to reliably and accurately record events that take place on very different time-scales: the laser pulse has a duration of only 5 ns, while the ions reaching the MCP have typical times-of-flight of tens of microseconds, a difference of nearly four orders of magnitude. During normal operations, the four channels on the scope are the external trigger from the ablation laser, the photodiode output, the pickoff from the target, and the pickoff from the MCP. When the scope is triggered, it records the signal on all four channels. The trace of each channel along with a timestamp is saved into CERN's network drive for subsequent analysis.

As mentioned previously in Section 9.1, the ablation laser is operated in Auto mode in order to minimize the variation in pulse energy from one shot to the next. This means that every ablation experiment has to be synchronized with the clock of the ablation laser that pulses at a regular interval of 50 ms. When an ablation measurement is desired, the laser shutter is opened, allowing a single pulse to pass through before the shutter is quickly closed. Generation of the laser pulse coincides with a digital trigger sent by the laser to the scope to start recording all subsequent events for the next $200\mu\text{s}$ in 1 ns time-steps.

The calibrated photodiode is used to measure the energy of each ablation pulse. By knowing the position of the focusing lens relative to the ablation target, the precise laser spot size can be calculated. The laser fluence is then simply obtained by dividing the pulse energy with the area of the laser spot.

For most of the ablation experiments, the target is positively biased at 40 V in order to prevent the release of photoelectrons and promote the creation of ions. As the ions quickly escape the target, the deficit of positive charge momentarily creates a negative potential with respect to the nominal bias voltage, on the order of a few tens to hundreds of mV. A low-voltage pick-off circuit is connected to the target in order to separate this signal due to ablation from the background bias potential. By knowing the capacitance of the circuit and the charge of the ablated species, the total number of ions leaving the target can be calculated.

A low-voltage pickoff circuit is also used to calculate the number of ions arriving at the MCP. The front of the MCP is biased at -600 V. As ions reach the surface of the MCP, the accumulation of charge shows up as an additional potential on top of the bias from the power supply. The pickoff circuit extracts this signal which can then be used to calculate the total number of ions received at the MCP. Unfortunately, due to significant electronic noise associated with the pickoff circuit and amplifier, a signal is only observed if more than 10^8 ions arrive at the MCP.

The final piece of data that is stored with every ablation pulse is the MCP image. As the ions hit the MCP phosphor screen, the photons generated make their way through the vacuum window and into the MCP camera. The camera, which is triggered at the same time as the scope, has pre-defined exposure duration of 5 ms. At the end of the exposure time, an axially-integrated two-dimensional image is obtained which is saved into CERN's network drive along with a timestamp for further study.

Chapter 10

Ion Generation and Trapping

The ablation apparatus and associated hardware and software described in Chapter 8 are designed to test, implement and qualify the functionality of the ion source, specifically for the production of beryllium ions that could be used for the sympathetic cooling of positrons in ALPHA. When the apparatus was first built, aluminum, instead of beryllium, was used as the ablation target in the source. There were two reasons for this choice of metal:

1. The availability of literature on the pulsed laser ablation of beryllium is very sparse [104]. In comparison, significantly more literature is available for metals such as aluminum [104, 105, 108, 111, 112, 118, 119], copper [107, 109, 110, 112, 119] and iron [104–106, 108]. In fact, many of these references study two or more of the aforementioned metals in parallel, allowing quantitative comparisons to be made between them. By investigating aluminum in our ablation chamber, we can benchmark our apparatus against these previous measurements. Once the behavior of aluminum in our setup is understood, we can then proceed towards the pulsed laser ablation of beryllium in a more informed manner.
2. Beryllium vapor in large amounts can be toxic and is a potential health hazard if inhaled. Attempting experiments with beryllium without a complete understanding of its ablation threshold and ion yield presents a risk of contaminating the internal components of the ablation chamber. If the contamination is significant, the exposed components may be rendered unusable for future experiments where beryllium contam-

ination is problematic.

It was therefore considered best to begin with aluminum and obtain a thorough understanding of the ablation process in our apparatus before switching to a beryllium target. In this chapter, I discuss the ablation measurements carried out with both aluminum and beryllium. I also compare the experimental results with particle trajectory simulations and extract relationships between laser fluence, ablation threshold, ion yield, and kinetic energy distribution of the ablated ions. Finally, I discuss attempts to trap beryllium ions and estimate the trapping efficiency in one specific case.

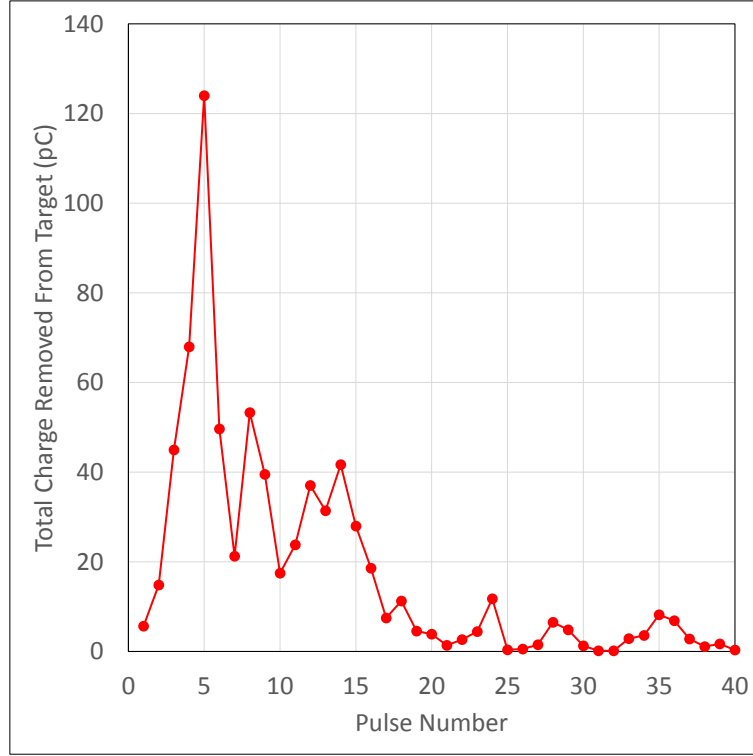
10.1 Laser Ablation of Aluminum

Pulsed laser ablation in our apparatus was first carried out on an aluminum target, a thin rectangular piece $15\text{mm} \times 13\text{mm}$ of at least 99.5% purity purchased from CERN Stores in Switzerland. After securely installing the target into the ion source, the source was inserted into the ablation chamber. The chamber was promptly vacuum-sealed and pumped down to pressures of $< 1 \times 10^{-7}$ mbar before ablation measurements were made.

10.1.1 Cleaning Pulses

During the first ablation experiments with aluminum in the ablation chamber, the chamber was still being actively pumped by a turbomolecular pump, with the pressure hovering around the 2×10^{-6} mbar level. Ablation was carried out by sending single pulses from the 355 nm laser to the target and recording the amount of charge removed as a result. It was observed that the amount of material ablated from the target varied significantly from shot-to-shot, and in particular the total charge removed for the first ~ 20 pulses was significantly higher than for any subsequent pulses. Figure 10.1 shows this effect for one of the measurement days.

Figure 10.1: Plot of the total charge removed from the aluminum target (ion yield) for 40 consecutive ablation pulses with approximately $62\mu\text{J}$ pulse energy. The first few pulses give a significantly higher ion yield than pulses that arrive later in the sequence.



The amount of charge ablated during the first 20 pulses (mean = 32.3 ± 6.3 pC) is an order of magnitude higher than the total charge ablated during the subsequent 20 pulses (mean = 3.2 ± 0.7 pC). The total shot-to-shot variation in the ablated charge for pulses 21 to 40 is also smaller. These observations can be explained by noting that at the prevalent pressure of this measurement ($\sim 2 \times 10^{-6}$ mbar), it takes only ~ 1.6 seconds for the formation of a monolayer of background gas particles (mainly oxygen) on the aluminum surface [120]. The first few pulses can therefore be understood to be acting as “cleaning” pulses before the bare aluminum surface becomes exposed to the laser pulses. Higher ion yields due to ablation of “unclean” aluminum surfaces have also been reported previously [121] with cleaning pulses similarly being used to effectively mitigate the issue.

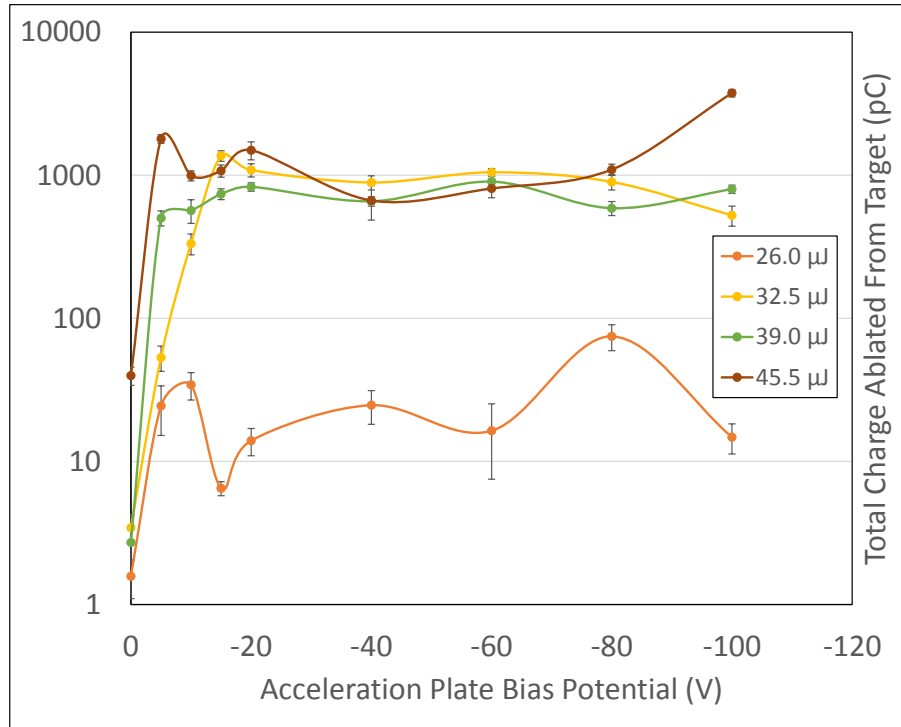
With this in mind, it was decided to carry out ablation with nominally 40 cleaning pulses before any measurement was carried out. Eventually, after the installation of an ion pump, the ablation chamber was kept at pressures $< 5 \times 10^{-8}$ and cleaning pulses were no longer necessary (> 60 s monolayer formation time [120]). Unless otherwise stated, all measurements discussed in this thesis were carried out at $< 5 \times 10^{-8}$ mbar and without any cleaning pulses.

10.1.2 Target and Acceleration Plate Bias Potentials

For our measurements, the first task was to establish the optimum bias potentials for the ablation target and the acceleration plate. Figure 10.2 shows the amount of charge removed from the target as a result of ablation at four different pulse energies. The acceleration plate bias was scanned between 0 V and -100 V, while the target potential was kept constant at $+40$ V. When the acceleration plate is not biased i.e. when it is held at 0 V, the amount of charge ablated from the target is at least an order of magnitude lower than when the acceleration plate is negatively biased. In fact, with an acceleration plate bias of just -10 V, the target ablation response appears to saturate, with higher bias potentials resulting in little or no increase in the total ablated charge.

This effect can be explained by looking at the laser ablation products. When the laser pulse arrives at the target, the ablation process results in the creation of positive ions and electrons both within the bulk of the target as well as at the surface. Depending on the target material as well as the energy and spot size of the laser pulse, a significant fraction of the products may be low-energy (thermal) ions. Ions and electrons at the surface with sufficient axial kinetic energy escape the target. When more ions leave than electrons, a momentary negative potential is created. This negative potential prevents ions with insufficient energy from escaping the target. These low-energy ions quickly recombine with electrons within the target to form neutral atoms and therefore do not contribute to the overall ablation signal.

Figure 10.2: The total charge removed from the aluminum target as a function of the applied bias potential on the acceleration plate. The application of a bias potential appears to dramatically increase the ion yield, which quickly saturates for potentials of magnitude -10 to -20 V. The vertical error bars, plotted on a log scale, are standard errors of the mean.



When an accelerating electric field is present however (e.g. by applying a negative potential at the acceleration plate), some of the thermal ions are able to overcome the potential barrier and escape, and fewer electrons leave the target surface. This is manifested as an increase in total positive charge removed in Figure 10.2 when only -5 V of bias is applied. For higher biases, this effect is amplified until one reaches the saturation point when all of the ions produced by ablation make it out of the target and almost all of the electrons remain confined. In Figure 10.2, this lies somewhere between -10 V and -20 V beyond which no significant increase is seen in the total charge ablated.

Based on the results from Figure 10.2, a bias of magnitude -10 V to -20

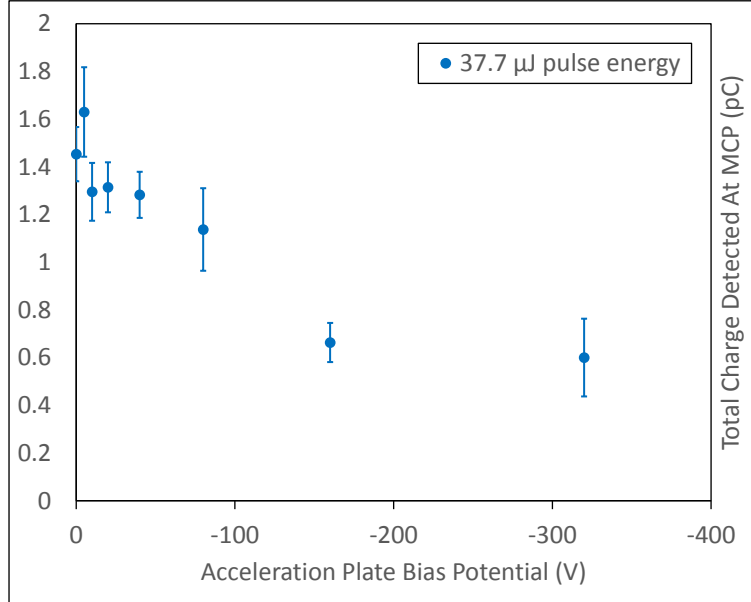
V would be sufficient for the acceleration plate. However, the impact of the specific value chosen on the ion transfer efficiency also needs to be considered, where we define the ion transfer efficiency as the fraction of ions produced at the target that successfully leave the aperture of the ion source and make it into the Penning-Malmberg trap. In order to estimate the ion transfer efficiency, we can look at the number of ions reaching the MCP when all the electrodes of the trap are grounded.

Figure 10.3 shows the amount of ablated charge detected at the MCP as the acceleration plate bias is scanned between 0 V and -320 V. Again, the target potential was kept constant at $+40$ V and the laser pulse energy was $37.7 \pm 3.8 \mu\text{J}$. We observe that the amount of charge detected at the MCP is maximum when the bias is -5 V. As the bias magnitude is increased, the amount of charge reaching the MCP decreases. Presumably fewer ions are able to pass through the aperture of the ion source when biases of higher magnitude are applied on the acceleration plate. This may be due to an increase in the transverse spread of the ion beam for larger acceleration potentials, resulting in ions missing the aperture and hitting the acceleration plate instead.

Having a high transfer efficiency of low energy ions from the ion source is crucial to eventually having a large number of trapped ions in the Penning-Malmberg trap. It is therefore preferable to keep the acceleration plate potential as low as possible. For this reason, an acceleration plate negative bias potential of magnitude 10 V is selected and is used for all subsequent measurements unless stated otherwise.

During the previous two measurements, the target bias was kept constant at $+40$ V. Figure 10.4 shows the effect on the total charge reaching the MCP for target bias potentials between 10 V and 320 V. For this measurement, the acceleration plate was held at -20 V and the laser pulse energy was $37.7 \pm 3.8 \mu\text{J}$. For a target bias between 0 V and 80 V, the effect is unclear since the data points do not exhibit any observable trend. Between 80 V and 320 V however, there is a clear increase in the charge arriving at the MCP for higher bias potentials.

Figure 10.3: The amount of positive charge detected at the MCP as a function of the applied bias potential on the acceleration plate. Larger biases appear to reduce the total charge that makes it to the MCP. The error bars are standard errors of the mean.

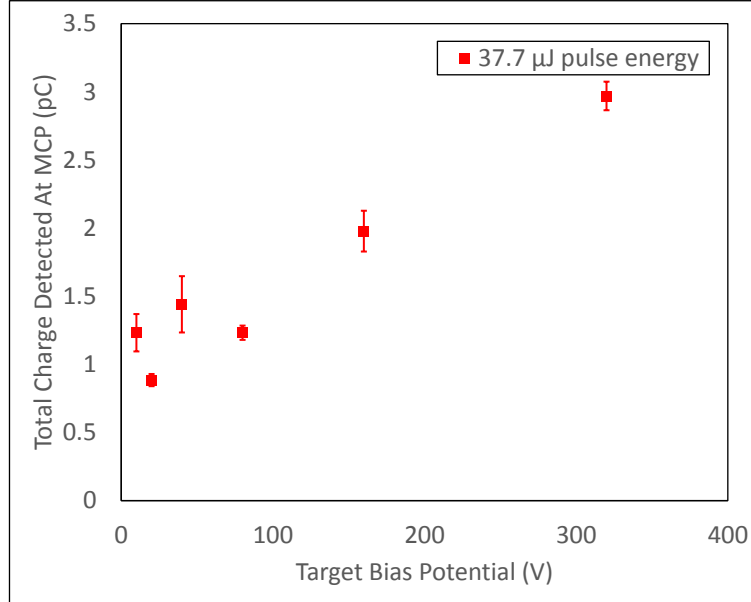


For ablation at a given fluence, increasing the target bias potential is expected to increase the axial kinetic energy of the ablated ions. This results in more ions with trajectories within the solid angle of the aperture, resulting in an increased transfer efficiency into the Penning trap and consequently an increased number of ions reaching the MCP as seen in Figure 10.4. Excessively high kinetic energies are however inadvisable for experiments intended to dynamically trap and subsequently laser-cool ions in a Penning trap. For this reason, we decide to continue with our somewhat arbitrary choice of +40 V as the target bias potential for all future experiments.

10.1.3 Al^+ Ablation Threshold and Ion Yield

In order to determine the minimum fluence required to produce singly charged aluminum ions, single pulses between $10\mu\text{J}$ and $30\mu\text{J}$ were used to

Figure 10.4: The amount of positive charge detected at the MCP as a function of the applied bias potential on the aluminum target. Larger biases appear to increase the total charge that makes it to the MCP. The error bars are standard errors of the mean.

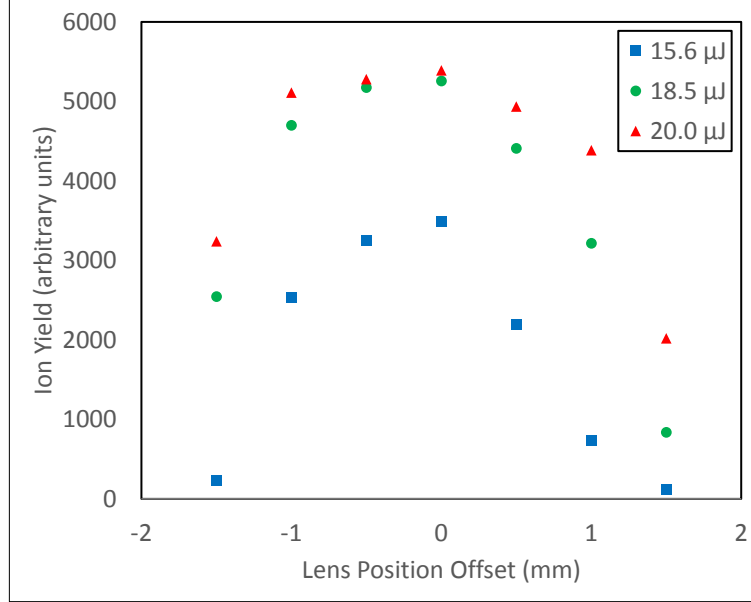


ablate the aluminum target. The target was biased at +40 V, the acceleration plate at -10 V and the shield at 0 V. The position of the focusing lens was also scanned in order to determine the location where the maximum number of ions are produced (see Figure 10.5). This corresponds to the location where the laser is exactly focused on the target; at the focus, the laser has a beam waist of $16.6 \pm 2.5 \mu\text{m}$ and a corresponding spot area of $(8.7 \pm 2.6) \times 10^{-6} \text{ cm}^2$ (beam waist measurement details in Reference [117]).

Figure 10.6 shows how the production of Al^+ ions varies with increasing laser fluence. For the lowest pulse energy of $10 \mu\text{J}$ used in this experiment, the ablation signal is already above background, with a yield of $\sim 0.2 \times 10^9$ ions. This implies that the ablation threshold for aluminum lies below 1.2 Jcm^{-2} , consistent with previously reported value of 0.8 Jcm^{-2} in [111] for similar laser parameters.

For fluences beyond 1.5 Jcm^{-2} , the ion yield increases dramatically, rising

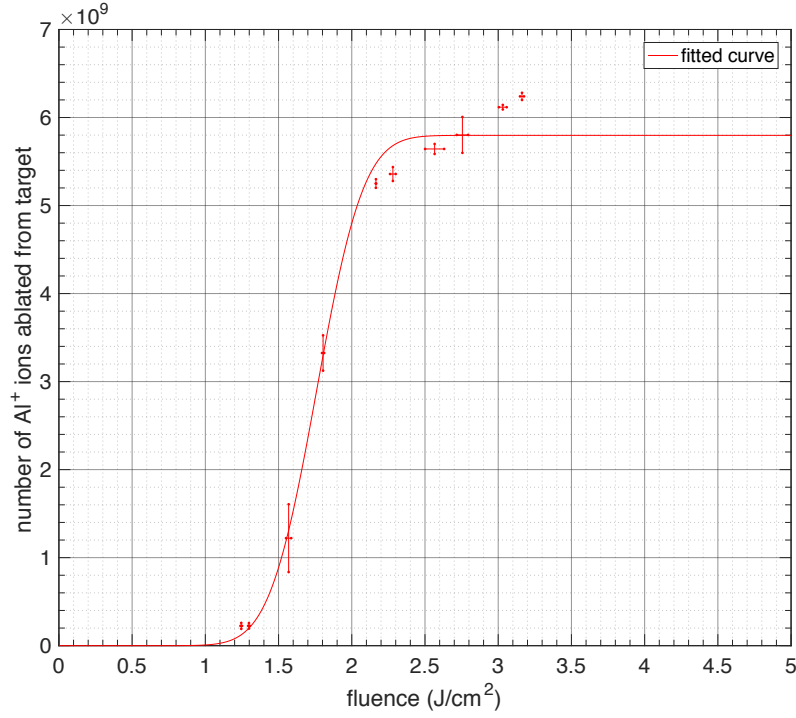
Figure 10.5: Variation in ion yield (at the target) as a function of lens offset from the focus for three different pulse energies. At the focus, the $15.6\mu\text{J}$, $18.5\mu\text{J}$, and $20.0\mu\text{J}$ pulse energies correspond to fluences of 1.8 Jcm^{-2} , 2.1 Jcm^{-2} , and 2.3 Jcm^{-2} respectively.



from just 1×10^9 at 1.5 Jcm^{-2} to nearly 5×10^9 at 2.0 Jcm^{-2} . For yields beyond $\sim 5 \times 10^9$ ions, the electronics used to pick off the ablation signal from the target begin to saturate. This is due to the inclusion of protection diodes (manufacturer: Vishay, model: 1N4937) in the low-voltage pickoff circuit. These diodes require a minimum forward bias of 1.2 V. Therefore, if the number of ablated ions is large enough to create a momentary potential difference of $> 1.2\text{ V}$ with respect to the applied target bias, the diodes begin conducting current, producing a saturation effect. In Figure 10.6, this is represented by the slow leveling-off of the signal for fluences of 2.2 Jcm^{-2} and larger. The saturation of the signal also allows us to fit a curve of the form $A \left(1 + \text{erf} \left(\frac{x-\mu}{\sigma\sqrt{2}} \right) \right)$. This curve crosses the 1×10^9 ion yield mark at a fluence of 1.5 Jcm^{-2} .

In Figure 10.7, we plot the same data along with four different exponential curves that pass through the $(x = 1.6, y = 1.2 \times 10^9)$ data point. Although the data is sparse, it is consistent with the observation that beyond the abla-

Figure 10.6: A plot of the number of ablated aluminum ions against laser fluence. The error bars at each data point are standard errors of the mean. A fitted curve of the form $A \left(1 + \operatorname{erf} \left(\frac{x-\mu}{\sigma\sqrt{2}} \right) \right)$ is also shown. The data suggests an Al^+ ablation threshold around 1 Jcm^{-2} .

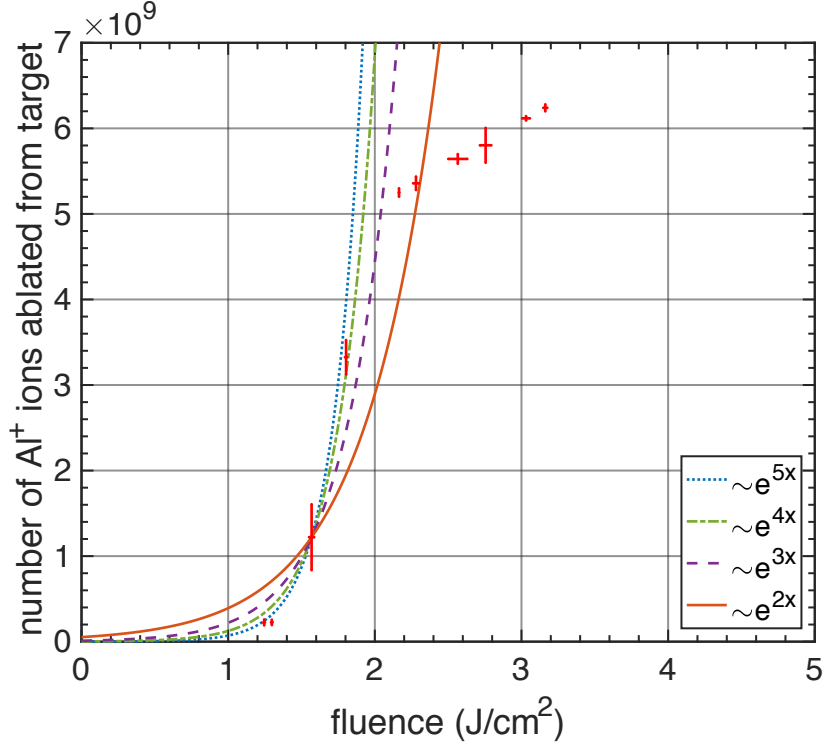


tion threshold, the number of ablated ions rises exponentially with increasing fluence. Similar sharp ion yield increases with fluence have previously been reported in [118]. The exponential increase in yield is attributed to an increase in the total number of ablated material coupled with more efficient multi-photon ionization of the ablated particles within the plume [118].

10.1.4 Al^+ Kinetic Energy

A large fraction of the ablated aluminum ions do not escape the ion source but instead only reach as far as the acceleration plate or shield where they implant themselves into the blocking surface. Only a small fraction of ions - the ones ablated normal to the surface - pass through the aperture opposite

Figure 10.7: A plot of the number of ablated aluminum ions against laser fluence. The error bars at each data point are standard errors of the mean. Theoretical exponential curves proportional to e^{2x} , e^{3x} , e^{4x} and e^{5x} are also shown. The ion yield is roughly consistent with an exponential rise with laser fluence.

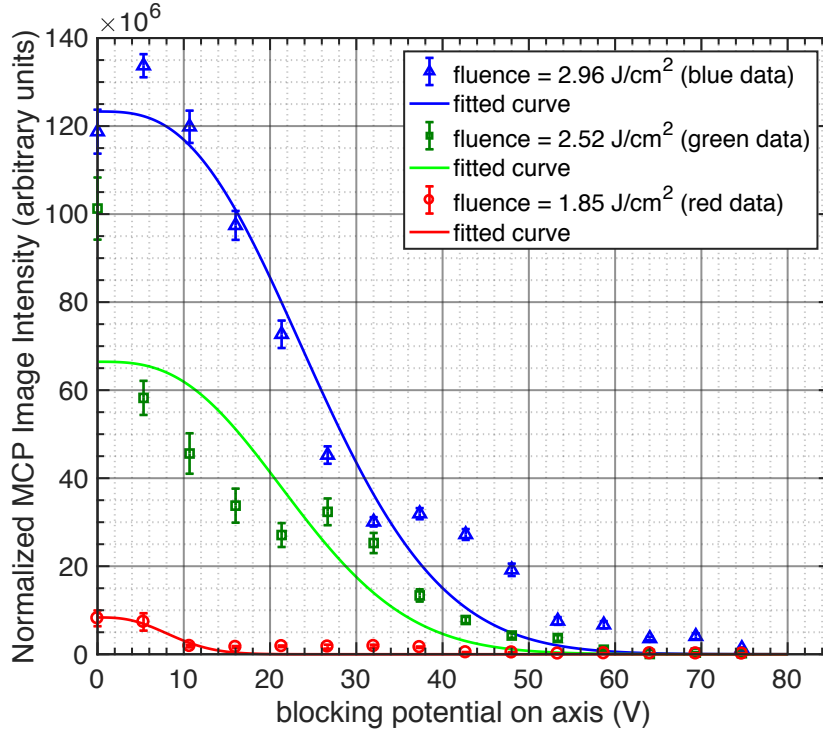


the target. In one experiment with 40 ablation pulses, we obtain this fraction to be $(1.4 \pm 0.2) \times 10^{-2}$. The small aperture guarantees that any ions exiting the ion source will remain unimpeded until they reach the electrode stack (the ion source aperture subtends a solid angle of 0.04 sr at the target, while the solid angle for the opening of the electrode stack from the target is 0.26 sr). The five electrodes of the stack along with the copper solenoid form a Penning-Malmberg trap. Ions reaching the electrodes may therefore be blocked, trapped, accelerated, or simply allowed to pass through unhindered.

One method of obtaining the kinetic energy distribution of the ions from the source is by selectively blocking ions below a certain kinetic energy while allowing more energetic ions to pass through and hit the MCP. The selective blocking can be achieved by applying a positive potential on one of the

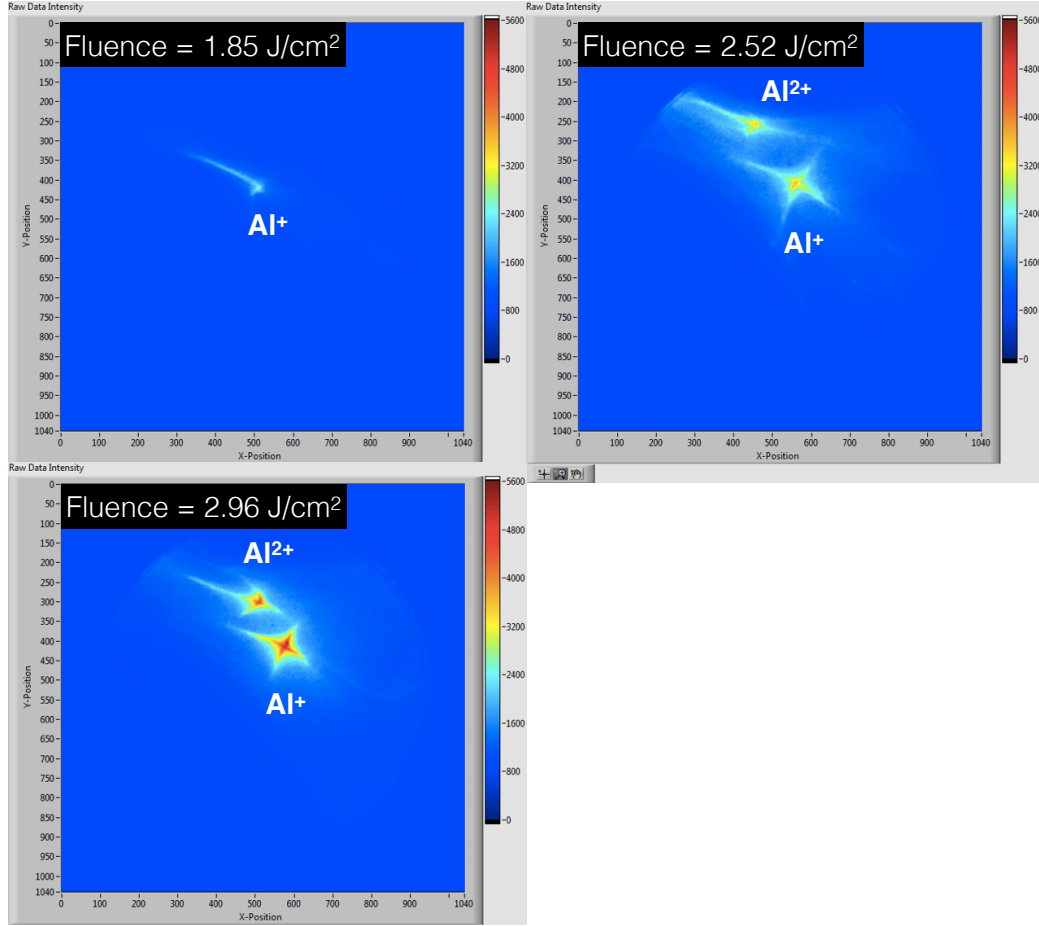
electrodes. By measuring the MCP signal for various blocking potentials, a cumulative distribution is obtained. Figure 10.8 shows such cumulative distribution curves at three different fluence values.

Figure 10.8: A plot of the the MCP image intensity of aluminum ions vs electrode potentials for three different laser fluences. The background MCP image intensity is subtracted from each data point. The error bars shown are standard errors of the mean. A Maxwell-Boltzmann cumulative distribution function of the form $A \left(1 + \sqrt{\frac{2}{\pi}} \frac{x e^{-\frac{x^2}{2a^2}}}{a} - \text{erf} \left(\frac{x}{\sqrt{2}a} \right) \right)$ is fitted to the data. For blocking potentials > 30 V, more than half of the ions are blocked for all fluences.



The total number of ions reaching the MCP is higher for higher fluences, indicative of the underlying higher ion yield. More importantly however, for the fluences used in this experiment, most of the ions generated have low energy, signified by the fact that more than half of the ions are blocked by potentials as low as 30 V. The presence of a high fraction of low energy ions is promising for our eventual goal of trapping and laser-cooling.

Figure 10.9: Typical MCP images for ablation carried out on an aluminum target at three different fluences.



The ion kinetic energy distribution is complicated however. The fitted curve shown, a Maxwell-Boltzmann cumulative distribution function of the form $A \left(1 + \sqrt{\frac{2}{\pi}} \frac{x e^{-\frac{x^2}{2a^2}}}{a} - \operatorname{erf} \left(\frac{x}{\sqrt{2}a} \right) \right)$, is a poor fit to the data: it overestimates the number of low-energy ions and underestimates the number of high energy ions. Assuming that the ablated ions in fact follow a single Maxwell-Boltzmann distribution, one potential explanation for this anomaly is the existence of doubly-charged ions, Al²⁺.

Figure 10.9 shows typical MCP images for ablation at the three fluences (and no blocking potential). For the 2.52 Jcm⁻² and 2.96 Jcm⁻² laser fluences, images from doubly-charged aluminum ions are clearly visible, suggesting an

ablation threshold for Al^{2+} of $< 2.5 \text{ Jcm}^{-2}$. This is roughly consistent with the previously reported threshold of $\sim 3 \text{ Jcm}^{-2}$ in Reference [118], but not the lower threshold of 1.4 Jcm^{-2} reported in Reference [111]. The observed cumulative data may therefore incorporate two independent Maxwell-Boltzmann distributions, one for each of the charged species. Further measurements with doubly-charged ions are needed in order to verify this hypothesis.

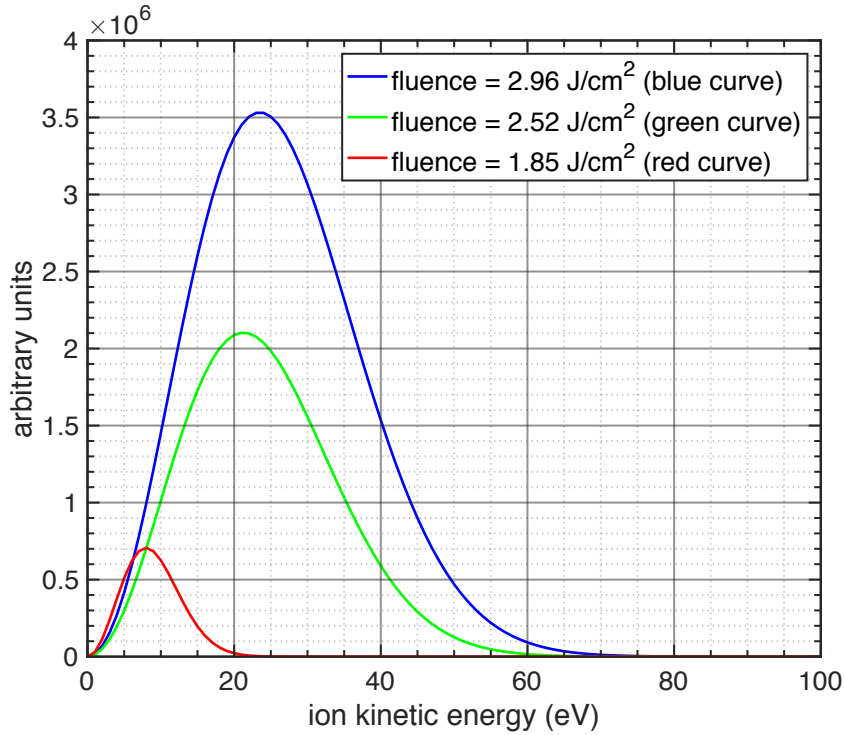
An alternative hypothesis is that the cumulative data represents a sum of several independent Maxwell-Boltzmann distributions that originate due to two separate ion formation processes. In Reference [111], aluminum ablation for fluences $> 1.3 \text{ Jcm}^{-2}$ is seen to produce Al^+ kinetic energy distributions with two separate Maxwell-Boltzmann components: a slower component due to multiphoton ionization of the ablation plume, and a faster component due to electron-impact ionization of neutrals within the target volume. For our data shown in Figure 10.8, the total distribution is most likely a convolution of both separate ion formation processes as well as the presence of both singly and doubly-charged species. Unfortunately the resolution of our detection mechanism is insufficient to deconvolve these two effects.

Using the cumulative distribution fits to the data collected above, and assuming that the distribution is dominated by Al^+ originating from a single formation process, we can plot Maxwell-Boltzmann distributions of the ion kinetic energies, as shown in Figure 10.10. The distributions are scaled relative to each other according to the scale factor A obtained from the fits. A lot fewer ions are produced at a laser fluence of 1.85 Jcm^{-2} compared to the two higher fluences. Furthermore, the average kinetic energy of the ions is seen to increase for higher fluences, as expected. The peak kinetic energies are 8 eV, 21 eV and 24 eV at 1.85 Jcm^{-2} , 2.52 Jcm^{-2} , and 2.96 Jcm^{-2} respectively. These kinetic energies are suitable for preliminary trapping experiments in the ablation chamber.

Our analysis relies on the assumption that the fitted curves obtained in Figure 10.8 accurately represent the kinetic energy distribution of singly-charged aluminum ions. Given the poor fit to the data at higher fluences, we unfor-

Unfortunately cannot trust the quantitative aspect of the analysis. We may comment on the overall qualitative behavior however, and note that the results are consistent with expectations that higher fluences result in the production of more energetic ions from the laser-ablation process.

Figure 10.10: A plot of the kinetic energy distribution of the ablated aluminum ions for three different laser fluences. The curves are based on the Maxwell-Boltzmann fits in Figure 10.8 and are scaled to reflect the relative ion production at each fluence. At a fluence of 1.85 Jcm^{-2} , almost all of the ions have axial kinetic energies less than 20 eV. This is not the case for the two higher fluences however, where more than two-thirds of the ablated ions have energies above 20 eV.



10.1.5 Vacuum Contamination from Ablation

An important consideration in the use of the ion source is the total amount of material ablated per pulse, and how that may impact the overall vacuum in the ALPHA trap. Minimizing background gas in the trap is essential for a long lifetime of antiprotons and positrons in order to create antihydrogen.

For the sympathetic cooling of positrons, we require approximately 5×10^5 trapped and laser-cooled Be^+ ions. Based on particle trajectory simulations of 10 eV thermal ions, we can estimate an ion source production efficiency of 1%, i.e. for every 100 ions ablated from the target, only 1 ion makes it through the aperture and into the electrode stack. The efficiency is much higher if ions are preferentially produced in the forward direction (i.e. they have an excess axial kinetic energy). When the trap is long enough relative to the ion beam bunch length, we can expect to have nearly a 100% trapping efficiency: this is expected to be the case in the main ALPHA Atom Trap where the trapping efficiency of antiprotons from the Catching Trap into the Atom Trap is nearly 100%. The trapping efficiency is much reduced if the bunch length exceeds the length of the trap (which is expected to be true in the Penning trap of the ablation chamber). Finally, we can pessimistically assume a laser-cooling efficiency of at least 5% i.e. for every 100 Be^+ ions loaded into the ALPHA Atom Trap, at least 5 are laser-cooled to below 5 K to eventually be used for the sympathetic cooling of positrons.

Multiplying all of these factors results in the requirement of producing 10^9 ions in order to have at least 5×10^5 trapped and laser-cooled ions in the ALPHA Atom Trap.

In the aluminum ablation experiments discussed in Section 10.1.3, a laser fluence of 3 Jcm^{-2} produces well over 10^9 Al^+ ions without any observable change in the vacuum chamber pressure ($\sim 5 \times 10^{-8}$ mbar background pressure). Later, in Section 10.2.1, we will note that in the case of beryllium, laser fluences as high as 5 Jcm^{-2} were attempted, again with no observable vacuum effects at the aforementioned background pressure, and yields well beyond the 10^9 requirement. Increases in the vacuum chamber pressure are only observed when ablation pulses have energies in excess of $200 \mu\text{J}$ (corresponding to fluences $> 23 \text{ Jcm}^{-2}$ at the focus). With such high pulse energies, we observe occasional pressure spikes up to a maximum of 2×10^{-7} mbar, which decay back to the background level in less than 10 seconds.

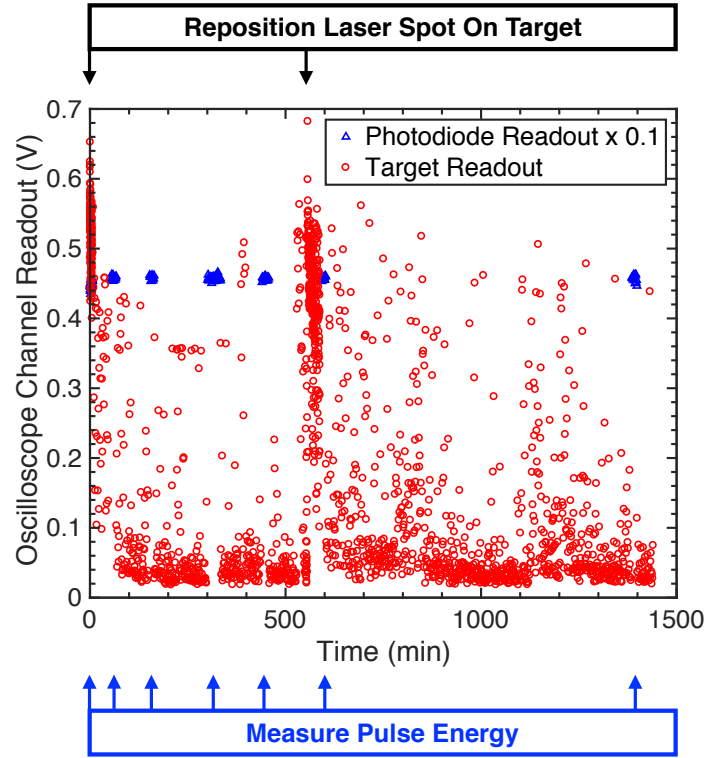
Laser ablation of metals usually results in the release of both ions and neu-

tral atoms. The ion fraction i.e. the ratio of ions to atoms from ablation depends on the laser pulse wavelength as well as the fluence used. Ablation with UV lasers produces a higher ion fraction compared with IR wavelengths [109]. Similarly, ablation at higher fluences produces a higher relative ion count than when ablating with lower fluences [113].

In order to rule out any vacuum effects from ablated neutral species, an estimate of the ablation ion-to-atom ratio needs to be determined. An experiment was carried out in the ablation chamber where a total of ~ 5.46 million pulses at $\sim 500\mu\text{J}$ of energy were used to ablate a freshly installed aluminum target over the course of 5 days. On each day, the laser would first be warmed up, aligned (using an alignment laser) and focused onto the ablation target. Next, the laser pulse energy would be measured and then ablation of the target would commence for an extended period of time at 20 pulses per second. The amount of charge removed due to an ablation pulse would be recorded each second for the first 5 minutes, and then every minute until the end of the measurement. The photodiode signal would also be periodically read out to ensure that the pulse energy delivered to the target was not drifting from the nominal value of $500\mu\text{J}$.

Figure 10.11 plots the oscilloscope readout for the third day of the measurement. At time $t = 0$ minutes, the laser pulse energy is measured and the laser is aligned and focused onto the ablation target. For the first ~ 5 minutes, the ablation target registers a signal of 0.5 ± 0.1 V, corresponding to the removal of ~ 540 pC of charge per pulse. In the subsequent 60 minutes, the magnitude of the charge removed begins to decline, and most of the readings lie scattered between 0.1 V and 0.4 V. After the 70-minute mark, hardly any charge is removed from the target with the target readout usually reading < 0.1 V. These low readings continue until the next instance of repositioning the laser spot on the target around $t = 550$ minutes. When the laser is repositioned at the ~ 550 -minute mark, the large 0.5 V ablation signal reappears, only to again decline during the subsequent 1 hour and remain low for the rest of the measurement.

Figure 10.11: Readout of the ablation signal as a result of ablating an aluminum target at a rate of 20 pulses per second in a 24-hour period. The black arrows above the plot indicate the times at which the laser spot on the target was repositioned to a fresh (previously un-ablated) location. The blue arrows below the plot indicate the times at which the photodiode signal was read out for several consecutive pulses in order to measure the pulse energy.



The observed behavior can be attributed to the creation of craters on the target. Initially, when the laser pulses arrive at a new location on the target surface, the ablation process results in the removal of atoms and ions. Over time, the repeated removal of ablated material results in the creation of a pit or crater. As the crater deepens, it becomes less and less likely for any ablated material to leave the target surface due to the reduction in solid angle.

During the 5-day period of this measurement, the removal of ions for a total of 6554 pulses was recorded with an integrated $1.13\mu\text{C}$ of charge removed. For 5.46 million pulses, this scales to 938 mC or 5.86×10^{15} singly-charged ions. The target was weighed both before and after the experiment, and the

decrease in mass was found to be within the precision of the scale (0.1 mg). Taking this value as an upper bound gives us a maximum of 2.23×10^{18} ablated particles. The ratio of the total number of singly-charged ions created to the maximum number of ablated particles gives us the lower bound on ion fraction at $\sim 0.25\%$. This estimate is consistent with a previous study of aluminum ablation with a UV laser where the reported ion fraction was 1.7% [108].

Due to the small aperture of the ions source, most of the ablated neutral species would simply end up on the internal surfaces of the ion source. For a 10 eV thermal distribution of ablated atoms, we expect at most 0.1% to exit the source and reach the electrode stack. This translates into a worst case scenario of 10 atoms for each ion loaded into the Penning-Malmberg trap. In the ALPHA experiment, this level of particle contamination is negligible and is not expected to adversely effect the antiproton or antihydrogen lifetime in the trap.

In addition to the immediate vacuum degradation that might be caused as a result of pulsed laser ablation, we also need to look into the health and safety aspects of producing beryllium vapor as a result of ablating a beryllium target. The Occupational Safety and Health Administration (OSHA) permits an average exposure of up to $2\mu\text{g m}^{-3}$ in an eight-hour working day, and a peak 30-minute exposure to not exceed $25\mu\text{g m}^{-3}$. The American Conference of Governmental Industrial Hygienists (ACGIH) suggests an even more conservative value of $0.05\mu\text{g m}^{-3}$. Exposure to beryllium vapor at concentrations below this threshold limit value poses no risk to an individual repeatedly working in such an environment.

Assuming the calculated production ion fraction and trapping efficiencies also hold true for beryllium, the desire to have 5×10^5 laser-cooled ions in the Atom Trap equates to producing at most 10^{11} ablated beryllium particles (atoms and ions) with every single pulse. This corresponds to an upper bound of 1.5×10^{-15} kg of beryllium vapor in each pulse or an equivalent of ~ 1.1 ng per 24-hour period assuming ablation is carried out every 2 minutes

(length of AD cycle). Given the large volume of the ALPHA experimental area and the air ventilation system of the AD hall, the production of beryllium vapor at 1.1 ng per day poses no health risk whatsoever to members of personnel as well as any visitors to the site.

10.2 Laser Ablation of Beryllium

Having measured the yields and kinetic energy distributions of ablated aluminum ions, and having estimated the ion fraction of ablated material per pulse, we have a better understanding of how ablation in our apparatus works. This allows us to proceed to our next target of interest: beryllium. For the next series of measurements, the ablation chamber was aired up to atmosphere, the aluminum target was replaced with a beryllium one, and then the chamber was pumped down to pressures of $< 5 \times 10^{-8}$ mbar. The target is again a thin rectangular piece of metal, in this case at least 99.8% purity beryllium purchased from Materion Electrofusion USA.

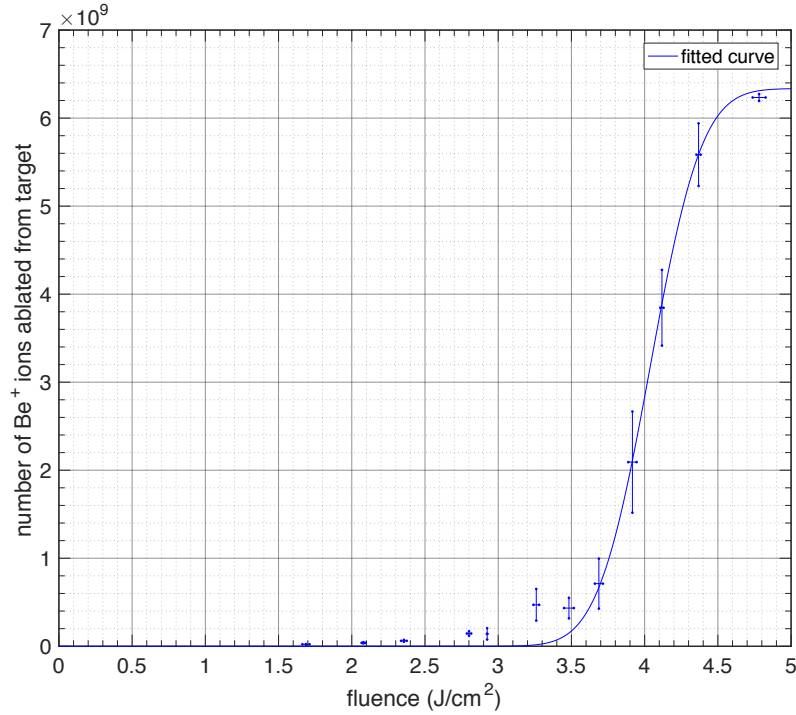
10.2.1 Be^+ Ablation Threshold and Ion Yield

As done previously with aluminum, single pulses of 355 nm light were focused onto the beryllium target and the resulting release of ions was measured and recorded by the data acquisition system. The target and acceleration plate were kept at +40 V and -10 V respectively. The fluence was varied between 1.5 Jcm^{-2} and 5 Jcm^{-2} , and the resulting ion yields were plotted on a graph, as shown in Figure 10.12.

Below a laser fluence of 2.0 Jcm^{-2} , the signals observed are consistent with noise in the readout. Close to $2.0 - 2.4 \text{ Jcm}^{-2}$, the signal is mostly indistinguishable from background, but occasionally a clear reading is observed. Beyond 2.5 Jcm^{-2} , the ablated ions are clearly visible in the data. These measurements suggest that in the case of beryllium, the ablation threshold for singly-charged ions lies just above 2 Jcm^{-2} , a factor of two larger than the

estimate for Al^+ . Given that the melting point of beryllium (1560 K) is a factor of ~ 1.67 larger than that of aluminum, the higher ablation threshold for beryllium is qualitatively consistent with earlier reports where metals with high melting points are observed to have relatively higher ablation thresholds for singly-charged ions [112].

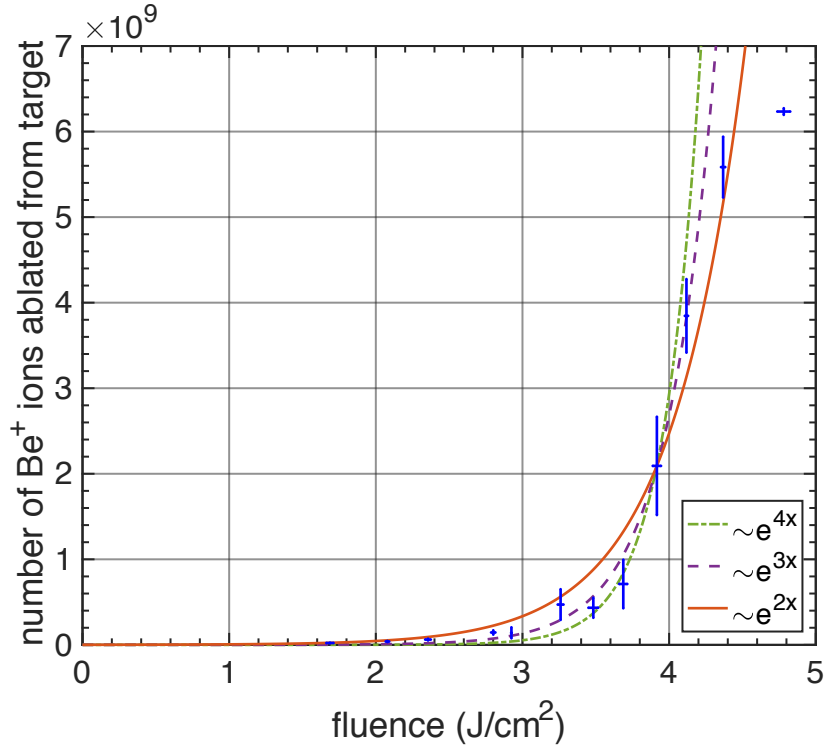
Figure 10.12: A plot of the number of ablated beryllium ions against laser fluence. The error bars at each data point are standard errors of the mean. A fitted curve of the form $A \left(1 + \operatorname{erf} \frac{x-\mu}{\sigma\sqrt{2}}\right)$ is also shown. The data suggests a Be^+ ablation threshold around 3 Jcm^{-2} . The signal begins to saturate and level off for yields $> 5 \times 10^9$ ions.



For fluences beyond the threshold, the yield rises slowly at first, only $\sim 0.7 \times 10^9$ ions at 3.7 Jcm^{-2} . Beyond 4 Jcm^{-2} however, the rise is dramatic. The data acquisition electronics again saturate when more than 5×10^9 ions are produced. In the case of beryllium, this happens at a laser fluence of 4.5 Jcm^{-2} . Beyond this point, the slow leveling-off of the signal is again observed. A curve of the form $A \left(1 + \operatorname{erf} \frac{x-\mu}{\sigma\sqrt{2}}\right)$ is fit to the data. It crosses

the 1×10^9 ion yield mark at a fluence of 3.75 Jcm^{-2} .

Figure 10.13: A plot of the number of ablated beryllium ions against laser fluence. The error bars at each data point are standard errors of the mean. Theoretical exponential curves proportional to e^{2x} , e^{3x} , and e^{4x} are also shown. The ion yield is consistent with an exponential rise with laser fluence.



In Figure 10.13, we plot the beryllium ion yield data along with three exponential curves that pass through the $(x = 3.9, y = 2.1 \times 10^9)$ data point. The ion yield clearly exhibits an exponential increase with increasing fluence; the $\sim e^{3x}$ curve having the best fit to the data. We are unsure however of the underlying physics process that is responsible for this exponential behavior and will only note that fluences of $> 3.5 \text{ Jcm}^{-2}$ most likely result in the generation of Be^{2+} as evidenced later in Figure 10.15.

10.2.2 Be⁺Kinetic Energy

Knowledge of the kinetic energy of the ablated beryllium ions is of utmost importance for their eventual use in the sympathetic cooling of positrons. The 313 nm cooling transition in Be⁺ ions exhibits a Doppler shift as a result of their non-zero axial velocities. The Doppler shift Δf is proportional to the velocity, and is given by

$$\Delta f = \frac{v_z}{c} f_0 \quad (10.1)$$

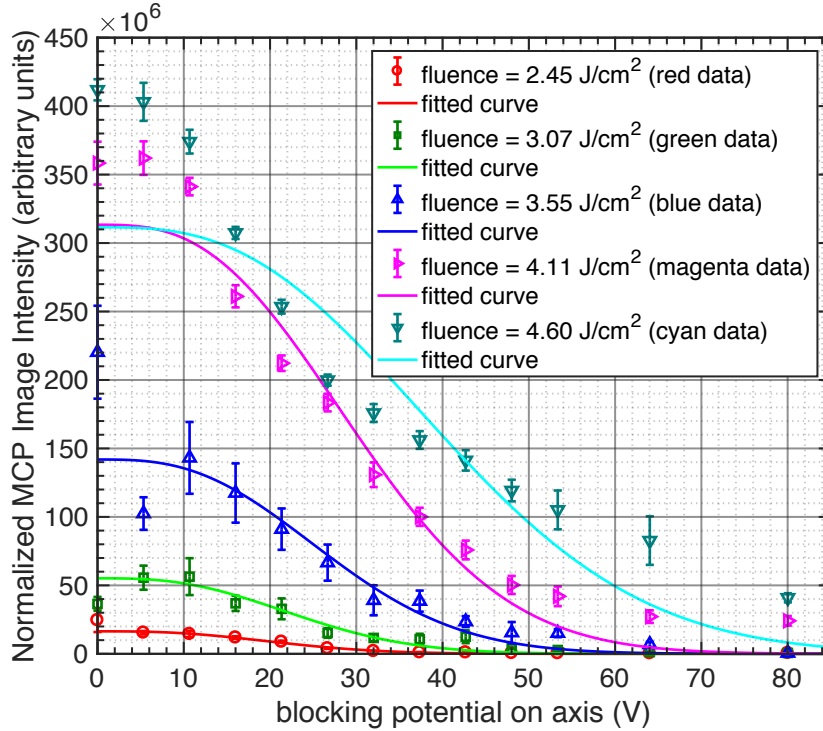
where f_0 is the frequency of the cooling transition, and c is the speed of light. Higher velocities therefore require the cooling laser to be red-shifted in frequency by larger amounts for the ion to be resonant with the laser. The extent to which the the laser's frequency can be modified then puts a limit on the maximum velocity an ion can have for it to be effectively cooled.

The same procedure that was used to obtain the kinetic energy profile of ablated aluminum ions was utilized for beryllium as well. A series of incrementally increasing blocking potentials were applied on electrode E3 and the resulting decrease in the intensity of the MCP image was recorded. The plot at five different fluences is shown in Figure 10.14.

For the range of fluences used in this measurement, a larger fluence clearly results in a higher overall ion yield, and no saturation is observed. Also, as reported earlier for aluminum, more than half of the ions produced, irrespective of fluence, have axial kinetic energies less than 30 eV. Such high fractions of low energy ions should hopefully translate into a high ion trapping rate in the Penning trap.

In Figure 10.14, Maxwell-Boltzmann cumulative distribution fits of the form $A \left(1 + \sqrt{\frac{2}{\pi}} \frac{x e^{-\frac{x^2}{2a^2}}}{a} - \text{erf} \left(\frac{x}{\sqrt{2}a} \right) \right)$ are also shown. For the three lower fluences, the fits appear to model the data reasonably well. For the two higher fluence of 4.11 Jcm⁻² and 4.60 Jcm⁻² however, the fits are poor, indicating that the kinetic energy distributions at these higher fluences are not necessarily

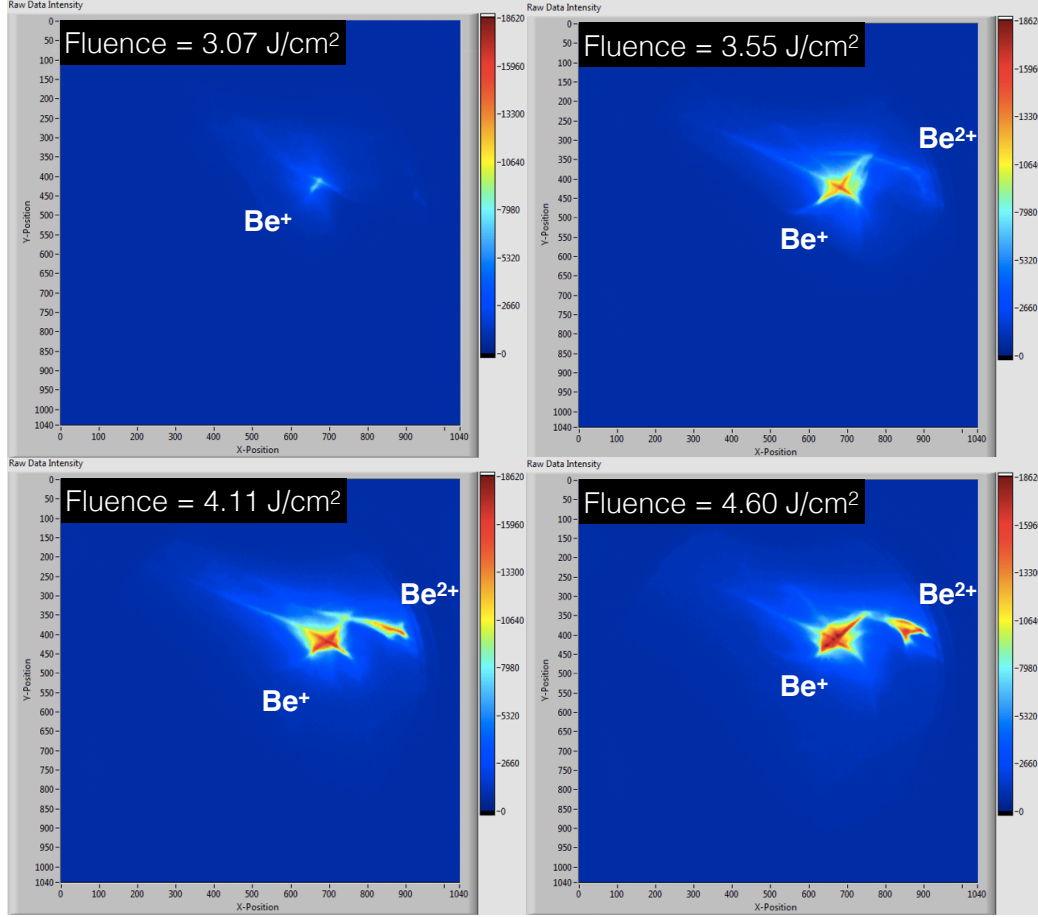
Figure 10.14: A plot of the MCP image intensity of beryllium ions vs electrode potentials for five different laser fluences. The background MCP image intensity is subtracted from each data point. The error bars shown are standard errors of the mean. A Maxwell-Boltzmann cumulative distribution function of the form $A \left(1 + \sqrt{\frac{2}{\pi}} \frac{x e^{-\frac{x^2}{2a^2}}}{a} - \text{erf} \left(\frac{x}{\sqrt{2}a} \right) \right)$ is fitted to the data. For blocking potentials > 30 V, more than half of the ions are blocked for all fluences.



a pure Maxwellian. Perhaps again, this could be due to the simultaneous production of some Be^{2+} ions at these higher fluences. Looking at the relevant MCP images in Figure 10.15, this conjecture is further reinforced since images of doubly-charged beryllium are clearly present. In fact, some of the images for 3.55 Jcm⁻² laser fluence also show Be^{2+} ions, but they are absent for fluences lower than that. This suggests a fluence threshold around 3.5 Jcm⁻² for Be^{2+} for 355 nm laser pulses.

Assuming that the image intensity recorded by the MCP is primarily that of singly-charged beryllium ions, we can use the Maxwell-Boltzmann cumulative distribution fits described above to estimate how both ion production

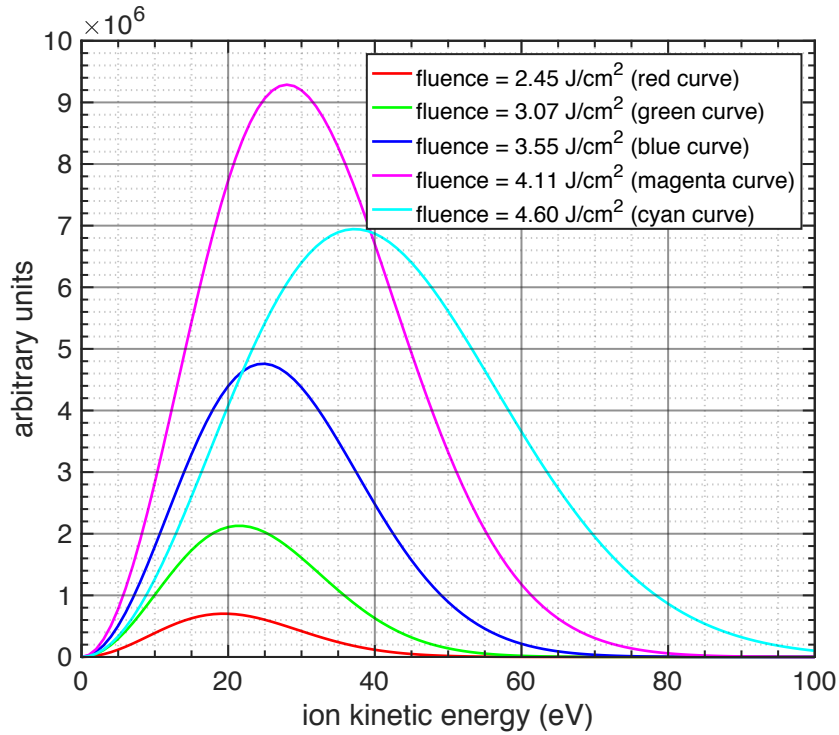
Figure 10.15: Typical MCP images for ablation carried out on a beryllium target at four different fluences.



and ion kinetic energy varies with laser fluence. The total number of ions produced (area under the curve) increases with increasing fluence. In fact, raising the laser fluence from 2.45 Jcm^{-2} to 3.07 Jcm^{-2} to 3.55 Jcm^{-2} and finally to 4.11 Jcm^{-2} results in the increased production of both low energy and high energy ions at each step. It is only in the last step change from 4.11 Jcm^{-2} to 4.60 Jcm^{-2} where a decrease in the proportion of ions below 40 eV is observed, indicating that at very high laser fluences, the production of low energy ions is significantly suppressed in favor of the production of much higher energy particles.

Figure 10.17 shows the Maxwell-Boltzmann kinetic energy distributions normalized by the total number of ions produced at each fluence. The average

Figure 10.16: A plot of the kinetic energy distribution of the ablated beryllium ions for five different laser fluences. The curves are based on the Maxwell-Boltzmann fits in Figure 10.14 and are scaled to reflect the relative ion production at each fluence. Between 2.45 Jcm^{-2} and 4.11 Jcm^{-2} , the total number of ions produced increases for all kinetic energies, resulting in an increased presence of both low energy and high energy ions. At the highest fluence of 4.60 Jcm^{-2} however, mostly high energy ions are produced.

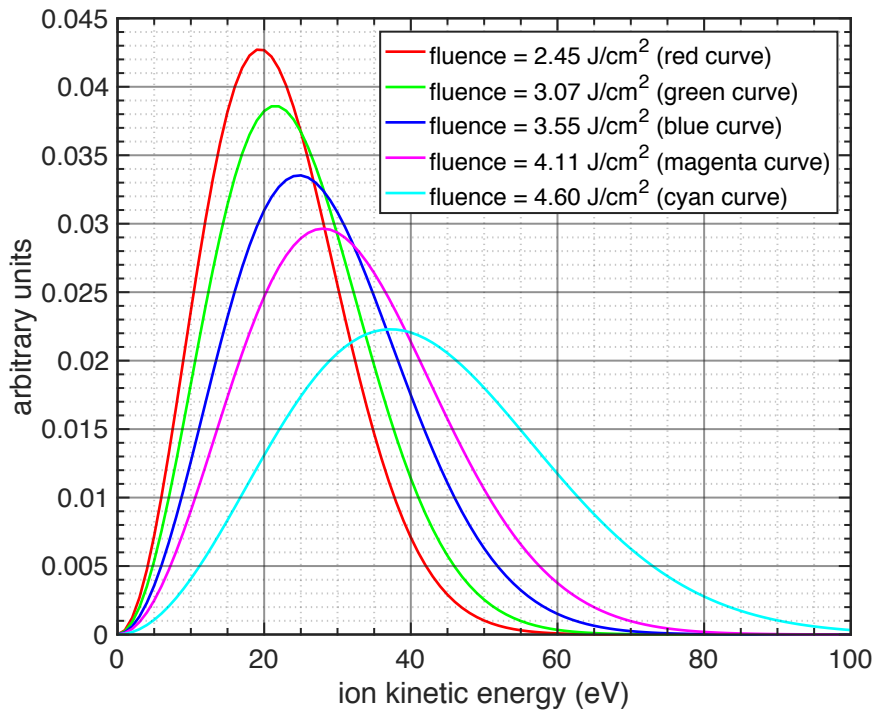


kinetic energy of the ions produced increases monotonically with increase in the laser fluence. At 2.45 Jcm^{-2} , the peak kinetic energy is just 19 eV, and almost all of the particles produced have energies less than 60 eV. For fluences at 3.07 Jcm^{-2} , 3.55 Jcm^{-2} and 4.11 Jcm^{-2} , the peak kinetic energies are 22 eV, 25 eV and 28 eV respectively, and almost all of the ions produced are within 80 eV. It is only at 4.60 Jcm^{-2} where the peak kinetic energy is $\sim 37 \text{ eV}$ and a significant fraction of ions are produced with energies greater than 80 eV.

The analysis presented here is only valid if the fitted curves obtained in

Figure 10.14 accurately represent the kinetic energy distribution of singly-charged ions. Given the poor fit to the data at higher fluences, we emphasize only the qualitative nature of our results and caution against their use in any quantitative manner.

Figure 10.17: A plot of the kinetic energy distribution of the ablated beryllium ions for the five different laser fluences used in Figure 10.16. The curves have been normalized against the total number of ions produced. With increasing fluence, an increase in the ion kinetic energy distribution is clearly evident.



10.2.3 Be^+ Bunch Length

For the beryllium ablation measurements discussed in the previous section, the blocking potential on a single electrode is applied before the ablation laser pulse is sent into the apparatus. When ablation takes place, the ions are released from the source towards the MCP. The more energetic ions are able to overcome the pre-erected potential barrier and move on to strike the

MCP. The less energetic ions are reflected back and eventually end up on one of several surfaces inside the ablation chamber. While this gives us the kinetic energy profile of the ions released by the source, it does not tell us the total length of the ion bunch that results from the ablation process.

The total ion bunch length can be experimentally measured by applying a potential on an electrode while the ions are still in transit towards the MCP. This can be done by following a three-step process

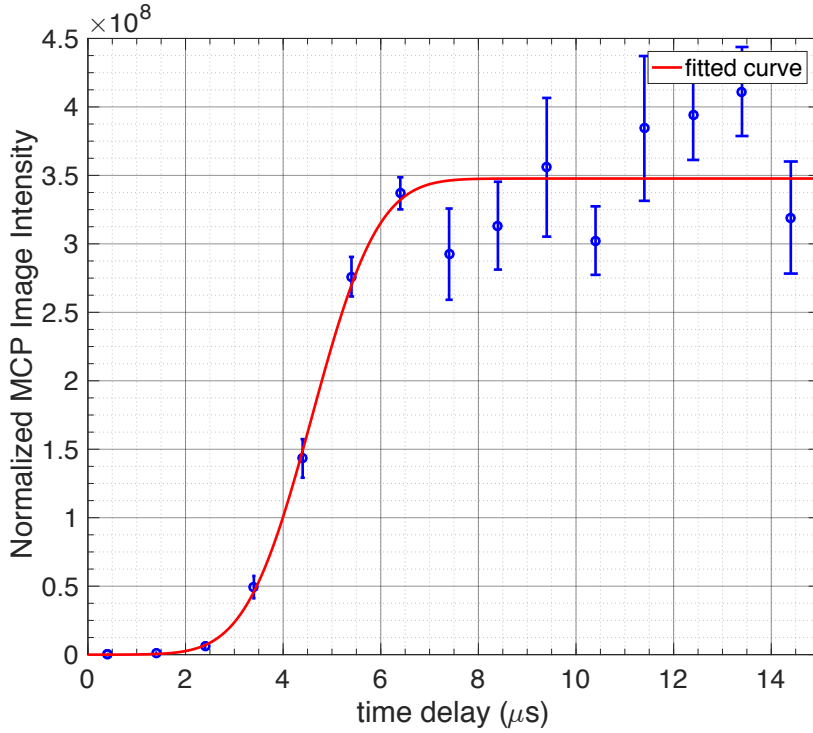
1. Send a single laser pulse to the ablation target to create ions
2. Wait a specified amount time before quickly applying an electrode potential (200 V) to one of the electrodes
3. Record the total number of ions arriving on the MCP

In this manner, only the ions that have traveled beyond the potential barrier before it is erected are able to reach the MCP. For a given time delay between the laser pulse and application of the potential barrier, the fraction of ions that reach the MCP is recorded. Scanning this delay time thus give us the overall temporal distribution of ions passing through the electrode.

Figure 10.18 shows the temporal distribution of ions for a fluence of $\sim 3.9 \text{ Jcm}^{-2}$ as they pass the middle electrode E3 in the ablation chamber. If the blocking potential is applied within $2\mu\text{s}$ of the laser pulse, no ions reach the MCP. If the delay is more than $3.5\mu\text{s}$, only the most energetic ions have passed the halfway point and reach the MCP. After approximately $7\mu\text{s}$, almost all ions have passed the middle of the trap.

A Gaussian cumulative distribution function of the form $\frac{A}{2} \left[1 + \text{erf} \left(\frac{x-\mu}{\sqrt{2}\sigma} \right) \right]$ is fit to this time-of-flight data. The curve is a good fit, indicating that the ablated ions have a normally-distributed temporal profile. The ions take an average time of $4.6 \pm 0.3\mu\text{s}$ to reach E3. Furthermore 90% of the ions pass the electrode between the $2.8\mu\text{s}$ ($\mu - 1.64\sigma$) and $6.3\mu\text{s}$ ($\mu + 1.64\sigma$), indicating a bunch length of $\sim 3.5 \pm 1.6\mu\text{s}$ (95% confidence level).

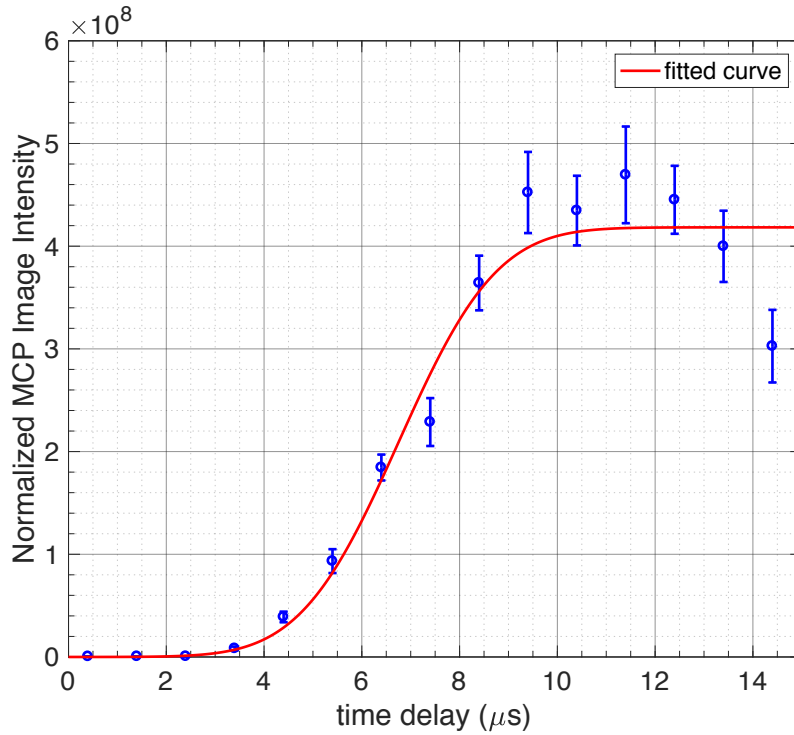
Figure 10.18: A plot showing the intensity of the MCP image for blocking potentials erected on electrode E3 at various times. The error bars are standard errors of the mean ($n = 10$). The MCP image intensity is proportional to the total number of ions reaching the MCP. For shorter delay times, only the fastest ions have passed E3 when the potential barrier is erected. The number is larger for longer delay times. A Gaussian cumulative distribution function is fitted to the data.



The measurement is then repeated for E5, the last electrode in the electrode stack. Figure 10.19 shows the associated time-of-flight data. The mean arrival time is $6.8 \pm 0.4 \mu\text{s}$ and the “90% bunch length” is $\sim 5.2 \pm 1.6 \mu\text{s}$. A Gaussian cumulative distribution function is also fitted to the data and appears to be a good fit as well. The last data point at $t \sim 14.5 \mu\text{s}$ is lower than expected due to a reduced pulse energy and consequently a lower ion yield for pulses at this particular time delay.

The observed time-of-flight data at E3 and E5 is consistent with the dynamics of a low energy ion beam. Ions ablated at the target travel $\sim 162 \text{ mm}$

Figure 10.19: A plot showing the intensity of the MCP image for blocking potentials erected on electrode E5 at various times. The error bars are standard errors of the mean ($n = 10$). The MCP image intensity is proportional to the total number of ions reaching the MCP. For delay times less than $4.2\mu\text{s}$, almost none of the ions have passed E5 when the potential barrier is erected. For delay times of $9.3\mu\text{s}$ or more, almost all ions have made it through. A Gaussian cumulative distribution function is fitted to the data.



to reach the center of E3. For fast ions that reach in $2.8\mu\text{s}$, this corresponds to an average drift velocity of 58 km/s . At this speed, the ion would reach the end of E5 (a total distance of 244 mm) in $4.2\mu\text{s}$, exactly as observed. Similarly, for slow ions that reach the middle of E3 in $6.3\mu\text{s}$, this corresponds to an average speed of 26 km/s , which would get the ions to the end of E5 in $9.4\mu\text{s}$, again within $0.1\mu\text{s}$ of the time observed.

Assuming the average drift velocities of the ablated ions are approximately constant over the length of the trap, we can extrapolate their trajectories to obtain an effective bunch length $\sim 1\text{ m}$ from the ablation target. This gives

us an expected bunch length of $21.5 \pm 6.6 \mu\text{s}$ if the same fluence conditions are used to load Be^+ ions in the ALPHA apparatus.

The fast and slow particles that flank each end of our ion bunch have axial kinetic energies of 156 eV and 31 eV respectively. These values are approximately a factor of 3 too large compared with the nearest Maxwell-Boltzmann fit shown in Figure 10.17 at 4.11 Jcm^{-2} (which give “90% bunch length” bounds at 50 eV and 11 eV). This is yet another indication that the ablated particles do not follow a simple Maxwell-Boltzmann distribution, but most likely encompass a superposition of two or more separate distributions.

10.2.4 MCP Calibration

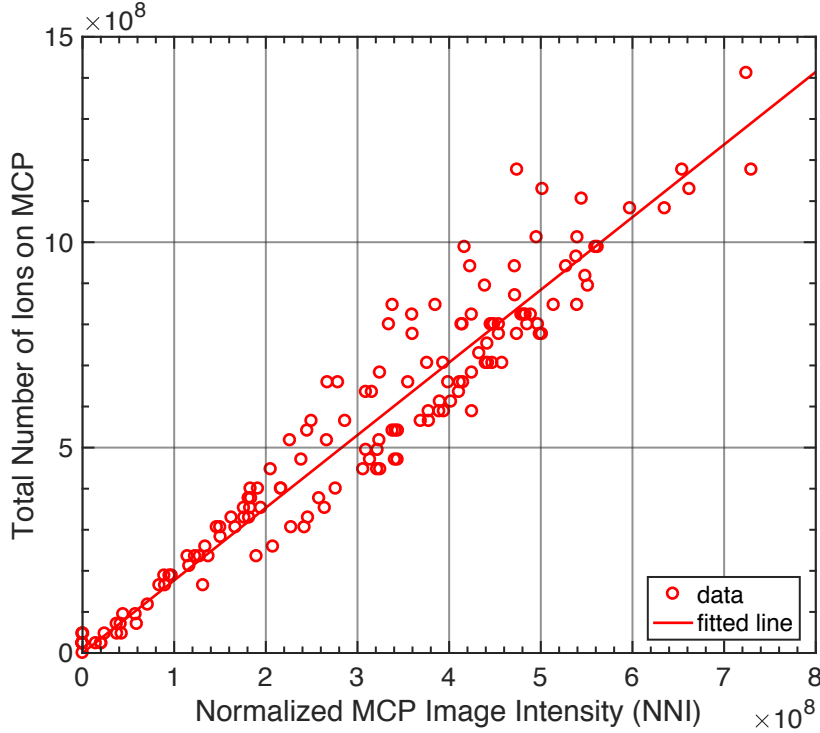
For the results discussed throughout this chapter, the MCP has been the primary detector to image the ions that arrive from the ion source. In order to determine the total number of ions reaching the MCP, it is important to calibrate the intensity of the MCP image against a particle counter. This is accomplished by operating the front of the MCP simultaneously as a Faraday Cup, similar to how the ablation target also doubles as a Faraday Cup.

For the time-of-flight data discussed in the previous chapter, the MCP was operated both as an imaging detector as well as a Faraday Cup. Figure 10.20 shows how the intensity of the MCP image scales with the total number of ions reaching the MCP. The MCP image intensity varies almost an order of magnitude from 1×10^8 to 8×10^8 . The number of ions detected by the Faraday Cup circuitry scales roughly linearly. The gradient of the fitted line is 1.77 ± 0.04 . The total number of detected ions is then given by the simple equation

$$\text{Number of Ions on MCP} = 1.77 \times \text{NNI} \quad (10.2)$$

where NNI is the Normalized MCP Image Intensity.

Figure 10.20: A plot showing the relationship between the normalized MCP image intensity and the total number of ions striking the MCP. For the experimental parameters investigated, the relationship is roughly proportional. The gradient of the fitted line, 1.77 ± 0.04 , is the proportionality constant.



Since the MCP Image Intensity has a much higher sensitivity and dynamic range than the Faraday Cup readout, this calibration allows us to work with much lower number of ions that would otherwise be undetectable by the Faraday cup. This will become important later on when we attempt to trap Be^+ ions and estimate the trapping efficiency.

10.2.5 Source Efficiency

Not all of the ions produced at the target leave the aperture of the ion source. The source efficiency, defined as the fraction of ions produced that escape the source from its aperture, depends on a number of parameters, in particular the bias of the target and the bias of the acceleration plate, both of which are

kept constant at 40 V and -10 V throughout. Using the MCP calibration from the previous section, we can determine the efficiency of the ion source.

Figure 10.21: A plot showing the relationship between the number of ions produced at the target and the total number of ions reaching the MCP. In the range of the measured data, the relationship is proportional. The gradient of the fitted line gives the production efficiency of the ion source: $(20 \pm 1)\%$

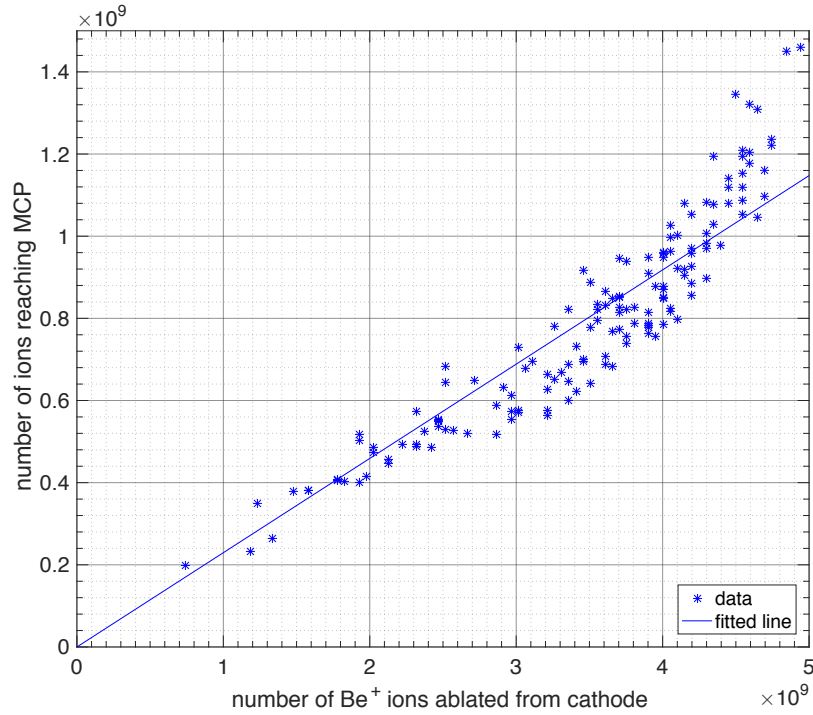


Figure 10.21 shows the variation in ion production due to the laser ablation of a beryllium target for 150 separate pulses. Laser fluences between 3.5 Jcm^{-2} and 4.3 Jcm^{-2} were used which resulted in the creation of between 1×10^9 and 5×10^9 ions at the target, shown on the horizontal axis of the plot. This is consistent with the ablation ion yields reported earlier in Figure 10.12. For each of these pulses, the ablated ions leaving the source were allowed to pass through the electrode stack unimpeded and an MCP image was recorded. Using the calibration from Eq. 10.2, the MCP image intensity was converted to total number of ions, shown on the vertical axis of the plot. The relationship between the two axes is approximately linear. A line

passing through the origin is fitted to the data. The gradient of the line is 0.20 with an error of ± 0.01 at the 95% confidence level. This indicates that the source produces useable ions with an approximate efficiency of 20%. In other words, approximately one-fifth of the ablated ions leave the aperture while the rest are lost within the volume of the source itself.

The 20% efficiency observed here is a factor of 20 larger than what is expected for 10 eV thermal ions. This is an indication that only a small fraction of useable ions are thermal ions, while the rest of the ions that leave the source aperture have relatively large axial velocities and highly directed trajectories in the axial direction.

10.2.6 Be⁺ Trapping

A simple four-step sequence can be used to dynamically trap ions from the ion source, followed by another two steps to observe the trapped ions on the MCP:

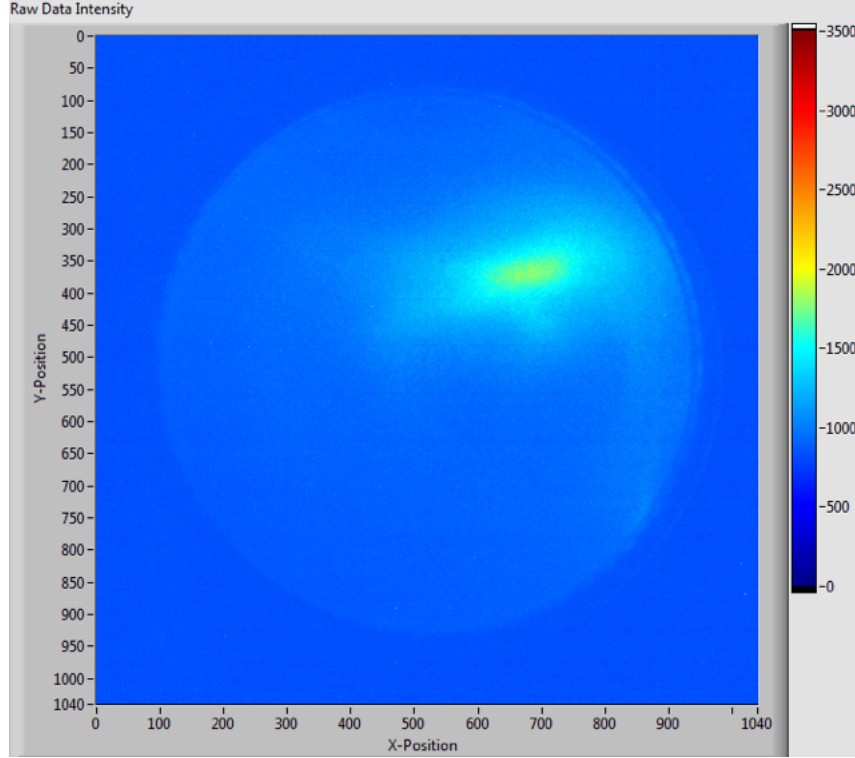
1. Pre-erect a potential barrier for the ions by applying an electric potential (200 V) on E5
2. Transmit a single laser pulse to the ablation target to create ions
3. Wait a specified amount time to allow ions from the ion source to enter the electrode stack
4. Quickly apply an electrode potential (200 V) on E1, effectively “closing the gate” on the ions located inside the stack.
5. Hold the trapped ions for the desired duration
6. Dump the trapped ions onto the MCP by quickly changing the potential on E5 to 0 V

From our ion bunch length measurements in Figures 10.18 and 10.19, we know that for the specific fluence of 3.9 Jcm^{-2} , it takes approximately $5.1 \mu\text{s}$

for the slowest moving (~ 31 eV) ions from the source to have passed E1. By closing the gate at this time, we ensure that most of the lowest energy ions are inside the electrode stack of the Penning-Malmberg trap. Unfortunately, since the ion bunch length is fairly long (~ 300 mm), the most energetic ions (i.e. with energy > 115 eV) have already reflected back from E5 and exited the stack before the gate is closed. Using these ablation parameters, it is not possible to obtain a 100% trapping efficiency.

Initially, trapping was attempted using a fluence of 3.9 Jcm^{-2} and a gate closing time of $8\mu\text{s}$. With these parameters, and following the trapping procedure described above with a holding time of 200 ms, no trapped ions were observed at the MCP. The gate closing time was then scanned from $0\mu\text{s}$ to $14\mu\text{s}$ in $1\mu\text{s}$ intervals, but still no trapping was observed. The number of trapped ions must be greater than 1×10^8 for the signal on the MCP to appear above the noise level. This would amount to a trapping efficiency of at least 10%. It was conjectured that the trapping efficiency is possibly much less than that, and therefore the trapped ions cannot be observed at the MCP with the given ablation parameters. By increasing either the laser fluence and/or the pulse energy, more ions can be produced from the ion source. Once the number of trapped ions exceeds the noise threshold, a signal on the MCP would confirm that beryllium ions can indeed be loaded into the Penning-Malmberg trap of the ablation chamber.

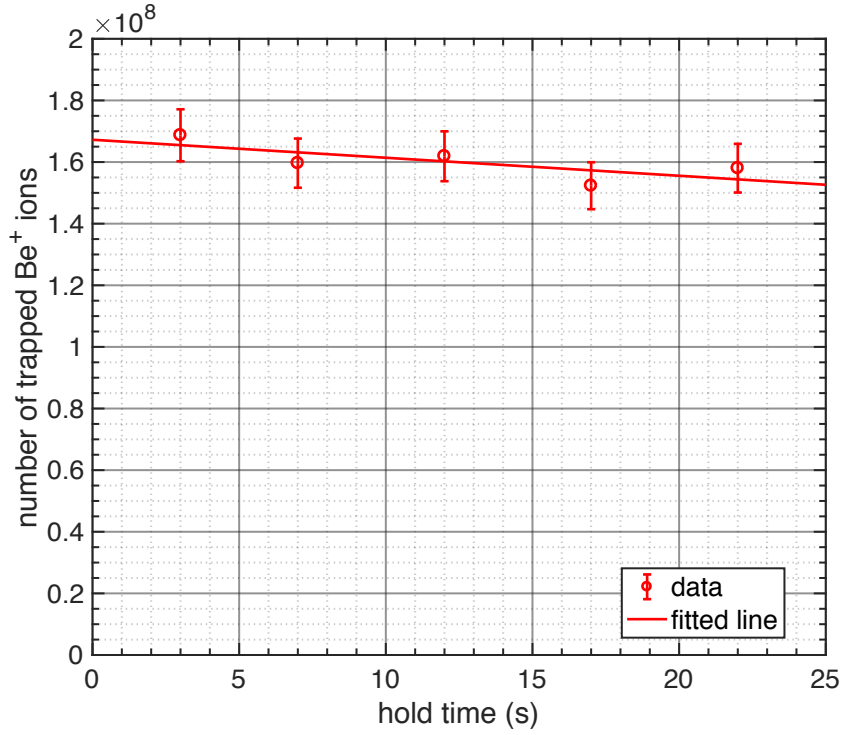
In order to produce more ions at the same laser fluence, the pulse energy was increased to $300\mu\text{J}$ and the focusing lens was translated by 7mm from its optimal position, resulting in a spot area of $8 \times 10^{-5}\text{cm}^{-2}$. With these modifications, the ablation pulse now had a fluence of 3.75 Jcm^{-2} but a significantly higher pulse energy. With the increase in the overall pulse energy, we would expect more ions to be produced and subsequently trapped in the Penning-Malmberg trap. Another attempt at trapping was therefore made. When the trapped ions are dumped onto the MCP, a faint MCP image (see Figure 10.22) is observed. Using the aforementioned MCP calibration, the number of trapped ions corresponds to approximately 1.7×10^8 , just above the detection threshold of the MCP.

Figure 10.22: An image of the trapped Be^+ ion cloud dumped on the MCP.

Ions trapped inside a Penning-Malmberg trap do not last forever. Magnetic field inhomogeneities and collisions with background gas particles may result in the expansion of the trapped plasma and the loss of particles. Since we ultimately intend to laser-cool the trapped Be^+ , it is imperative to ensure that the ion cloud can be trapped long enough for the individual particles to be laser-cooled. In Chapter 6, we established that laser-cooling would take on the order of a few ms, and mixing between the axial and transverse degrees of freedom takes place in a few seconds. It would therefore be ideal if we are able to keep the ions for at least ten seconds without significant degradation to the trapped cloud.

Figure 10.23 shows the number of trapped ions as a function of holding time. Beryllium ions were held for a minimum of 3 seconds up to 22 seconds, without any significant reduction in the number of trapped ions at these long holding times. The average trapping rate is 1.6×10^8 ions, more than two

Figure 10.23: A plot of the total number of trapped Be^+ ions for various durations of trapping. The average trapping rate is 1.6×10^8 and the total number of trapped ions falls very slowly over time, indicating that the ions are stably trapped.



orders of magnitude higher than what is required for the sympathetic cooling of positrons in ALPHA.

In order to estimate the trapping efficiency, we note that when all of the ions are allowed to strike the MCP, the MCP image is saturated, with the image intensity corresponding to the arrival of about 3×10^9 ions. Using this value, we determine an upper bound for the trapping efficiency of $\frac{1.6 \times 10^8}{3 \times 10^9} \times 100\% = 5.3\%$.

Although the absolute number of ions trapped is more than what is needed for sympathetic cooling of positrons in ALPHA, the obtained trapping efficiency of 5.3% is perhaps an order of magnitude lower than one would expect given the size of the Penning trap and the bunch length of the ion cloud. Given the

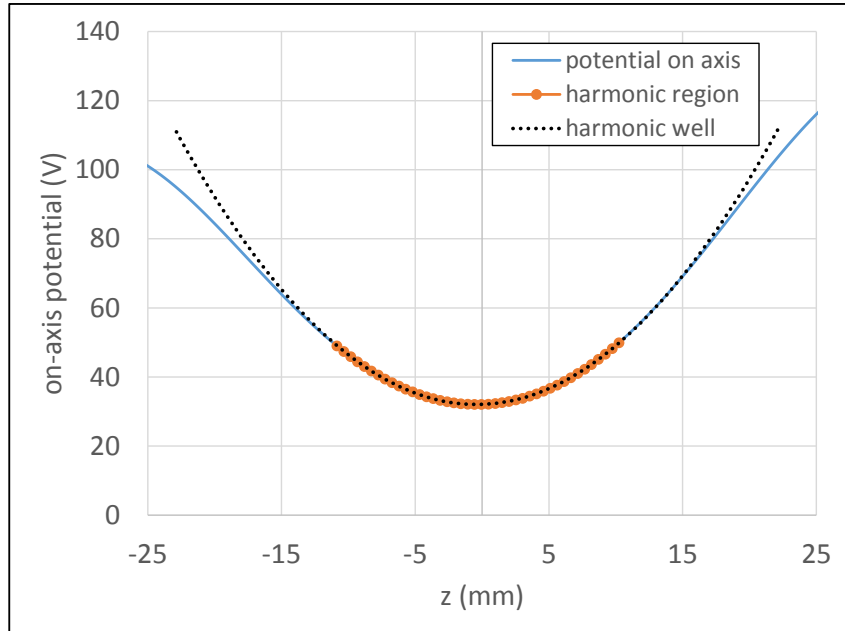
consistency of the trapping between 3 s and 22 s however, we suspect that most of the ions lost have transverse kinetic energies that are too large to be trapped in the nominal ~ 0.07 T magnetic field produced by 40 A in the solenoid. In other words, the Penning trap is unstable for beryllium ions at such a weak magnetic field. This can be verified by looking at the stability criterion given in Equation 2.8

From Equation 2.5, we can immediately calculate the cyclotron frequency for beryllium ions in the center of the Penning trap where the magnetic field strength is 0.07 T.

$$f_c = \frac{\omega_c}{2\pi} = \frac{qB_z}{2\pi m} \simeq 120 \text{ kHz} \quad (10.3)$$

At the ends of the trap, the magnetic field falls off to 0.05 T where the cyclotron frequency is only 85 kHz.

Figure 10.24: Graph of the electric potential on axis (blue curve) as a function of distance from the trap center. The potential at the center of the confining well (orange curve) can be approximated as a harmonic potential (black curve).



The axial bounce frequency, f_z , is a little more complicated to derive analytically since it depends on the shape of the electric potential well. For applied confining potentials of 200 V, the on-axis potential is shown in Figure 10.24. For ions confined in the center of the well, we can calculate the axial bounce frequency by approximating the axial potential as a harmonic well given by $U(z) = \frac{1}{2}k_z z^2$. By substituting the value of k_z into Equation 2.4, we have

$$f_z = \frac{\omega_z}{2\pi} = \frac{1}{2\pi} \sqrt{\frac{qk_z}{m}} \simeq 146 \text{ kHz} \quad (10.4)$$

In this case, $\sqrt{2}f_z > f_c$, indicating that the Penning trap is unstable. This explains the low trapping rate observed. For the potential well shown in Figure 10.24, stable trapping would require $B_z > 0.12 \text{ T}$. This is not currently achievable in our existing setup and requires replacement of the magnet power supply as well as improved water-cooling to keep the solenoid temperature within acceptable limits.

10.3 Summary

In summary, we have successfully demonstrated the production of aluminum and beryllium ions using single ablation pulses from a 355 nm laser. The production of Al^+ requires at least 1.0 Jcm^{-2} laser fluence whereas the ablation threshold for Al^{2+} is around 2.5 Jcm^{-2} . In our experiments, we observe higher ion yields and higher kinetic energies with increasing laser fluence. For fluences up to 3.0 Jcm^{-2} , most of the aluminum ions produced have axial kinetic energies $< 80 \text{ eV}$, low enough for them to be blocked and trapped in our apparatus. The observed fluence thresholds and kinetic energy distributions are consistent with available literature on aluminum ablation, making our test apparatus a reliable tool to study other metals of interest, in particular beryllium.

For beryllium, we observe ablation thresholds around 2.0 Jcm^{-2} and 3.5 Jcm^{-2} for Be^+ and Be^{2+} respectively. As with aluminum, using higher fluences for ablation result in more energetic ions and larger ion yields. Fluences

up to 4.0 Jcm^{-2} may be employed to produce ions that have kinetic energies less than 80 eV and ion yields in excess of 10^9 , suitable for trapping. This corresponds to pulse energies of only $35 \mu\text{J}$, thus very low laser powers are required to meet the production requirement in ALPHA. We obtain an ion source production efficiency of about 20%, so only a fifth of all ions produced actually leave the source.

By choosing the appropriate electrode potentials and gate timing, ions released by the source may be trapped in a Penning-Malmberg trap. For the trap used in our experiment, the electrode stack is relatively short compared to the nominal bunch length of the ion cloud, therefore only a small fraction of the ions can be trapped. Additionally, the trap is unstable for beryllium ions due to the low axial magnetic field. For a specific set of parameters, we obtain an upper limit on the trapping efficiency of beryllium ions at 5.3%.

Repeatedly using laser ablation to produce ions does not adversely affect the trap vacuum. Additionally, the number of particles produced as a result of ablating beryllium is well below the recommended occupational exposure limits. Since we intend to install the ion source in the stick of the main ALPHA apparatus, this is an important result that allows us to proceed with the program of sympathetically cooling positrons in ALPHA using laser-cooled beryllium ions.

Chapter 11

Conclusions and Future Work

One of the best ways to study antimatter is to investigate antihydrogen. With our detailed knowledge of the hydrogen atom, both theoretically and experimentally, probing antihydrogen allows us to make high-precision comparisons between the two. Members of the ALPHA collaboration, including this author, have recently measured the charge of antihydrogen to < 0.71 parts per billion [36] and performed laser spectroscopy on trapped antihydrogen with a relative precision of about 2×10^{-10} [37]. At the time of writing this thesis, the full line-shape measurement of the 1S-2S transition in antihydrogen is underway.

Antihydrogen is synthesized at ALPHA by merging a cloud of antiprotons with a cold positron plasma. In the magnetic atom trap, only a fraction of the coldest atoms remain trapped, while the rest escape and annihilate on the trap walls. In Chapter 5, I have summarized previous studies that outline the positron temperature dependence of antihydrogen formation, and have also analyzed new data from the 2016 run which is consistent with previous results. In addition, I have analyzed and presented the first empirical scaling law of antihydrogen trapping rate with positron plasma temperature: $T^{-1.8 \pm 0.6}$.

Since colder positrons lead to higher antihydrogen trapping rates, the possibility of sympathetically cooling positrons was theoretically investigated by

Madsen *et al.* and deemed plausible with laser-cooled beryllium ions. That work has been reviewed in Chapter 6. The main focus of this PhD has been to demonstrate a beryllium ion source that uses pulsed laser ablation as the mechanism of ion production and that is simultaneously compatible with the strict electromechanical, thermal, vacuum and space constraints of the ALPHA apparatus. That work has been detailed in Chapters 8, 9, and 10. In short, the suitability of the ion source has been verified and a smaller version of the source has been integrated into ALPHA.

Figure 11.1 shows the design of the modified ion source that has been installed. This design is slightly smaller than the original version in order to fit into the vacuum chamber, and a mounting manifold has been added in order for the source to be attached to a new stick. Figure 11.2 shows the design of the new stick, along with the ion source, an electron gun, an MCP, and a fluorescence detector that are mounted as shown. These components are part of an upgrade package intended to improve the functionality and diagnostic ability of the experiment.

Figure 11.1: The modified beryllium ion source intended for use in the main ALPHA apparatus.

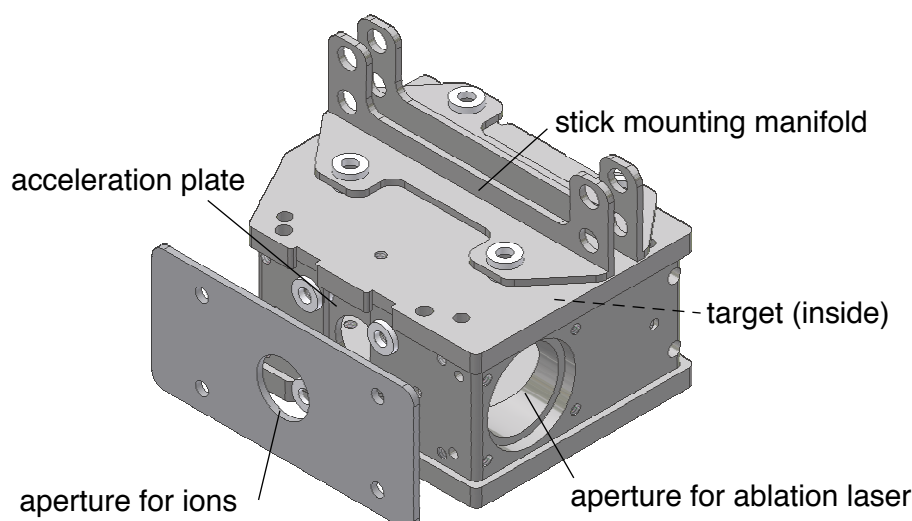
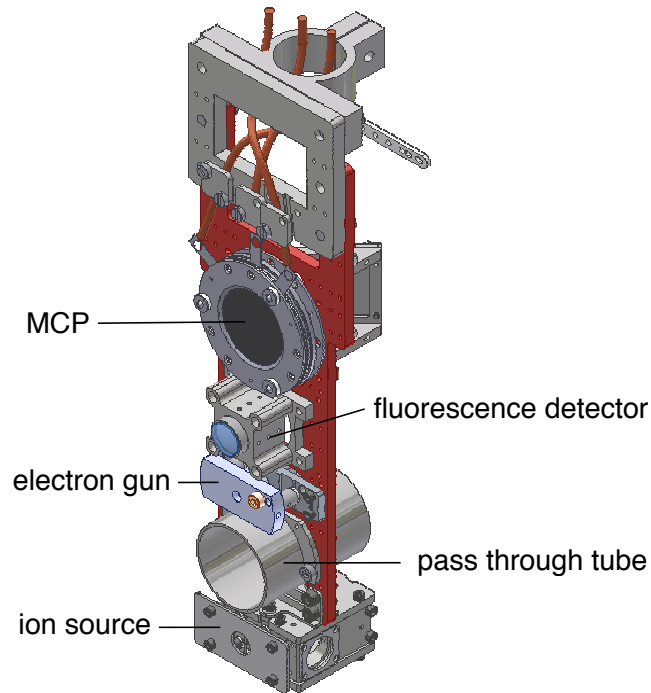


Figure 11.2: The new stick that will be installed upstream of the Atom Trap.



Once the new stick is commissioned, steps will be taken to

- trap beryllium ions in the Mixing Trap and obtain the trapping efficiency,
- laser-cool the beryllium ions using the 313 nm laser, and
- sympathetically cool the positrons using the laser-cooled beryllium ions.

Once sympathetic cooling of positrons has been demonstrated in the Atom Trap, it can be implemented in the antihydrogen production sequences. The positron preparation steps would be as follows:

1. Catching and cooling of positrons in the the positron trap
2. SDREVC of positrons
3. Catching and cooling beryllium ions with the positron cloud
4. Rotating-wall compression and cyclotron cooling of the mixed plasma

5. Transfer of mixed plasma to the mixing trap
6. Laser-cooling of beryllium ions and sympathetic cooling of positrons

As the sympathetic cooling stage progresses, the beryllium ions would slowly move outwards and centrifugally separate from the positrons. Antiprotons may then be axially injected into the plasma to create and trap antihydrogen atoms as before.

By cooling the positrons to temperatures of ~ 5 K before mixing, we expect to see an increase in the antihydrogen trapping rate by at least a factor of 10. This would be an unprecedented improvement in our ability to make physics measurements: instead of taking several days or weeks, we would be able to perform some experiments in a single day, or perhaps even within a single trapping cycle. Although not strictly necessary, an improved antihydrogen trapping rate would also be invaluable in measuring the behavior of antihydrogen in the Earth's gravitational field [122].

With the first precision measurements on antihydrogen, the commissioning of the ELENA ring, the approval of ALPHA-g, and the maturation of several of the AD experiments, we have entered an exciting new phase where long-standing questions about the properties of antimatter can finally be answered.

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