

PIONEERING EXPERIMENTAL APPLICATIONS USING TIME PROJECTION CHAMBERS

By

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A DISSERTATION

Submitted to
Michigan State University
in partial fulfillment of the requirements
for the degree of

Physics—Doctor of Philosophy

August 14, 2026

ABSTRACT

A Time Projection Chamber (TPC) is a gaseous detector that is used to reconstruct 3D tracks from various nuclear reactions. Charged particles interact with a detection gas inside the TPC and ionize electrons, which are then drifted towards a pad plane via an applied $\mathbf{E}$ -field to be collected on a time-based readout. Fine granulation of pads on the pad plane provide two dimensional tracking, and electron drift time coupled with the charge collected on pads as a function of time provides the third dimension. Particle tracks can then be inferred using this information. Nuclear reactions may be induced outside of the TPC by placing a target in front of the detector on which an energetic beam is impinged, for example. By using the detection gas itself as the target, a TPC can also serve as an “active target” to induce reactions inside the detector. Further information about the measured particle tracks can be obtained by placing the TPC inside a $\mathbf{B}$ -field, so that the charged particles follow curved paths in relation to their rigidity.

This dissertation will detail pioneering applications of Time Projection Chambers to probe highly excited nuclei and nuclear matter created in two distinct experiments. In the first experiment, detailed in Part I, the Active-Target Time Projection Chamber (AT-TPC) is employed to measure helium induced fusion-fission of neutron deficient isotopes near the lead region. Results of this experiment will provide cross-section and charge distribution measurements along the transition region from symmetric to asymmetric fission. In the second experiment, detailed in Part II, the $S\pi$ RIT TPC is employed to measure charged pion production resulting from heavy ion collisions (HICs) between Xe beams and Sn targets. Results of this experiment will provide constraints to the symmetry energy term of the Equation of State (EoS) for nuclear matter above saturation density.

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ACKNOWLEDGEMENTS

Firstly, I would like to thank my guidance committee members: Professors Alex Brown, Kyle Brown, Daniel Hayden, Bill Lynch, and Betty Tsang, for being patient with my academic journey. The proposed contents of this dissertation have changed over the years, and ultimately resulted in the combination of two separate projects, which is not the most typical case. Through the mentoring of my advisor Bill, I was able to participate in several experimental campaigns and develop the skills and experience which have made me into the physicist I am today.

This research would also not have been possible without the support of several collaborations and their wonderful scientists, who should also be acknowledged. Thank you to the previous members of the **High Resolution Array (HiRA)** who were present during my time in the group and have helped me in one way or another. It has been a pleasure to work with you all.

- **The postdocs:** Curtis Hunt, Genie Jhang, Taras Lokotko, Chenyang Niu, Som Paneru, and Rensheng Wang.
- **The graduate students:** Adam Anthony, Jon Barney, Justin Estee, Sean Sweany, Fanurs Teh, Tommy Tsang, and Kuan Zhu.
- **The undergraduate students:** Olivia Gunter, Bergen Kendziorski, Andrew Pype, Sarah Wegert, and Angus Wong.

Thank you to our collaborators at **Western Michigan University**, who shared their lab resources and valuable time with us.

- **The graduate students:** Jesús Diaz Aponte, Kin Tam, and Teajun Kim.
- **The professor:** Zbigniew Chajecki.

Thank you to the many past and present members of the international **S π RIT collaboration** and **BIGRips** team, who are too numerous to list in their entirety. There are a few who I believe deserve specific mentions for the great assistance they have personally given to me.

- **Korea University:** JungWoo Lee and Jeonghyeok Park.
- **RIKEN Nishina Center:** Tadaaki Isobe and Mizuki Kurata-Nishimura.
- **King Mongkut's University of Technology Thonburi:** Suwat Tangwanchaoen.

Thank you to the many members of the **Transport Model and Evaluation Project (TMEP)**, who have taken the time to simulate our experiments with their transport models, and bravely shared their predictions with us before seeing any results. Special thanks to Dan Cozma, Natsumi Ikeno, and Akira Ono for providing the comparisons that supplement this dissertation.

It was a privilege to work with my fellow classmates who braved the trials of graduate courses and preparing for subject exams with me; we supported each other and exchanged ideas often, and you taught me different ways of thinking. Some of you I also had the opportunity of living with throughout my time at MSU, and sharing weekly home cooked meals with; our gatherings gave me a sense of community. My various FRIB Room 2019 office mates were the best anyone could have asked for; you made me feel welcome and included, which I know is not something everybody can claim.

My loving and supportive family was there for me at every step in this journey. Thank you for asking about my projects and trying your best to understand them, and for always believing in me. To my Grandmother, Aunts, Uncles, Cousins, Parents, and Sister, I hope to make you all proud.

Finally, thank you to my beeb Sydney. I most certainly could not have done this without you by my side.

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CHAPTER 1

DISSERTATION ORGANIZATION

In this dissertation, we make pioneering applications of Time Projection Chambers to probe highly excited nuclei and nuclear matter created in two distinct experiments. The first experiment detailed in Part I focuses on fusing heavy radioactive beam isotopes with light helium, such as ^{200}Bi to form ^{204}At compound nuclei at a low excitation energy of 20-40 MeV. This experiment studies fission decays induced by large Coulomb forces, which make the highly excited combined compound nucleus unstable with respect to fission. In principle, such data are sensitive to the mean field potential that binds the nucleus and to its behavior at low densities. Since this fission experiment also contributed to the dissertation of another student [17], the focus on this research topic is on the investigation and correction of space charge effects within the detector used to measure the fission. This correction method was recently published with the current PhD student as first author [18].

Chapter 2 introduces fission as a nuclear process and motivates the fission experiment and its particular region of study. Chapter 3 explains the NSCL accelerator facility and detector systems used to create and measure the radioactive beam (HEIST), and record the fission reactions (AT-TPC). Chapter 4 details how the radioactive beam isotopes were tagged and identified in flight pre-fission, and how charge state contamination was measured. Section 4.4 focuses on the methods developed to correct for the beam induced space charge effects.

The major focus of this dissertation in Part II is on an experiment to constrain the EoS of dense matter by studying central collisions of two heavy nuclei of approximately equal mass at higher incident energies of 320 A MeV. In these central collisions, some of the nuclear matter is compressed to approximately $0.21 \text{ nucleons/fm}^3$, which corresponds to about 1.5 times the saturation density of $0.14 \text{ nucleons/fm}^3$. In contrast to the fission work, portions of the compressed nuclear systems can be excited to excitation energies in excess of 20 MeV/nucleon before expanding and disintegrating. From these measurements, constraints on the Equation of State of Nuclear matter at densities of $0.21 \text{ nucleons/fm}^3$ are obtained.

Chapter 5 introduces the nuclear matter Equation of State and Symmetry Energy, and how

charged pions can serve as a probe. The context of a previous measurement is also given, which motivates this more recent experiment. Chapter 6 explains the RIKEN RIBF facility used to create the heavy ion collisions measured in the experiment, and the operation of the S π RIT TPC used to record them. Chapter 7 details the analysis pipeline for reconstructing charged particle tracks from these collisions, and how we select the more central collisions. Chapter 8 explains how pions are isolated in the data, and what corrections need to be applied in order to obtain absolute yields. Preliminary pions spectra and yield results are presented, and compared with some preliminary predictions from a few transport codes.

Part I

Fission with the AT-TPC

CHAPTER 2

INTRODUCTION

In this chapter, the nuclear process of fission is introduced along with different ways to model it. Recent observations that motivated the experiment discussed in this part of the thesis are mentioned.

2.1 Background

In this section, I introduce some important properties of *fission* as a nuclear process and highlight recent groundbreaking results that have sparked increased interest in studying fission of exotic nuclei.

2.1.1 History of Fission

Fission is essentially the splitting of one atomic nucleus into two smaller (lighter) nuclei. It was first observed by two Berlin chemists in 1938 [19], Otto Hahn and Fritz Strassmann, when they detected barium ($Z = 56$) after bombarding uranium ($Z = 92$) with slow neutrons. Unsure of how to physically interpret such a result, they sent their findings to two physicists, Lise Meitner and Otto Frisch. Meitner and Frisch proposed a model that assumed the Uranium split into two nuclei of roughly comparable charge. Frisch chose to name the process *fission* after cell division, and later in 1944, Hahn was granted the Nobel Chemistry Prize. It is interesting that Fermi actually observed fission earlier, but interpreted the measurement incorrectly and this wrongly interpreted discovery was included among the accomplishments that earned Fermi a Nobel prize. The chemical studies of Hahn and Strassmann were essential to identifying that this reaction was definitively fission and not some other process.

Figure 2.1 illustrates the Uranium fission decay process measured by Hahn and Strassmann. It starts when a slow neutron is captured by a Uranium target nucleus. This compound nucleus (CN), formed from the neutron and Uranium nucleus, then elongates and eventually fissions into two lighter fission fragments plus a few neutrons. One fission fragment was observed to be an isotope of barium and by conservation of charge, the other fission fragment must have been an isotope of bromine. Since its discovery, fission of many different heavy nuclei has been induced via many different reaction mechanisms, including cases of spontaneous fission of an isolated heavy

fissionable nucleus, or fusion-fission reactions that occur when a light-ion and a heavy-ion collide and form a CN that subsequently undergoes fission. Much remains to be learned. For example, little is known about the fission of nuclei produced in the astrophysical r-process [20, 21]. In the fusion-fission example, the role of fission in the r-process largely remains unquantified because many of these nuclei are not experimentally accessible.

Now, we turn to the energetics of the fission process, which most often occurs in nuclei that are axially deformed with a long axis, which helps to minimize the influence of the repulsive Coulomb force. It helps to understand these heavy nuclei as liquid droplets of nuclear matter, with central densities of approximately $0.16 \text{ nucleons/fm}^3$. Such nuclei can fission when capture of a neutron heats the nuclear liquid and the increase in pressure makes it expand slightly in the direction of its longer axes.

This increase in total energy with elongation ends when it becomes energetically more favorable for the nucleus to fission into two fission fragments. The onset of fission decays as a function of excitation energy defines a threshold excitation energy for fission decay and an associated *fission barrier* height.

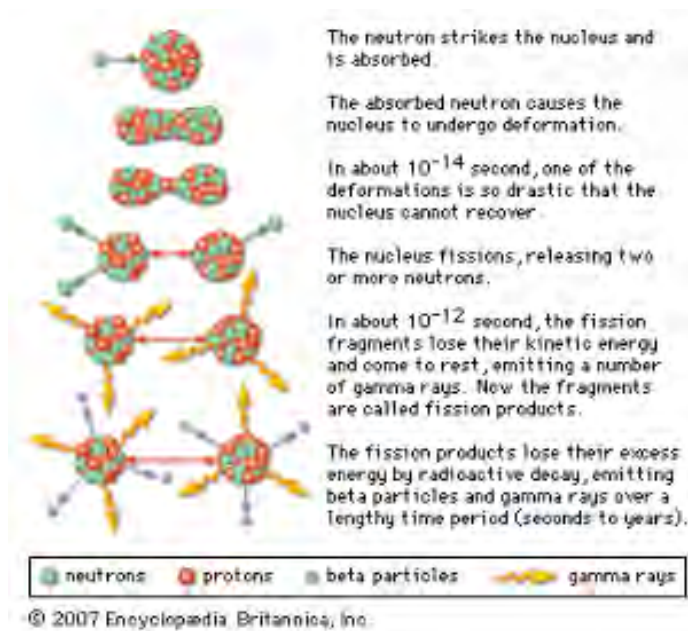


Figure 2.1 Illustration of fission induced by slow neutron bombardment. Because of the low kinetic energy of the neutron, it is more likely to be captured by the nucleus than to simply scatter.

2.1.2 The Liquid Drop Model

The energetics of the fission barrier are important to understand and this requires a model. The nuclear **Liquid Drop Model** (LDM) approximates the nucleus as a drop of a charged incompressible liquid. This model provides predictions for the measured mass of a nucleus as the difference between the mass of its constituent nucleons, and a simple expression for the so-called ‘*Binding Energy*’ (BE), which is the difference between the rest energy of free nucleons and the rest energy of the nucleus:

$$m = N * m_n + Z * m_p - \frac{BE(N, Z)}{c^2} \quad (2.1)$$

This liquid-drop or Semi-Empirical Mass Formula (SEMF) provides a simple but fairly accurate formula for the binding energy, BE(N,Z), as the sum of four terms that depend on nucleon numbers plus an additional pairing term. The form of these terms is inspired by the concept of nuclear matter being a liquid comprised of nucleons, where each nucleon can be in one of four possible states labeled by the spin projection +1/2 (neutrons) or -1/2 (protons) along some spatial direction and the isospin of the nucleon.

If we neglect the Coulomb, spin, and isospin dependent forces, all nuclei with same number of nucleons A will have the same mass. Summing of all nucleons this “volume” or “bulk” yields a binding energy proportional to ‘ A ’. For simplicity, this has been proposed as an approximation and values for the **Volume** term have been extracted. In detail, however, nuclei are finite. Thus, nucleons at the outer surface of the sphere have fewer neighbors to interact with, and while there, the binding energy will be reduced by a correction to the Volume term called the **Surface** energy term. The density of matter in the interiors of heavy nuclei is 0.16 nucleons/fm³. If the nucleus is a sphere of volume ‘ A ’ nucleons, and therefore with radius ‘ $A^{1/3}$ ’ nucleons, the Surface term should scale like surface area: ‘ $A^{2/3}$ ’.

Next, we consider the effect of Coulomb repulsion between protons. The classical electrostatic potential between two charges scales as the product of the two charges divided by the distance between them, i.e. $q_1 q_2 / r$. Each proton (Z) is repelled by the remaining ($Z-1$) protons, and counting the number of protons leads to **Coulomb** energy that is proportional to ‘ $Z(Z-1)/A^{1/3}$ ’.

There is no Coulomb force between neutrons, as they have no charge, but there is another effect stemming from the Pauli Exclusion principle, which depends on both the neutron and proton number and depends on proton-neutron asymmetry. Because no two identical fermions can occupy the same state within the nucleus, a system of a given number of nucleons will be forced to occupy higher and higher energy levels for larger nuclear asymmetry, defined by $|(N - Z)|$. This effect should depend on the absolute difference in number of protons and neutrons (no sign dependence), and the same level of asymmetry should have a smaller effect on systems with a larger number of total nucleons. A reasonable accurate estimate of this *Asymmetry* term can be achieved if it scales like $(N-Z)^2/A$.

Finally, there is a small effect coming from spin coupling of nucleons. For atoms with an odd number of total nucleons, this effect is zero. For even (N,Z) binding energy is added, and for odd (N,Z) it is subtracted. This *Pairing* term is thus represented via a Kronecker-Delta symbol (δ), and does not depend on total nucleon numbers.

Altogether then, the BE can be represented as

$$BE(N, Z) = a_V A - a_S A^{2/3} - a_C \frac{Z(Z-1)}{A^{1/3}} - a_A \frac{(N-Z)^2}{A} \pm \delta(N, Z), \quad (2.2)$$

where the constants a_V , a_S , a_C , and a_A , correspond to the Volume, Surface, Coulomb, and Asymmetry terms, respectively. Equation 2.2 is also known as the Bethe–Weizsäcker formula [22]. To facilitate conceptualization of these different terms, a graphic has been provided in Figure 2.2. The values of these terms for various nuclei as described by Equation 2.2 are shown in Figure 2.3. It is possible to predict values for these constants that will reproduce measured values to within a few MeV, but in practice they are fitted to experiment. The agreement between the predicted and measured values is demonstrated in Figure 2.4.

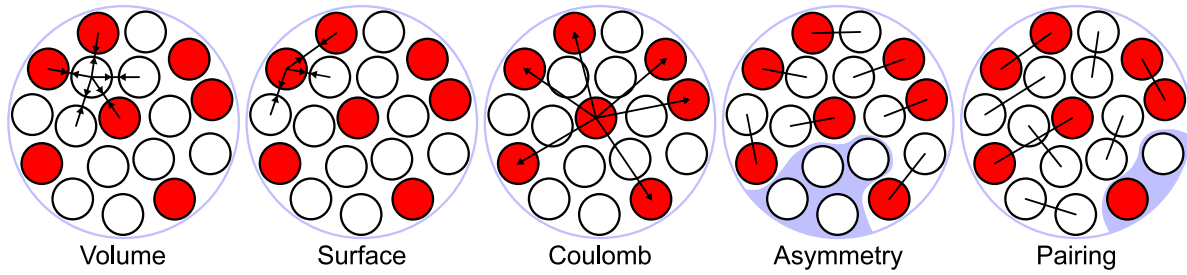


Figure 2.2 A graphical illustration of the various terms in the SEMF, where red circles indicate protons and white circles indicate neutrons. Figure from [4].

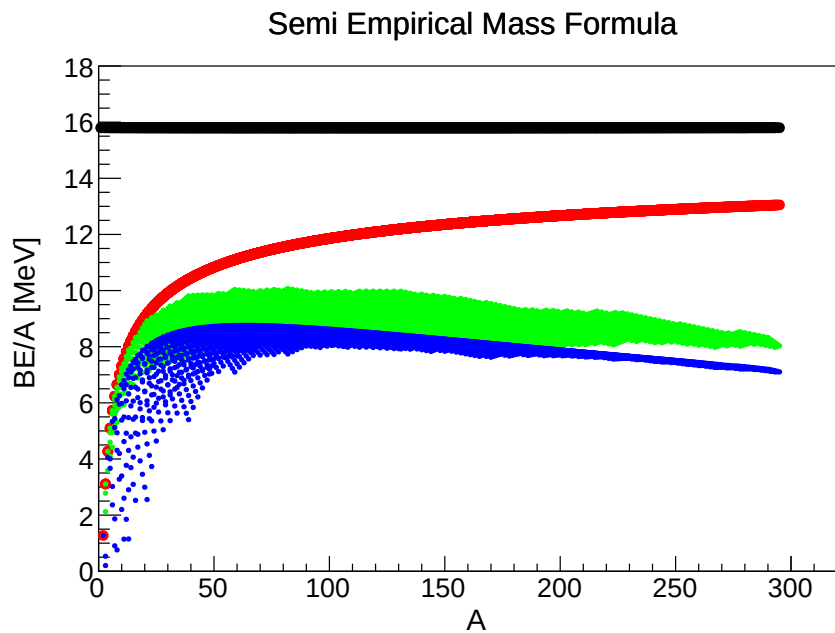


Figure 2.3 Numerical values for the terms of the SEMF for a selection of isotopes. (Black) The volume term. (Red) Volume + surface terms. (Green) Volume + surface + Coulomb terms. (Blue) Volume + surface + Coulomb + Asymmetry terms.

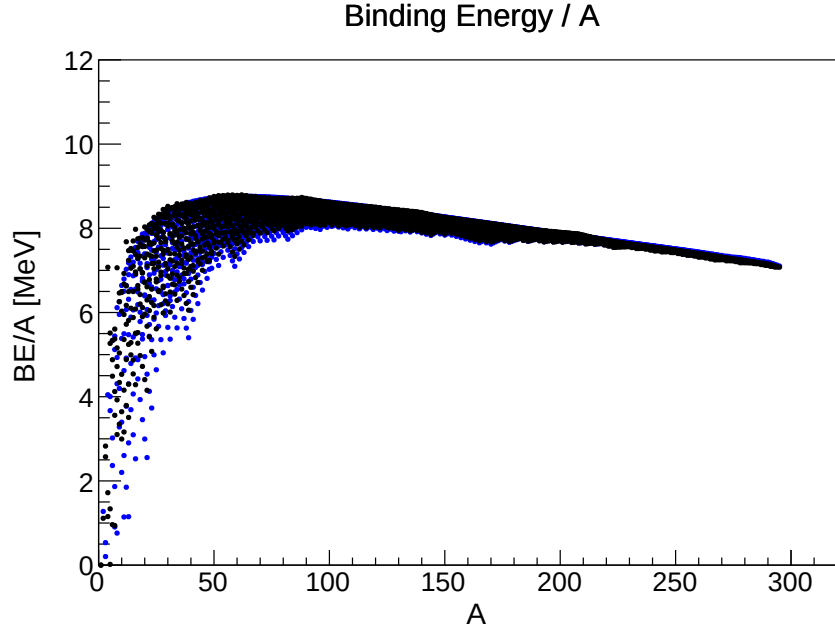


Figure 2.4 An overlay of binding energies for the most stable nuclei. The LDM prediction is shown in blue, while the experimental measurements are shown in black. Experimental values are taken from [5].

2.1.3 Physics Motivation

The SEMF contains the main geometric factors that govern fission. However, it is too simplistic to predict the astrophysical relevance of fission, which plays an important role in the rapid neutron capture process (r-process) that governs the creation of the heaviest elements. This empirical scaling is not sufficiently accurate to predict the energetics and probability for many of the important nuclei in the r-process, and this makes it too imprecise to accurately predict the role of fission in the cosmos. Figure 2.5 shows the potential energy surface (PES) for ^{238}U in two dimensions (elongation and mass asymmetry) [6]. In the left panel, the PES considering only the SEMF in the LDM would predict symmetric fission into two fragments equal in mass. In the right panel, nuclear shells are pairing effects are taken into account by the PES to show a more realistic result where the result of fission is two fragments with different masses, favored by energy minimization through closed shells in one of the products.

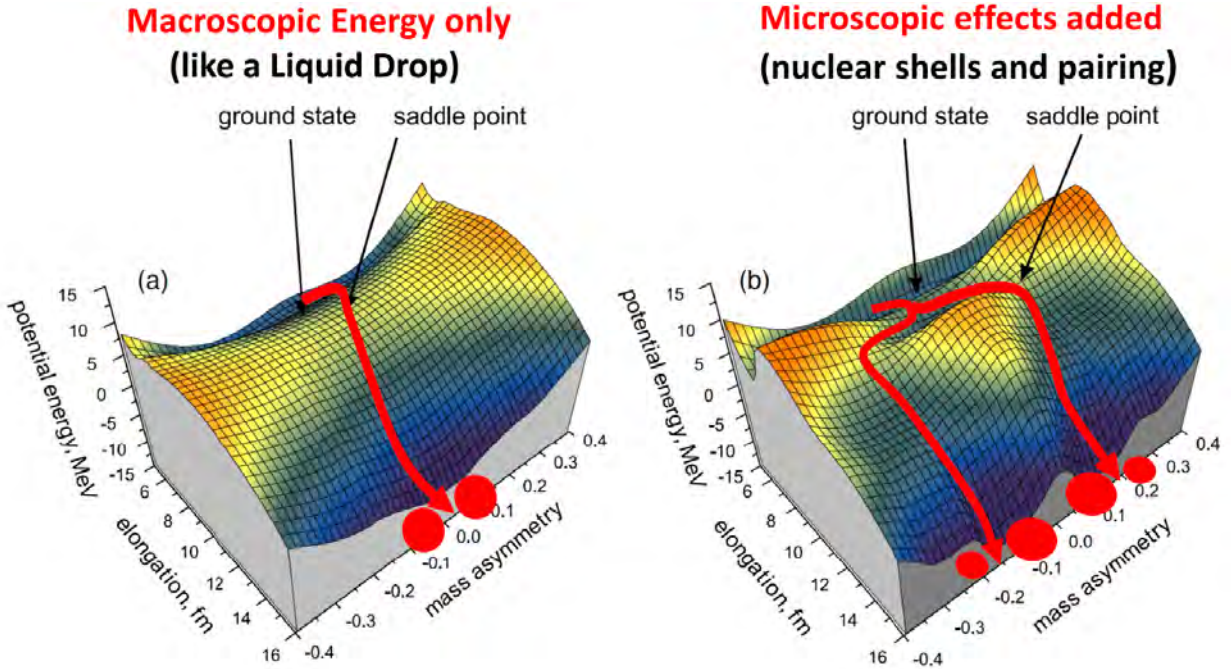


Figure 2.5 Plotted is the potential energy surfaces (PES) for ^{238}U in two dimensions (elongation and mass asymmetry). The left plot shows the PES in the liquid drop model. The right plot shows the PES from the liquid drop model plus microscopic (shell) effects. The most probable fission paths are shown by the red lines with arrows and follow the lowest energy pathway. Figure from [6].

The main driving force behind the asymmetric fission of actinides like ^{238}U is understood to be from shell effects, with one of the daughter products favoring the availability of the doubly magic ^{132}Sn with a closed shell of both protons and neutrons. Until recently, it was expected that lighter nuclei below the actinide region would favor symmetric fission. A study in 2010 observed that the fission fragment mass distributions of ^{180}Hg were peaked near ^{80}Kr and ^{100}Ru , despite the availability for a symmetric split into two ^{90}Zr fragments which would have a closed shell of neutrons [23]. This surprising observation sparked discussion among theorists to find an explanation for this new mode of asymmetric fission.

2.2 Experiment Proposal

This section explains the feasibility of measuring fission in exotic nuclei at the NSCL, and the isotopic region of study.

2.2.1 Recent Developments

An interesting opportunity to explore this question became available late in the operational lifetime of the Coupled cyclotron facility, when it became feasible to propose experiments using secondary beams that could be produced at the MSU Coupled Cyclotron Facility with a ^{208}Pb primary beam. This allowed, for the first time, the production of secondary beams that could enable fission measurements near the ^{180}Hg region. The alternative possibility of earlier experiments studying fission via a deep inelastic entrance channel was considered, but rejected because fission production via the heavy-ion deep inelastic scattering entrance channel is irrelevant to our current understanding the astrophysical r-process.

2.2.2 Location in Chart of Nuclides

Motivated by the recent observation of asymmetric fission in ^{180}Hg , our goal was to study isotopes near the transition region between observed symmetric and asymmetric fission. Figure 2.6 shows theoretical calculations for the expected mass distributions of isotopes with $74 \leq Z \leq 94$ and $90 \leq N \leq 150$, where the color scale denotes expected symmetric (blue) and asymmetric (red) yields. The black oval near the center indicates the potential region of study for our NSCL fission experiment.

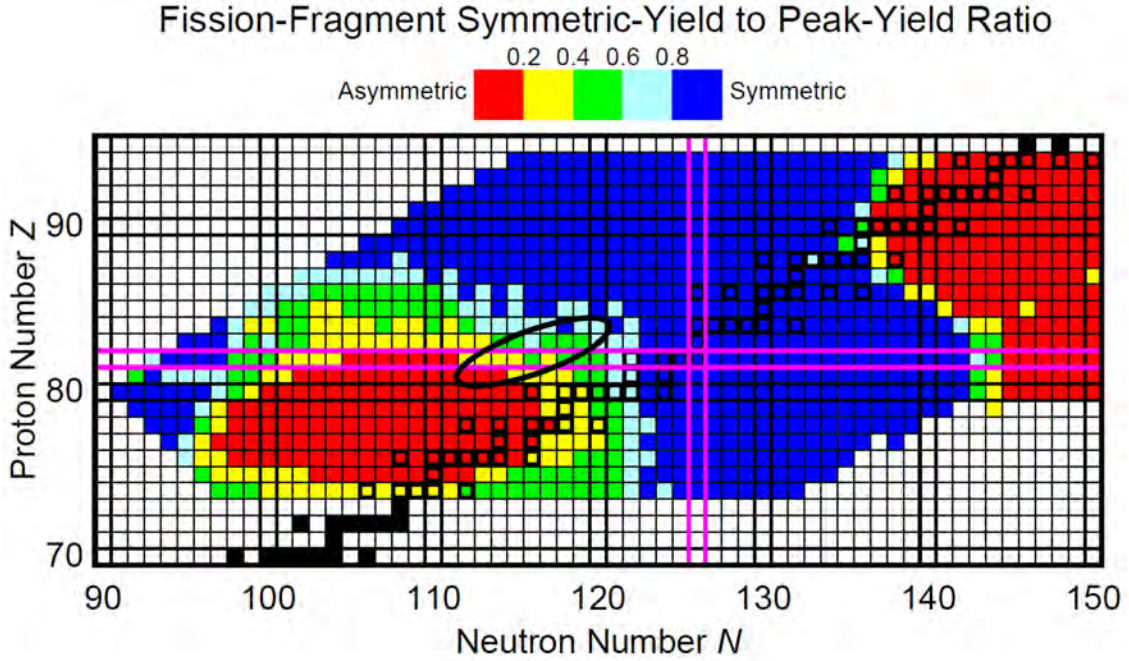


Figure 2.6 Theoretical fission yield calculations for isotopes with $74 \leq Z \leq 94$ and $90 \leq N \leq 150$. The black oval near the center shows the region of study for our fission experiment. Figure from [7].

2.2.3 Novelty of Study and Challenges to Solve

To obtain information about the barriers to fission, the excitation energies of the formed CN should ideally be as low as possible. Heavy beams are often studied at places like GSI [24] at energies ≥ 200 AMeV, which are well above the typical fission barrier. This experiment utilizes fusion-induced fission via inverse kinematics of a RIB in the neutron deficient Pb region at energies around 70 AMeV with helium gas, which is then further degraded to lower energies through added material in the beam path and/or energy deposition with the gas before fusion. As a result, we will be able to provide experimental measurements to compare with of the calculated fragment yields in Figure 2.6, and obtain fission barriers.

Well after the experiment was successfully proposed, and just before this experiment was supposed to run, the NSCL beam development group deemed the task of identifying such heavy RIB isotopes at such low energies to be too difficult with their instrumentation. This task had been accomplished before, at higher beam energies where the beam ions are completely stripped of their

electrons. To achieve particle identification (PID), the HEavy ISotope Tagging (HEIST) detector system system was designed specifically for resolving heavy nuclei that are not fully stripped [25]. This involved the development and construction of an ionization chamber capable of sufficient energy resolution for this task, which will be discussed further in Section 3.2, as well as a way to estimate the contamination from isotopic charge states which will be discussed in Section 4.2.

CHAPTER 3

EXPERIMENTAL SETUP

This Chapter explains the accelerator facility used to create the RIB used in the fission experiment, as well as the detectors employed to make the measurements.

3.1 The NSCL

The fission experiment E12014 was conducted in 2020 at the National Superconducting Cyclotron Laboratory (NSCL), since decommissioned and replaced by the Facility for Rare Isotope Beams (FRIB). The NSCL as a whole had many diverse capabilities to support a vast range of experimental studies; here I will only focus on the relevant components to this thesis. The layout of all beam lines can be seen in Fig 3.1; we utilized only a very small fraction of the facility.

In this section, I will discuss the role of the NSCL in producing the radioactive beams used in E12014. Namely, the acceleration of a stable primary beam into a production target, followed by tuning of the resulting cocktail of radioactive isotopes.

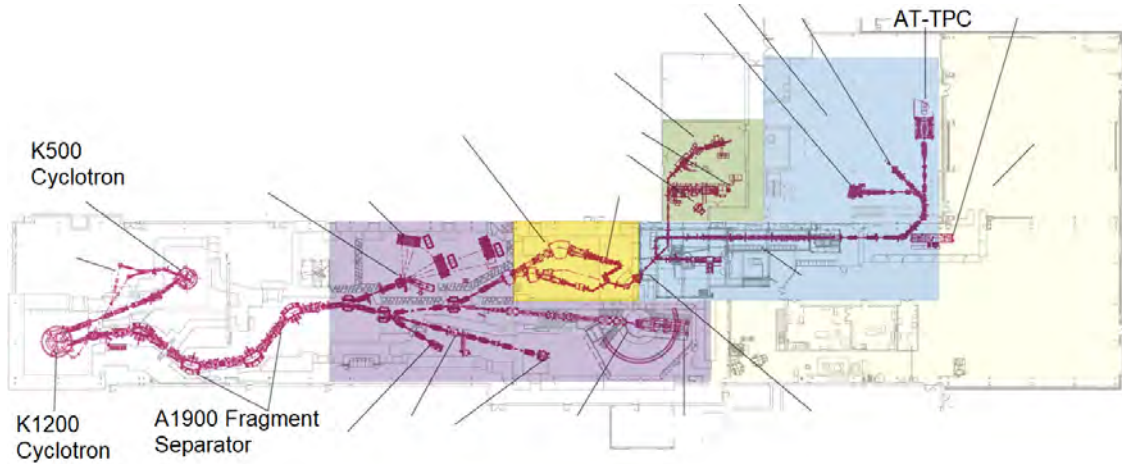


Figure 3.1 Layout of the NSCL facility beam lines, with only the components relevant to E12014 labeled. Modified from [8].

3.1.1 Radioactive Isotope Beam Production

Beam production begins in the K500 and K1200 cyclotrons at the NSCL, where a primary beam of the stable isotope ^{208}Pb is accelerated and stripped of electrons. Ions of Pb ($Z=82$) that have been completely stripped of electrons are denoted by the *charge state* label ' $82+$ '; Pb ions with a single

remaining electron are denoted ‘81+’, and so on. A useful metric when discussing projectiles in a particle accelerator is their resistance to deflection by a magnetic field, or their *magnetic rigidity* ($B\rho$):

$$B\rho = p/q \quad , \quad (3.1)$$

where p is the projectile momentum and q is its charge.

A list of the stable beams employed in E12014 is given in Table 3.1. While $^{208}\text{Pb}^{82+}$ was used to produce the radioactive isotopes, the other charge states $^{208}\text{Pb}^{63+,79+,80+,81+}$ were utilized to test and calibrate our detector systems.

Once the beam is accelerated to the desired production energy, it is sent to collide with a 22.5 mg/cm² target of ^9Be , resulting in the creation of many (≈ 100) unstable isotopes. For the purpose of extracting physics measurements, it is useful to select a subset of these isotopes near the area of interest within the chart of nucliides (≈ 30). This is done in practice by appropriately tuning the A1900 fragment separator. Beam particles are tuned via focusing magnets, and restrictions are placed on their momenta using an aperture slit.

Beam type	$B\rho$ (Tm)	Charge states
Pilot beam	4.4785	63 ⁺
Calibration beam	2.5782	79 ⁺ , 80 ⁺ , 81 ⁺ , 82 ⁺

Table 3.1 A list of the stable Pb beams used in E12014 for calibration purposes.

3.1.2 A1900 Fragment Separator

During its operation, the A1900 fragment separator has been used for multiple purposes, e.g. 1) such as magnetic spectrography or 2) as a monochromator for primary beams. In the case of E12014, however, it served in its 3d and main role as an achromatic separator for radioactive secondary beams. Figure 3.2 shows a schematic diagram of the A1900 layout in this mode. Four superconducting dipole magnets (large “D” shapes) provide the main bending power need to select the rare isotope beam and direct it through 24 superconducting quadrupole magnets (smaller boxes) to the reaction target. The beam momentum was defined by magnetic fields of the dipoles and the location of aperture slit located at Image2 (I2) in the A1900 separator.

Unstable beam particles resulting from collisions of the primary beam with the production target are sent directly to the A1900. The initial $B\rho$ acceptance is such that any unreacted primary ^{208}Pb beam is not transmitted. The aperture slit at I2 is set to a gap of 2 mm, which selects projectiles with $B\rho = 3.271 \text{ Tm}$. This significantly reduces the number of isotopes that make it to our detectors for study, from the initial hundreds down to the desired isotopes for our experiment.

At this stage along the beam line, the identity of each isotope is still unknown despite the strict selections that have been applied to momentum. Isotopes that survive tuning in the A1900 will then reach the first of our detectors specific to E12014.

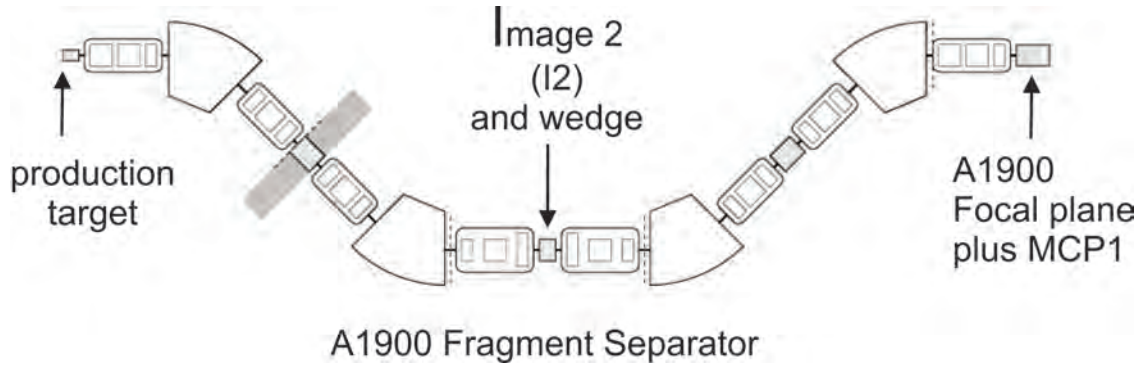


Figure 3.2 Schematic of the dipole and quadrupole magnets in the A1900 fragment separator.

3.2 HEIST Detector System

The HEavy ISotope Tagger (HEIST) system is a collection of detectors that were designed and manufactured specifically to combat some of the unique challenges presented when identifying radioactive isotope beams in inverse kinematic experiments [25]. There are three main detector systems tasked with (1) providing particle time of flight (ToF), (2) beam tracking and energy deposition (ΔE), and (3) collection of gamma rays emitted from isomer states. This method of particle identification, in addition to selections made during beam tuning, is commonly referred to as the $\Delta E - \text{ToF} - B\rho$ method. The MUSIC provides ΔE , the start and stop MCP provide ToF , and the momentum cuts made in the A1900 fragment separator give $B\rho$. Collection of isomer gamma rays provides another tool for isotope identification.

ToF is measured by a start and stop micro channel plate (MCP) detector. The start MCP is placed just after the A1900 fragment separator, and the stop MCP is located approximately 32 m

further down the beam line, just in front of the other HEIST detectors. The difference in the time signals recorded by the two MCP's gives an uncalibrated preliminary ToF value for each beam particle, which is then shifted by a constant to obtain the true ToF in ns.

Energy deposition and basic particle tracking are provided by the multi-sampling ionization chamber (MUSIC). This detector is located after the stop MCP, and contains an active volume of gas which the beam particles will deposit a small amount of their energy in through ionizing electrons. The charge is collected over several rectangular pads, with triangular pads at the beginning and end to facilitate particle tracking.

After the MUSIC detector lies an aluminum beam stop, which can be inserted into the beam line or retracted to allow for transmission to the final detector. When inserted, the beam will be fully stopped next to a high purity germanium (HPGe) crystal detector that collects gamma rays emitted by the stopped fragments. These gamma rays are essentially unique to each isotope, and so can be used for exact identification.

An overview of these detectors can be seen in Figure 3.3, where the radioactive isotopes exiting tuning from the A1900 enter from the left, and travel towards the right to sequentially pass through each of the HEIST detector systems. The next few subsections will go into further detail about the operation of each.

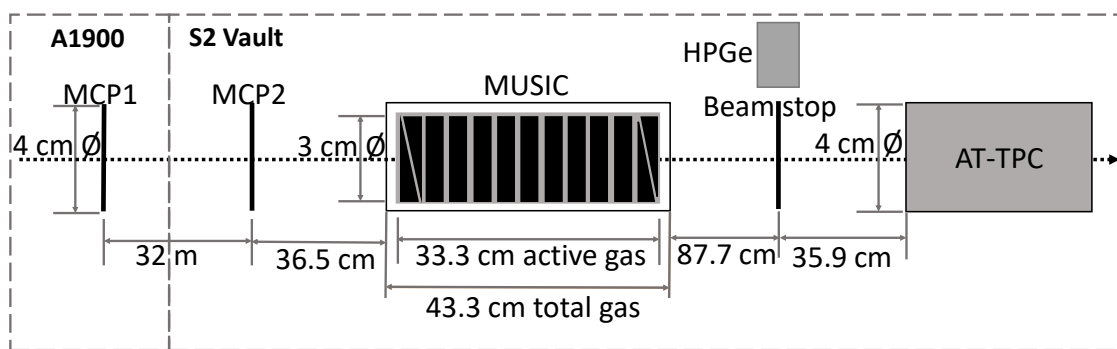


Figure 3.3 The complete HEIST detector system.

3.2.1 MCP

A diagram of the operating principles behind the MCP detector is shown in Figure 3.4. Beam particles pass through a thin emitting polypropylene foil, producing electrons that travel under the

influence of an $\mathbf{E} \times \mathbf{B}$ field to be collected on the MCP anode. The $\mathbf{E}$ field is created by biasing copper plates, and the magnetic field is supplied by fixing small permanent neodymium magnets on the top and bottom plates of the MCP housing.

Beam rates in E12014 can potentially reach 10^4 , so a counter is needed which can ideally support even higher rates. Typical fast counting detectors like a plastic scintillator would not survive the heavy secondary beam. Because the physical detector lies outside the beam path in an MCP, the only limiting factor is the timing resolution. This has the additional benefit of negligible changes to beam energy and transmission, since the polypropylene foil is basically invisible to the beam. Thus, the MCP provides a convenient balance between high rate capabilities with minimal insertion of material into the beam path.

Figure 3.5 shows the ToF spectra for the primary beams utilized in E12014, which are listed in Table 3.1. The top panel shows our pilot beam of $^{208}\text{Pb}^{63+}$ with $B\rho = 4.4785$ Tm, which is the beam initially used by the tuning operators to roughly center the beam focus. The scale on the x-axis has been adjusted from the raw time difference to represent the actual elapsed time between the two MCPs. Calculations from the LISE++ software [26] are used to predict the expected ToF for the pilot beam, and an adjustment offset is applied to match the raw spectra. The bottom panel shows several charge states for the primary ^{208}Pb beam with $B\rho = 2.5782$ Tm, which provide further calibration points for the ToF spectra.

A small detail to take note of here that will be important for the secondary beam, is that neighboring charge states of a given isotope will differ in their ToF by only a few ns. With only the primary beam present, it is simple to distinguish the different charge states of the beam. When the multiple species of the secondary beam are present, this requires more measurements to disentangle.

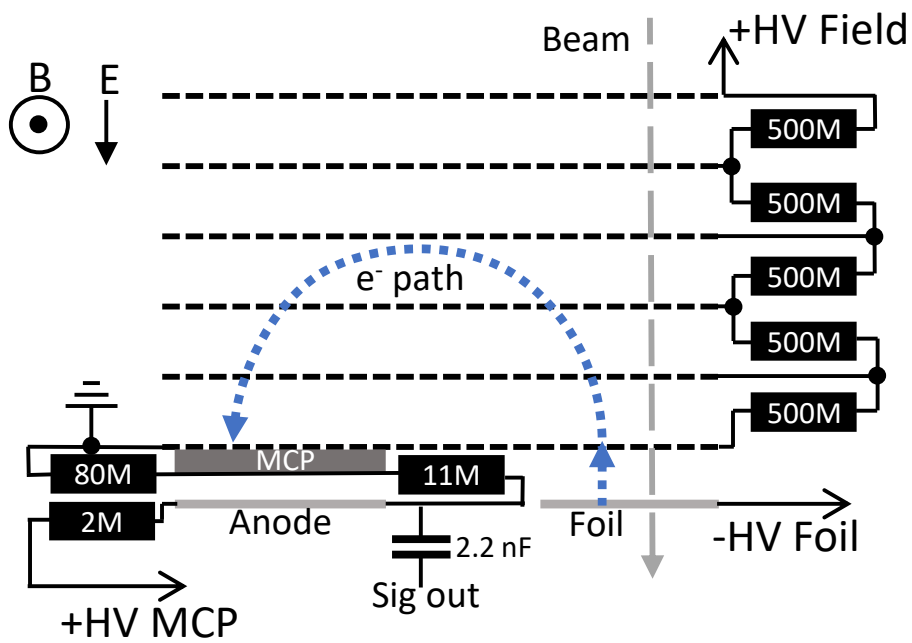


Figure 3.4 An schematic illustration of the MCP detector.

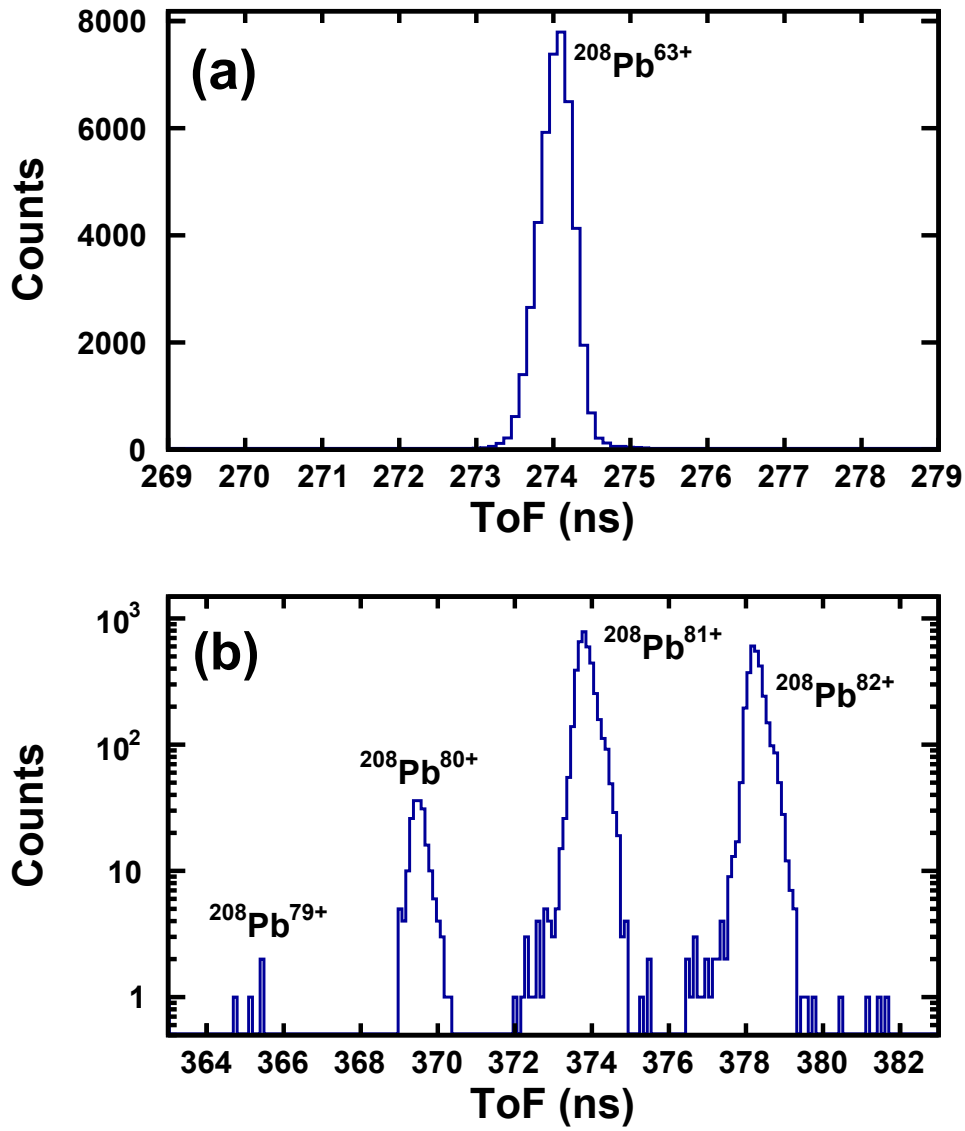


Figure 3.5 ToF spectra from the MCP detectors for the primary and calibration stable beams.

3.2.2 MUSIC

After passing through the stop MCP, beam particles will next enter the MUSIC detector. In general, an ionization chamber is a volume of detecting gas held to constant pressure through which charged particles can traverse and deposit their energy by liberating electrons from the detecting gas. These liberated electrons are drifted orthogonally from the beam path to a collection anode by the application of a uniform electric field within the volume, enforced by a field cage. The amount of collected electrons is proportional to the energy deposited by the charged particle.

In the case of E12014, it was necessary to develop an ionization chamber capable of resolving the typical energy deposition of a heavy radioactive isotope beam. At other facilities (GSI, RIBF) heavy radioactive secondary beams are usually fully stripped of their electrons. Here at the NSCL, the energies are such that most ions will retain their 1s shell of electrons ($Q/e = Z-2$), with a smaller but still relevant probability for the neighboring $Z-1$ and $Z-3$ charge states to be populated. Furthermore, in the energy regime of our study, the velocity distributions of these heavy isotopes and their charge states can overlap.

An operating schematic of the MUSIC detector can be seen in the top panel of Figure 3.6, and the bottom panel shows a cross-sectional inside view from during construction. There are two main advancements to the standard ionization chamber design: the use of a segmented anode plane, and the presence of a *Frisch* grid wire plane.

The 9 rectangular segments of the anode provide finer energy resolution than would a single pad, and the 4 triangular pads at the entrance and exit of the active volume allow for primitive particle tracking. For beam particles not exactly centered, more charge will be deposited on one of the triangular pads. The ratio of charges on these pads can be converted to a horizontal position within the MUSIC. Furthermore, vertical position information can be inferred from the known drift properties of electrons in the detection gas under the applied $\mathbf{E}$ field.

For the experiment, CF_4 was chosen as the detection gas for MUSIC. The advantages of this particular gas choice are that CF_4 is non-explosive, has a fast drift velocity, and a large stopping power. These characteristics are good for safety, minimize pile-up effects, and provide solid energy deposition resolution, respectively. Information about the detection gas settings used for MUSIC in this experiment can be found in Table 3.2.

Table 3.2 Selected gas properties of CF_4 at running conditions.

Gas (torr)	Cathode voltage (V)	E_C (V/cm)	e^- drift velocity (cm/ μs)	Avg energy per ion pair (eV)
CF_4 150	-1000	192	10.3	34.3
CF_4 300	-1000	192	8.3	34.3

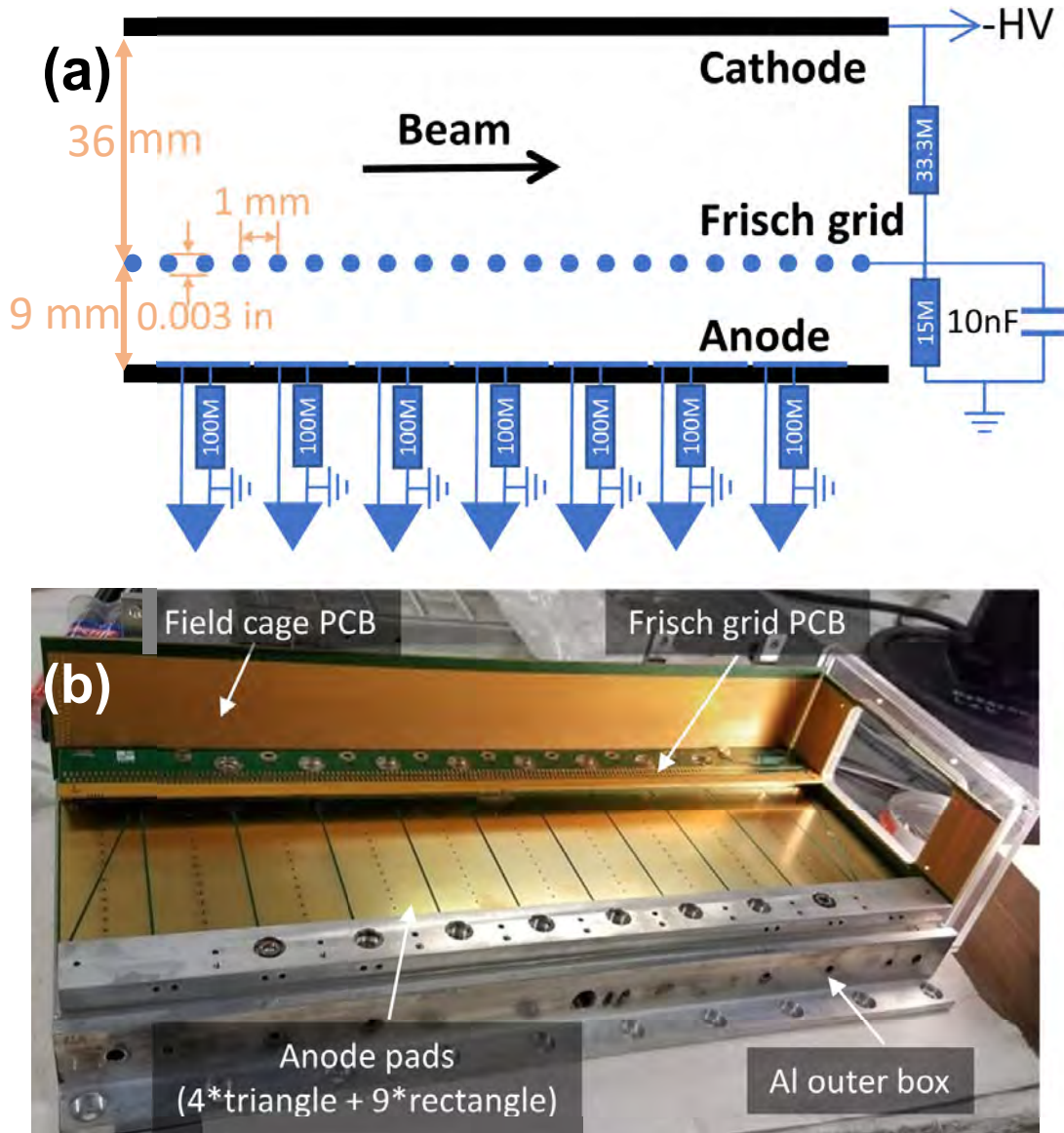


Figure 3.6 (TOP) An illustration of the operating principles of MUSIC. (Bottom) The inside of the MUSIC detector, before complete assembly.

3.2.3 Constructing Music

Figures 3.7 and 3.8 show pictures taken at two stages of the MUSIC construction. Paramount to the energy resolution of the detector are the requirements of even spacing between the wires in the Frisch grid, and the levelness of the anode plane. The energy resolution is discussed further in Section 4.2.1.



Figure 3.7 Picture taken during the gluing process. The anode is held level on a vacuum plate in preparation, while epoxy is applied to the aluminum base plate by the author of this dissertation. The vacuum plate was then lowered onto the bas plate to allow the epoxy to bond.

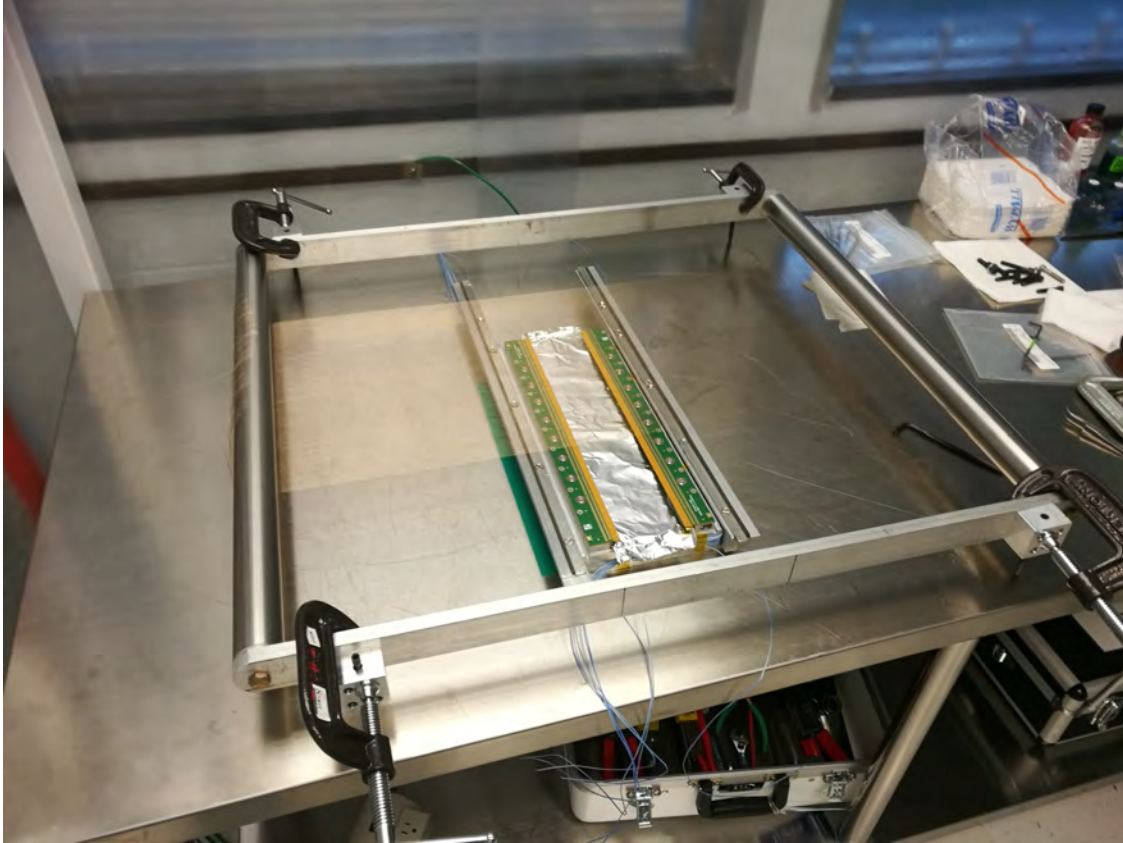


Figure 3.8 A Picture showing the alignment of the anode wire plane before soldering. The wires are wound around the support bars using a step motor to achieve uniform 1 mm spacing.

3.2.4 HPGe

When the aluminum beam stop is inserted, gamma rays from the fragments will be collected by the high purity germanium (HPGe) crystal detector. The 80.8 mm $\times$ 105.8 mm crystal is cooled by a liquid nitrogen dewar and connected an ORTEC 120 preamplifier through a high voltage filter. FWHM resolution of 1.1 keV was achieved over the range of measured gamma rays, from 300 - 1000 keV. The energy scale of the HPGe was calibrated using the known emission spectra of a ^{125}Sb source.



Figure 3.9 An external view of the HPGe detector. The liquid nitrogen Dewar can be seen.

These gamma rays come from the decay of long lived excited nuclear states called *isomers*, which are populated during acceleration, transmission, and tuning. Since the typical flight time of the radioactive isotopes is around 300 ns, only isomers with lifetimes of $\leq 10^2$ ns can be efficiently detected. This sensitivity to isomer lifetime means that a stronger gamma ray spectra will be obtained from some isotopes than others.

Because the isomer states depend on shell effects within the nucleus, the energy and lifetime of the gamma-ray from an isomer decay will be unique to that nucleus and can be used to associate the unique arrival times with specific isotopes in the beam. Once that arrival time for a given isotope is established, all future arrivals at that time can be assigned to that specific isotope.

3.3 The AT-TPC

Finally, after traversing the A1900 fragment separator and all of the HEIST detectors, the beam particles will arrive at the active target time projection chamber (AT-TPC) where they will have the chance to undergo fusion-fission. In general, a time projection chamber (TPC) provides a method of 3-D spatial reconstruction for charged particle tracking. The AT-TPC is an extension of the standard TPC wherein the detection volume also serves as the target. In this section I will discuss the operating principles and setup of the AT-TPC as it pertains to E12014.

3.3.1 Operating Principles

A schematic diagram of the AT-TPC is shown in Figure 3.10, along with the coordinate system that will be used when discussing geometry. The beam, traveling in the negative $\hat{z}$ direction, enters the detection volume through the window and some foils. In events of interest, fusion-fission occurs where the beam particle fuses with an atom of the target gas to form a compound nucleus (CN) which then promptly fissions into two charged daughter products. Both the initial beam particle and the charged products will ionize the target gas as they travel the active volume, creating free electrons just as in the MUSIC. An applied electric field will drift these electrons parallel to the beam direction towards a highly segmented readout anode.

2-D information about particle tracking in the x-y plane is inferred from which pads recorded charge during an event. The center of the anode plane consists of a hexagonal region of small triangular pads with height 0.5 cm, surrounded by larger triangular pads with height 1 cm with a total pad count of 10240 [27]. This provides a finer spatial resolution for tracks in the small pads. Tracks are then “unfolded” in the $\hat{z}$ -direction using the known drift properties of the electrons under the applied $\mathbf{E}$ -field of 700 V/cm. General Electronics for TPCs (GET) is used to operate the 10240 channels associated with the pads; more info about GET can be found in [28].

For E12014, helium gas was pumped in and cycled at a pressure of 800 atm (slightly over atmosphere) in order to avoid inflow contamination inside the detection volume. As unreacted beam particles travel further into the 1 m long active volume, they will continuously lose energy through ionization of helium electrons. This means that the excitation energy of a fission event depends on the depth where the CN is formed. Sampling events with vertices at different locations within the AT-TPC conveniently allows for sampling of the fission barrier.

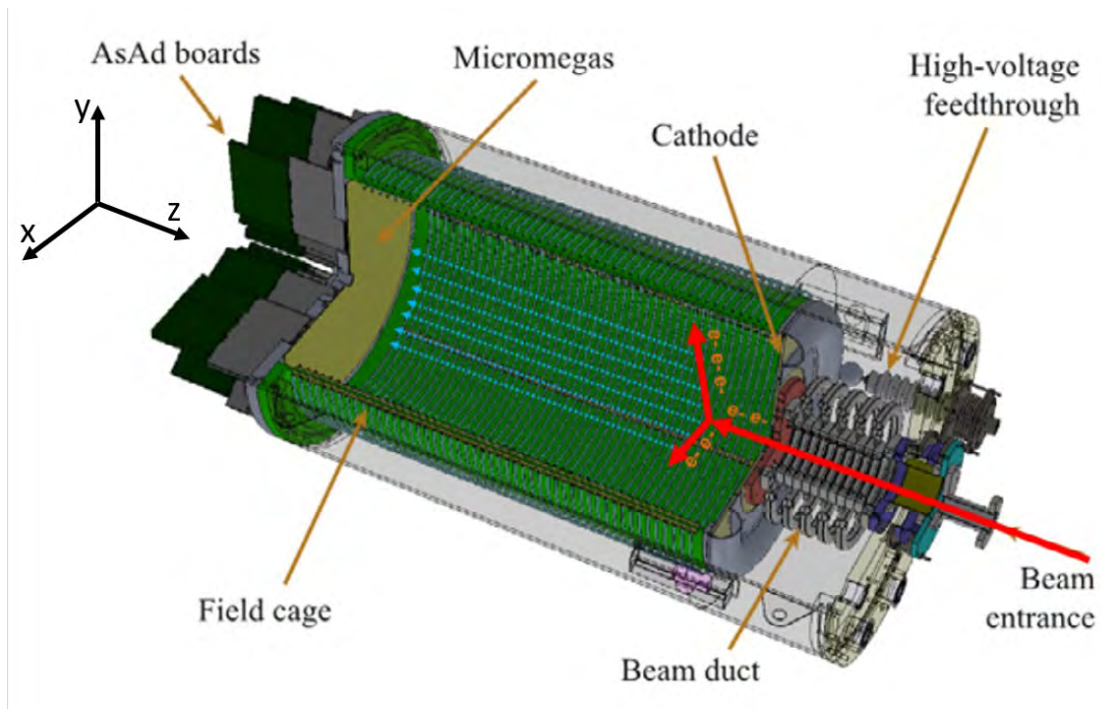


Figure 3.10 Diagram of the AT-TPC with labeled parts.

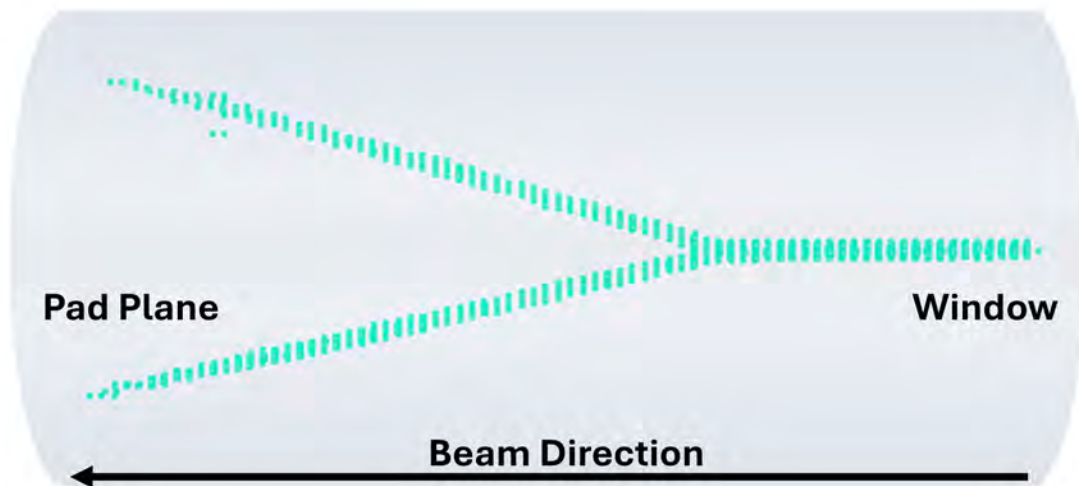


Figure 3.11 Terminology and orientation within the TPC to be used for reference later.

3.3.2 Smart-Zap Region

Of course, we are only interested in events where the beam particle forms a CN and undergoes fission, which occurs very rarely. The vast majority of beam fragments will simply pass through

the AT-TPC. To facilitate recording only events of interest, a selection of anode pads in a 50 mm diameter circular area around the beam spot, called the “*smart-zap*” region, were operated at a much lower gain. These pads are noted by the red dots in Figure 3.12. This way, the unreacted beam particles on their own don’t satisfy the trigger threshold conditions, and so will not be recorded. If a CN is formed and fission takes place, the idea is that the fission products will create tracks which travel outside of the smart zap region, over pads at nominal gain values that will meet trigger thresholds.

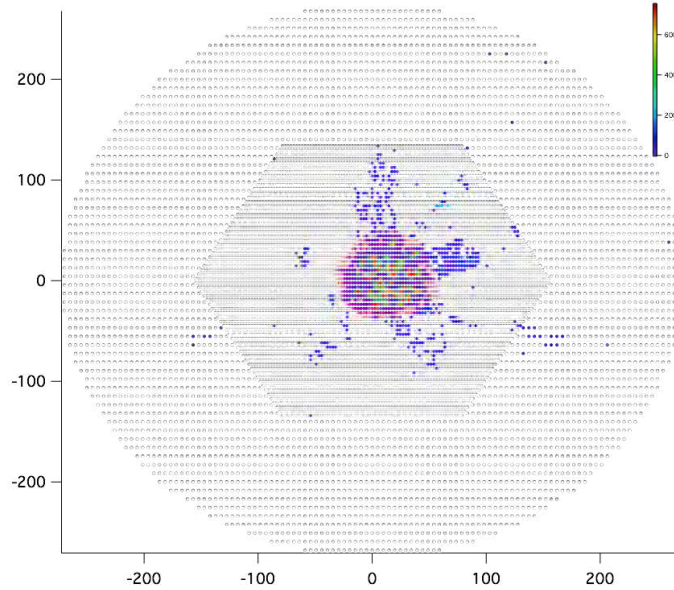


Figure 3.12 A view of the AT-TPC pad plane where the smart-zap region of pads are shown by the red dots in the very center. The hexagonal area of small pads in the anode and surrounding region of larger pads are also visible.

In practice, many unreacted beam particle events were prevented from triggering the data acquisition (DAQ), consistent with this concept. However, we still recorded a large number of non-fission events due to electronic noise and multiply scattered beam particles that nevertheless escaped the smart zap region. A sequential condition algorithm was developed to sort the events into categories that was later supplemented by a machine learning routine, which will be discussed in more detail in Section 4.3.

CHAPTER 4

ANALYSIS

This chapter is dedicated to detailing the steps of analysis required to obtain the desired fission observables in E12014, including identification of the pre-fission isotopes in the RIB, measurements of species purity in the data, performance of the HEIST detector system, classification sorting of fission events, and the development of a space charge correction procedure for space charge effects.

4.1 Beam Identification

Before the interesting fission events can be properly analyzed, beam fragments must be tagged in particle identification (PID). Due to the nature of how the radioactive secondary beam is created, it is not possible to know precisely which isotopes are created and transmitted down the beam line. That task is precisely the reason that HEIST was developed, and in the following subsections I will detail how the different detector systems work together.

4.1.1 PID From ΔE -ToF- $B\rho$

To separate out the isotopes by mass and charge, we construct a PID plot using the ToF information from the MCPs and the energy deposition from MUSIC. The ToF is accurate to the level of picoseconds, and is given by the difference between the start MCP and stop MCP timing. Calibration to absolute ToF is done by shifting the spectra for the pilot/calibration beams to match the expected results from a LISE++ calculation, based on the velocity of those beams at the given energies, and on the known distance of the beam line between the MCPs. Such resolution allows for neighboring isotopes of an element to be separated by mass within the energy regime of E12014.

Energy deposition inside the MUSIC obeys the Bethe-Bloch formula, which goes like

$$\Delta E \approx \frac{Z^2}{v^2} f(v) \quad , \quad (4.1)$$

where Z is the proton number, v is the particle velocity, and $f(v)$ is a slowly varying function of v . Selections on $B\rho$ in the A1900 essentially fix v for the particles that make it to our detector, so to simplify even further

$$\Delta E \approx Z^2 \quad . \quad (4.2)$$

For each isotope, filling a 2-D histogram with “ ΔE vs ToF” will result in the plot shown in Figure 4.1. Each of the peaks correspond to a separate isotope, where the color scale is related to the total yield. Elements are grouped within ΔE into Z-lines, where lighter isotopes travel slightly faster down the beam line and deposit slightly less energy. Within ToF, particles are grouped into lines of equal A/Q.

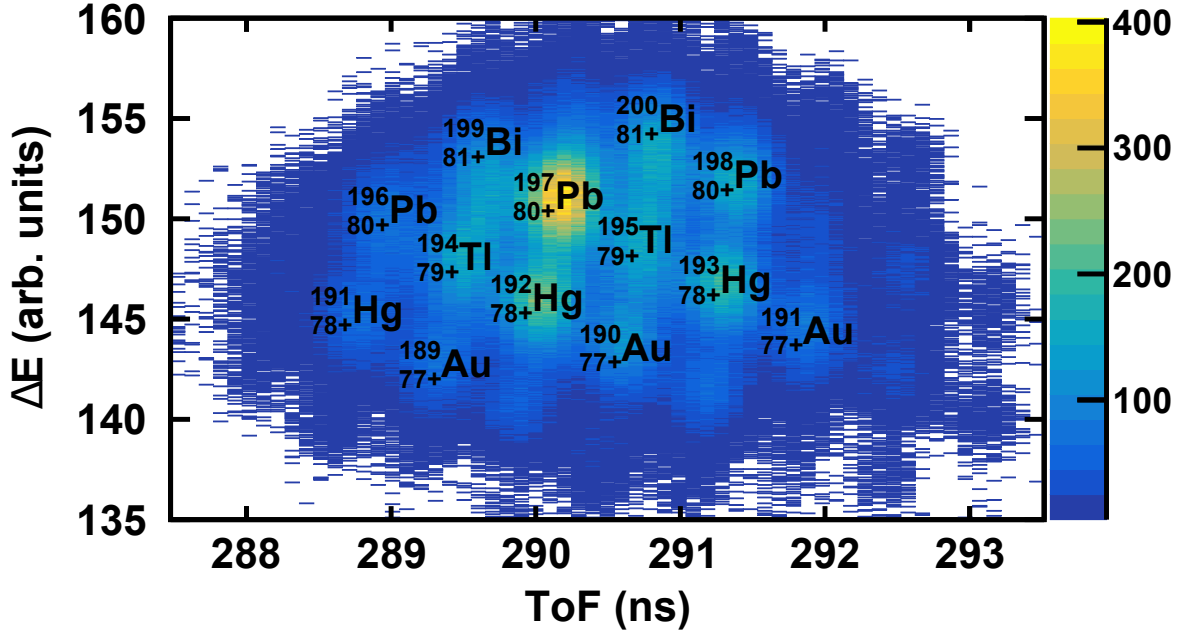


Figure 4.1 PID with labeled isotopes.

4.1.2 Isomer Tagging

The total ToF from the production target to the HEIST detectors is around 590 ns for the isotopes seen in Figure 4.1. The decay chain of isomers with long enough lifetimes to survive the trip down the beam line can be measured by the HPGe detector when the aluminum beam stop is inserted. By placing a restrictive gate on the PID plot around one of the isotopes, the corresponding γ -ray spectra will show a unique collection of peaks that each relate to a step in an isomer decay chain. By fitting the energies of these peaks, an unknown isotope in our PID can be cross referenced with the National Nuclear Data Center (NNDC) database for measured and expected nuclear decay chains [29] to get a positive identification.

Those isotopes which are more likely to occupy isomer states (that have compatible lifetimes

with our measurement scheme) will naturally have stronger peaks in the γ -ray spectra. A selection of such isotopes for E12014 are gated on, and their background-subtracted γ -ray spectra shown, in Figure 4.2. Solid white circles in the top panel indicate example gates on ΔE and ToF , with the resulting spectra in the bottom panel. Peaks matching decay chains found in the NNDC database are labeled with their energies; one can easily see how the isomers leave a unique fingerprint for each isotope.

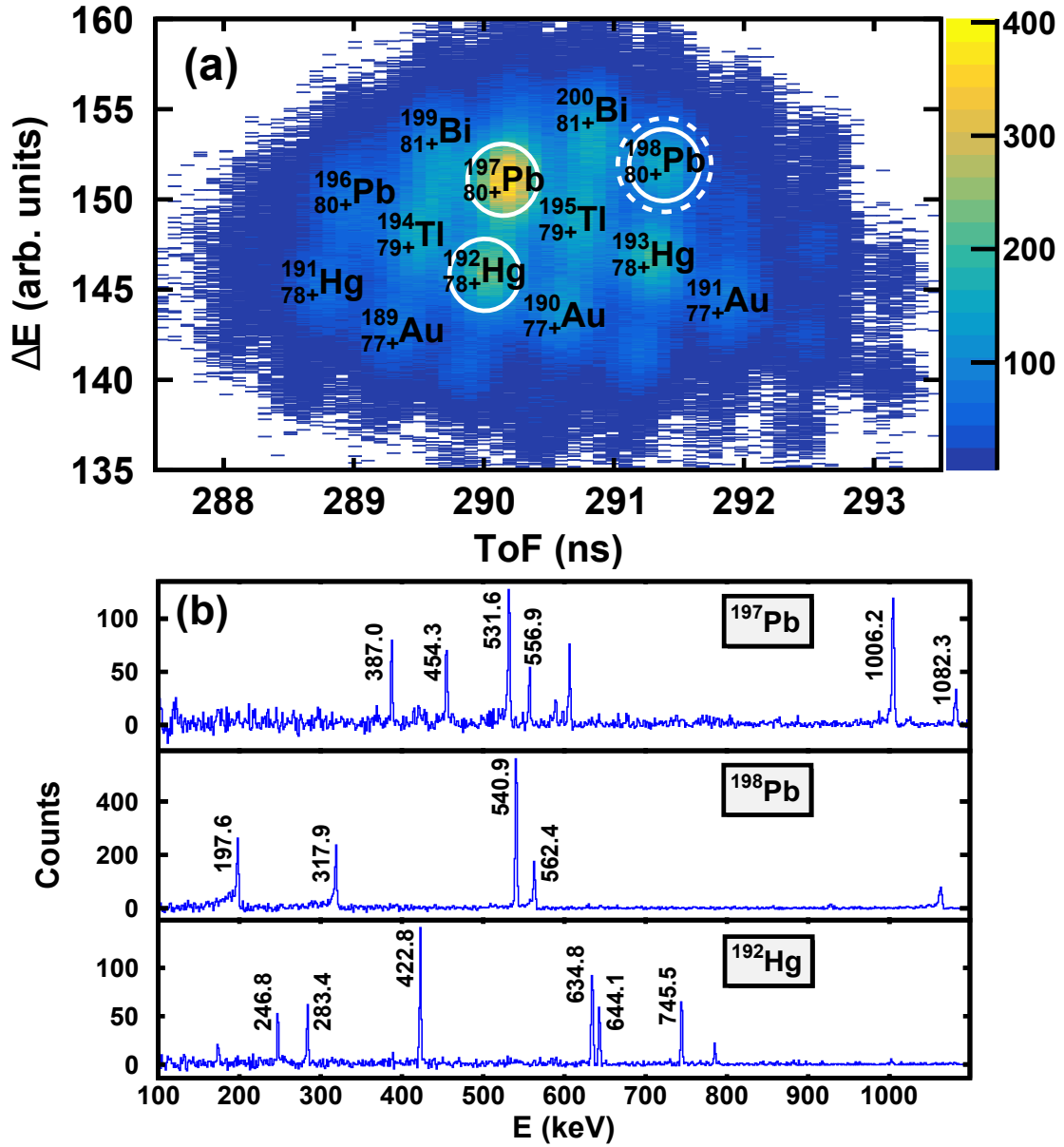


Figure 4.2 An example of the isomer tagging process.

Two important details should be emphasized in relation to isomer tagging. Firstly, is that in order to obtain γ -ray spectra, the beam is completely stopped in front of the HPGe. This means that no beam will be transmitted into the AT-TPC, and therefore no fission will take place. Isomer tagging can only be performed in this situation, and the resulting “map” is applied to standard data runs with full transmission. Secondly, not every isotope will have suitable isomers for which we can measure the decay chains. However, only a few of the isotopes in the PID need to be tagged before the rest can be known by counting. A list of the positively tagged isotopes is shown in Table 4.1.

Isomer-tagged Isotopes	Typical γ -rays (keV)	Charge States Identified
¹⁹⁶ Pb	337.3 , 371.9 , 689.0	79+ , 80+
¹⁹⁷ Pb	387.0 , 531.6 , 589.0	80+
¹⁹⁸ Pb	197.6 , 317.9 , 450.9	80+ , 81+
¹⁹² Hg	283.4 , 422.8 , 745.5	78+
¹⁹⁴ Hg	280.2 , 427.9 , 636.3	78+
¹⁹⁰ Au	232.3 , 854.9 , 427.8	77+
¹⁹¹ Au	264.0 , 420.1 , 725.2	77+ , 78+
¹⁹⁹ Bi	145.7 , 499.6 , 1002.2	81+
²⁰¹ Bi	185.8 , 411.7 , 967.4	81+

Table 4.1 A list of isotopes that have been confirmed by isomer tagging, their corresponding charge states measured in the PID, and some characteristic γ -rays seen in the HPGe.

4.2 Beam Purity

In this section, I will discuss the isotopic purity achievable through data selection for the HEIST system. Meaning that for a given gate on ΔE and ToF , one must determine which isotopes will be present and in what fractional percentage within the cut. This essentially relies on two main factors: the resolution of the MCP/MUSIC detectors in time/energy, and the abundance of charge states other than the helium-like configuration.

4.2.1 HEIST Resolution

At the energies available in the NSCL, the isotopes present in E12014 are not completely separated in ToF or ΔE as they might be at other facilities (GSI). This means that it is not possible to make a gate on the PID with sufficient statistics that includes only one species. The extent of

overlapping in ToF and ΔE can be seen in Figures 4.3 and 4.4, respectively. By taking thin slices indicated by the red dashed lines and projecting them to the energy or time axis, a multi-gaussian function can be fit and converted to a resolution for each of the isotopic peaks.

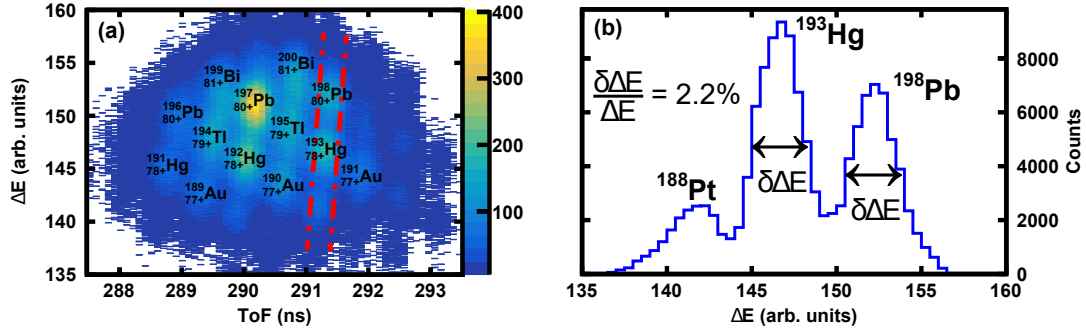


Figure 4.3 Resolution of the IC in energy deposition.

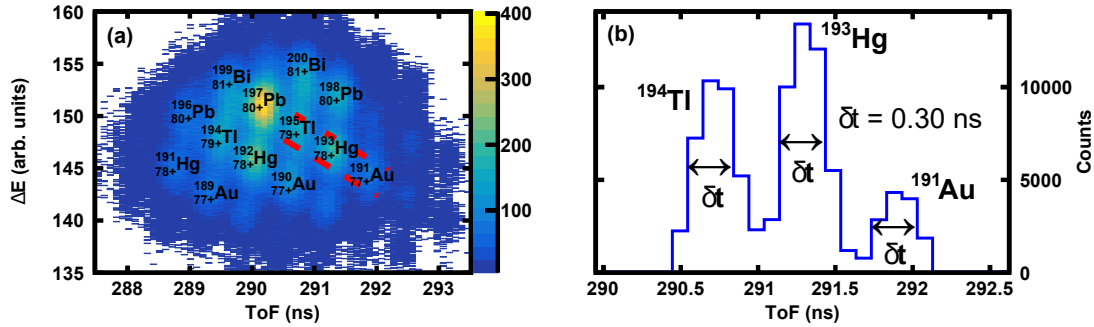


Figure 4.4 Resolution of the MCP in ToF.

The latent timing resolution of the MCP detectors is much finer than the separation of 0.30 ns FWHM achieved by HEIST in E12014 due to the overlapping velocity distributions of selected fragments. In theory, greater ToF separation in the PID could be achieved by placing the start and stop MCPs further apart. In practice, they were placed as far away from each other as possible, with approximately 32 m of open beam line between them. At any rate, the ToF separation achieved was sufficient for the experimental goals.

Specifically designed and manufactured for measuring the slow-moving heavy ions of interest in E12014, the MUSIC detector has a latent potential dE/E resolution of 0.85 %. The two main geometrical factors that influence this resolution are the size/spacing of the Frisch grid wires, and

the degree of uniformity in the spacing between the anode and cathode planes. When assembling MUSIC, efforts were made to control these two factors as much as possible. Vacuum plates were used for leveling the anode and cathode plates while they were epoxy-ed to the aluminum frame.

A level height gauge was used to measure the height discrepancy across pads, with the maximum height difference across plates found to be within $100\ \mu\text{m}$. Spacing for the Frisch grid wires was ensured by mechanically winding around a frame using fine stepping motors. The wires were soldered in place while attached to the frame, then cut free afterwards. More information about the construction process can be found in [17].

4.2.2 Charge State Contamination

As mentioned in previous sections, it is possible for other charge states to be populated than just the main helium-like (Z-2) configuration. In the HEIST detectors, these other charge states will deposit nearly the same energy in MUSIC but will enter with a notably different ToF. An example of this can be seen in Figure 4.5, where coincidence γ -ray gates have been placed on the HPGe data to select the ^{196}Pb isotope in the top panel. The main helium-like charge state is seen at 289 ns, with the lithium-like (Z-3) state coming in at 292 ns. In the full PID in the bottom panel, the Z-3 state of ^{196}Pb can potentially be a source of contamination in a gate on ^{198}Pb .

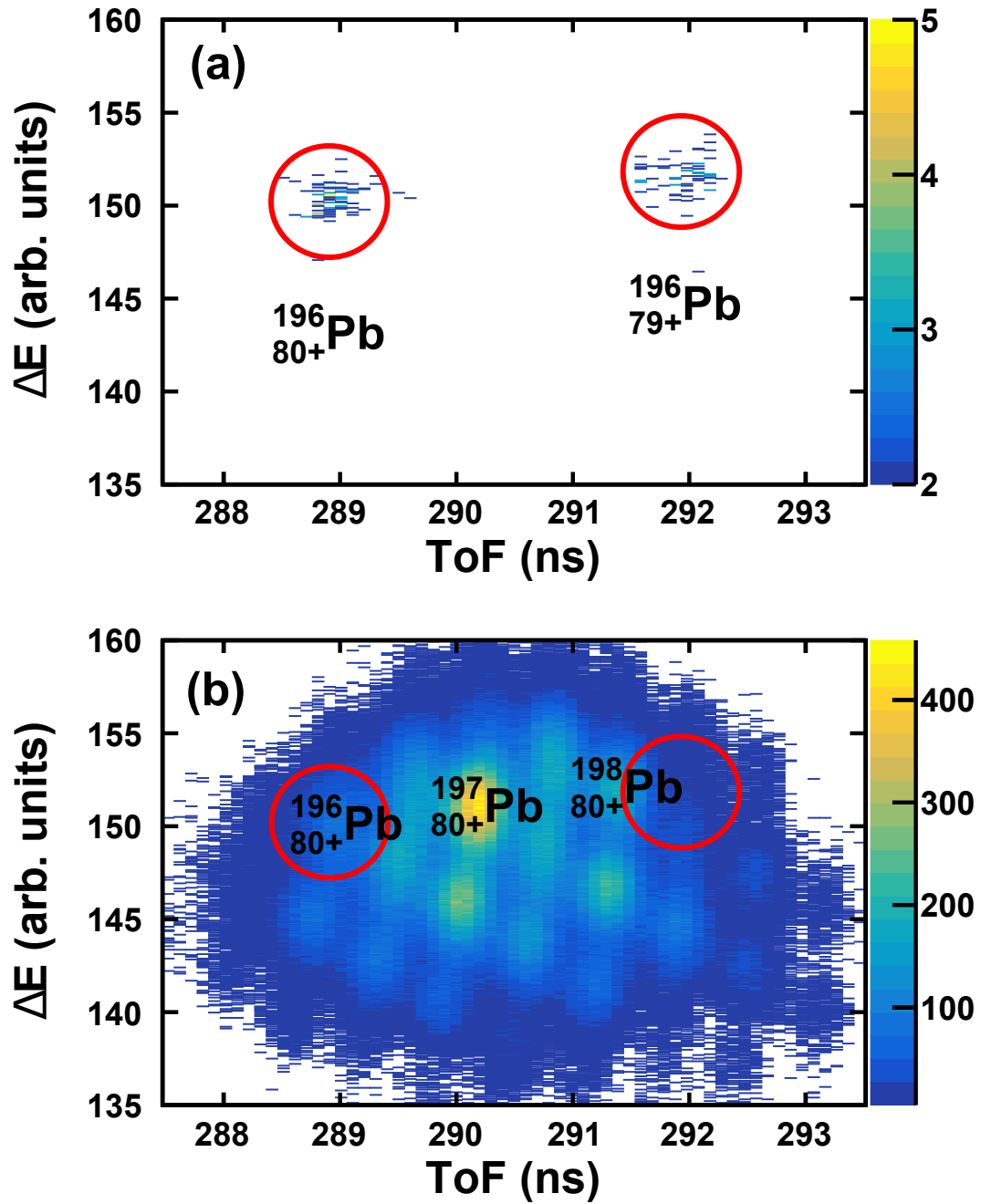


Figure 4.5 Example of charge state contamination as it appears in the PID.

One possible way to directly estimate the amount of charge state contamination in a given PID gate is by counting the number of characteristic γ -rays present in the cut for the suspected contaminate, compared to the numbers for a gate on the main charge state of that suspect. This method will obviously not work if the isotope in question doesn't have isomers that we can measure

in the HEIST set up. In this case, we could potentially rely on predictions from LISE++. There is limited experimental data for the region of study in E12014, such that the GLOBAL models in LISE++ cannot be immediately trusted in regards to yield information.

Focusing on the top panel of Figure 4.6, we see the predicted population of two charge states of ^{196}Pb as a function of $B\rho$, as calculated by LISE. The blue line represents the main helium-like configuration, and the red line shows the lithium-like. The vertical gray dashed line indicates the experimental $B\rho$ setting for E12014 of 3.271 Tm, for which LISE predicts the helium-like state should dominate. An experimental measurement obtained by counting γ -rays is shown by the red square and blue circle, which suggest nearly equal amounts of each charge state. By adjusting the effective $B\rho$ value in the LISE++ simulations to a best reproduction of 3.283 Tm, the observed distribution is recovered. The bottom panel shows a similar analysis for the ^{191}Au isotope.

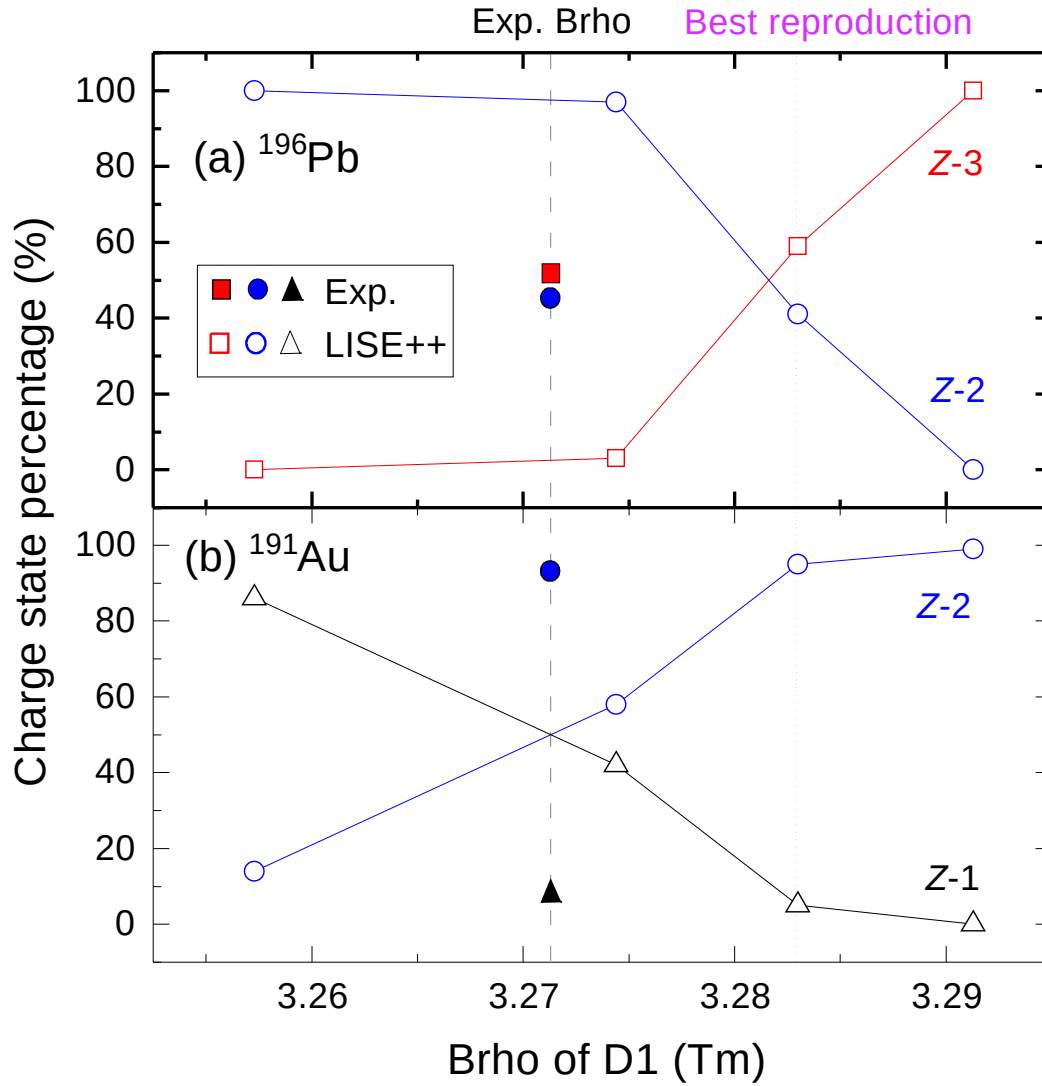


Figure 4.6 Expected population of charge states as a function of magnetic rigidity.

Table 4.2 show a summary of the measured charge state distributions for isotopes that have suitable isomers, compared to predictions from LISE++ using the best reproduction value of $B\rho = 3.283$ Tm. Some information about charge states for other isotopes than cannot be measured using the HPGe detectors may be inferred in this way. Simulations can also be performed to estimate where in the PID these other charge states would show up.

Isotopes	Calculation			Experiment		
	Z - 3	Z - 2	Z - 1	Z - 3	Z - 2	Z - 1
¹⁹⁶ Pb	59%	41%		52% ± 3%	48% ± 3%	
¹⁹⁷ Pb		100%			100%	
¹⁹⁸ Pb		100%			99% ± 2%	1% ± 1%
¹⁹¹ Hg	27%	73%		-	-	-
¹⁹² Hg		100%			100%	
¹⁹³ Hg		100%		-	-	-
¹⁸⁹ Au	2%	98%		-	-	-
¹⁹⁰ Au		100%			100%	
¹⁹¹ Au		95%	5%		95% ± 5%	5% ± 1%

Table 4.2 A comparison between the theoretical and measured charge state distributions of isotopes Pb, Hg, and Au with the charge state $Q/e = Z-3$, $Z-2$, and $Z-1$. The theoretical charge state distributions were calculated using the GLOBAL model with LISE++ at the adjusted $B\rho$ of 3.283 Tm. The experimental charge state distributions were extracted by gating on the corresponding γ -rays from known isomers.

4.3 Event Sorting

In this section, I will discuss the methods used to sort out the interesting fission events in the data from among the undesirable noise and unreacted beam events. While the smart-zap region of pads was successful to an extent in filtering out some of the unreacted beam events from the DAQ, there were still many instances where the trigger conditions were met despite the absence of a fission reaction. In fact, the vast majority of the recorded data is unreacted beam, with fission making up less than 1% of the total events.

4.3.1 Classification as Event Types

In order to efficiently sort through the data and find the events of interest, a labeling system was developed that is illustrated in Figure 4.7. Although clear fission is readily discernible by the human eye when viewing a 3D event reconstruction, it is not feasible to look through the hundreds of thousands of events recorded over the experimental campaign. Furthermore as was mentioned in Section 3.3.2, the trigger system designed to avoid recording unreacted beam events did not always succeed. Without it, orders of magnitude more unwanted data would have been recorded, but the cross section of fusion-fission is small to begin with.

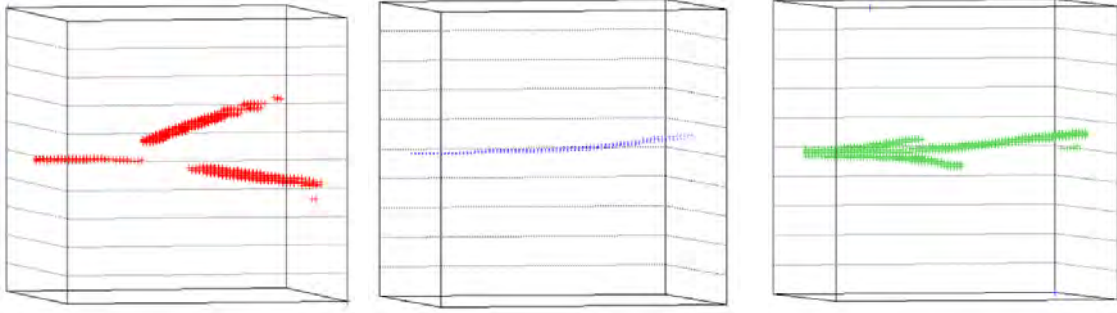


Figure 4.7 Examples of the three categories for sorting the recorded AT-TPC events. (Left) A clear example of fusion-fission. (Middle) An unreacted beam event that has strayed outside of the smart-zap region to pass the trigger conditions for detection. (Right) Likely a pile-up event where multiple RIB particles came through during one cycle of the DAQ. Clear distinction here between fission and unreacted beam is hard to make.

A sorting algorithm that relies on numerical characteristics unique to unreacted beam and fission was developed to classify the events according to Figure 4.7. An illustration of one such metric used in this algorithm is shown in Figure 4.8. Even with the forward boost to the fission fragments coming from the incident RIB velocity, there is still a significant folding angle in the laboratory frame. This folding angle manifests in the data as the spread of charge across the AT-TPC pad plane orthogonal to the z-axis.

Tracking will be described in Section 4.4.2, which involves fitting straight lines to a reconstructed event. Other metrics are involved in the classification algorithm which make use of tracking results, such as the total number of straight lines found and their beginning and ending location inside the AT-TPC. For example, a fission event should ideally be tracked into three separate lines, with one for the incoming beam and then two for the fragment products. An unreacted beam event will likely be tracked as a single line that starts near the entrance window and ends at the far side of the detection volume.

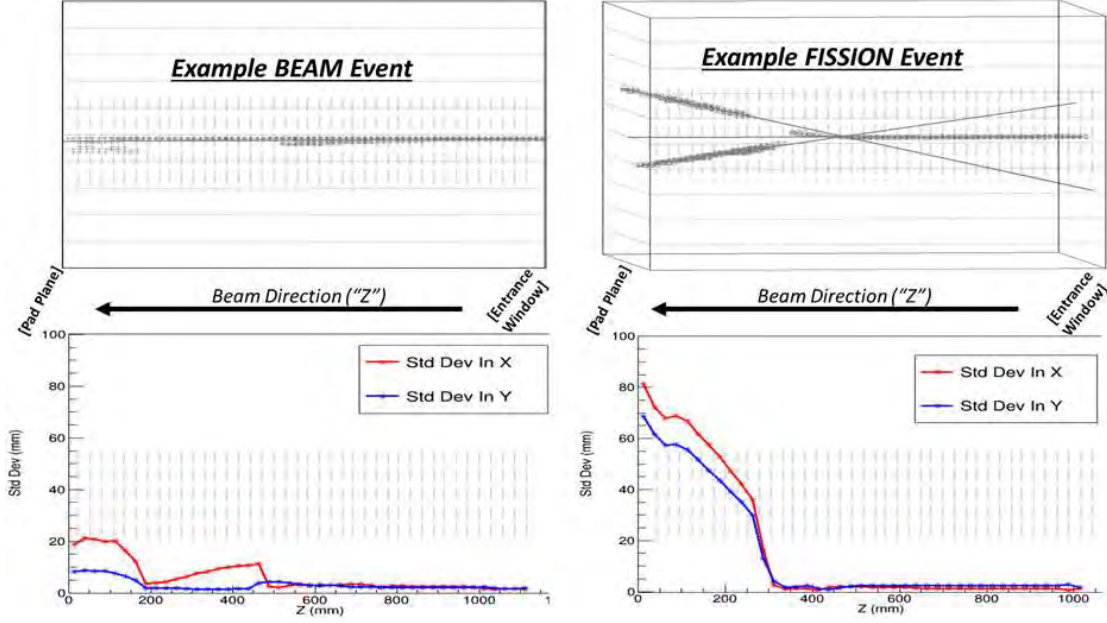


Figure 4.8 (Bottom) Standard deviation of collected charge in x (red) and y (blue) for an unreacted beam (left) and fission (right) event. The standard deviation of the charge increases dramatically after the reaction vertex in fission events, but remains small for unreacted beam.

4.3.2 Further Labeling with Machine Learning

In order keep all potential fission events and only discard unreacted beam, the thresholds for the various metrics of the sorting algorithm were tuned accordingly. If a clear result could not be obtained, the event was left unclassified. This primarily occurred when the fission occurred near the end of the AT-TPC, resulting in short fission fragment tracks. A machine learning project focused on further classifying such events that could not confidently be labeled as unreacted beam or fission was performed using labels from the algorithm described above as a starting point. More details about this project and the results of the sorting algorithm described above can be found in [30]. As discussed in that paper, a point cloud model was trained to identify fission by training on data that was confidently labeled, so that unclassified events could be potentially sorted.

4.4 Space Charge

In this section, I report on the nature of the space charge problem as it pertains to the AT-TPC, and how I dealt with it [18]. I will explain the nature of the space charge problem, and describe how a method was obtained to model, and correct for, the effects of the positive ion space charge column.

Characteristics of the fission data that can be used to validate these corrections are detailed.

4.4.1 The Space Charge Problem

In the context of this fission experiment, the term *space charge* describes a net background density of slowly moving positive charge. In the fission experiment, the space charge drifts throughout the helium TPC gas after it is ionized by unreacted beam particles. Although only a small fraction of beam particles have a nuclear reaction with a helium nucleus in the TPC gas, all beam particles ionize the gas, producing pairs of electrons and positive helium ions. The stopping power of beam ions is proportional to $(Z_{eff}/v_{ion})^2$. Fortunately, both this ratio and the stopping power remain approximately constant as the ions slow down in the gas. The isotopic mixture of beam particles also remains constant, so the positive ion production rate per unit length of track is independent of the distance to the anode and consequently the total space charge remains proportional to the total beam intensity.

The electrons, with their higher drift velocities (0.82 cm/ μ s), quickly reach the pad plane in 122 μ s, after which no electrons remain. The higher drift velocities of the electrons means that they never create a significant space charge effect. On the other hand, the positive helium ions drift much more slowly (<40 m/s) and take more than 25 ms before they reach the cathode near the TPC entrance window. As a result, the helium ions remain in the gas for extended periods, forming a column of positive charge along the beam path. This positively charged column attracts electrons ionized by the fission fragments, distorting their trajectories and making the fission tracks appear curved near in the high space charge region near the reaction vertex, even though the actual fragment paths remain straight. An example of this distortion and its rate dependence can be seen in Figure 4.9.

The distortion is most pronounced for fission fragment tracks in or very near the beam path in the vicinity of the cathode. There, the ion column is densest because most ions are produced by the beam and all ions drift towards the cathode. The ionic charge density strongly decreases as the fission fragments move farther from the beam axis. Beyond approximately 10 cm from the axis, the space charge effects become negligible, and the tracks straighten out. For the parts of the

track that are closer to the beam axis, distortions associated with a particular z position are strongly influence by the space charge distribution of the beam. Correcting this ion-induced distortion is critical for accurately reconstructing the fission tracks back to the reaction vertex, the position of which determines the excitation energy of the compound nucleus formed in the fusion of beam particle and He nucleus.

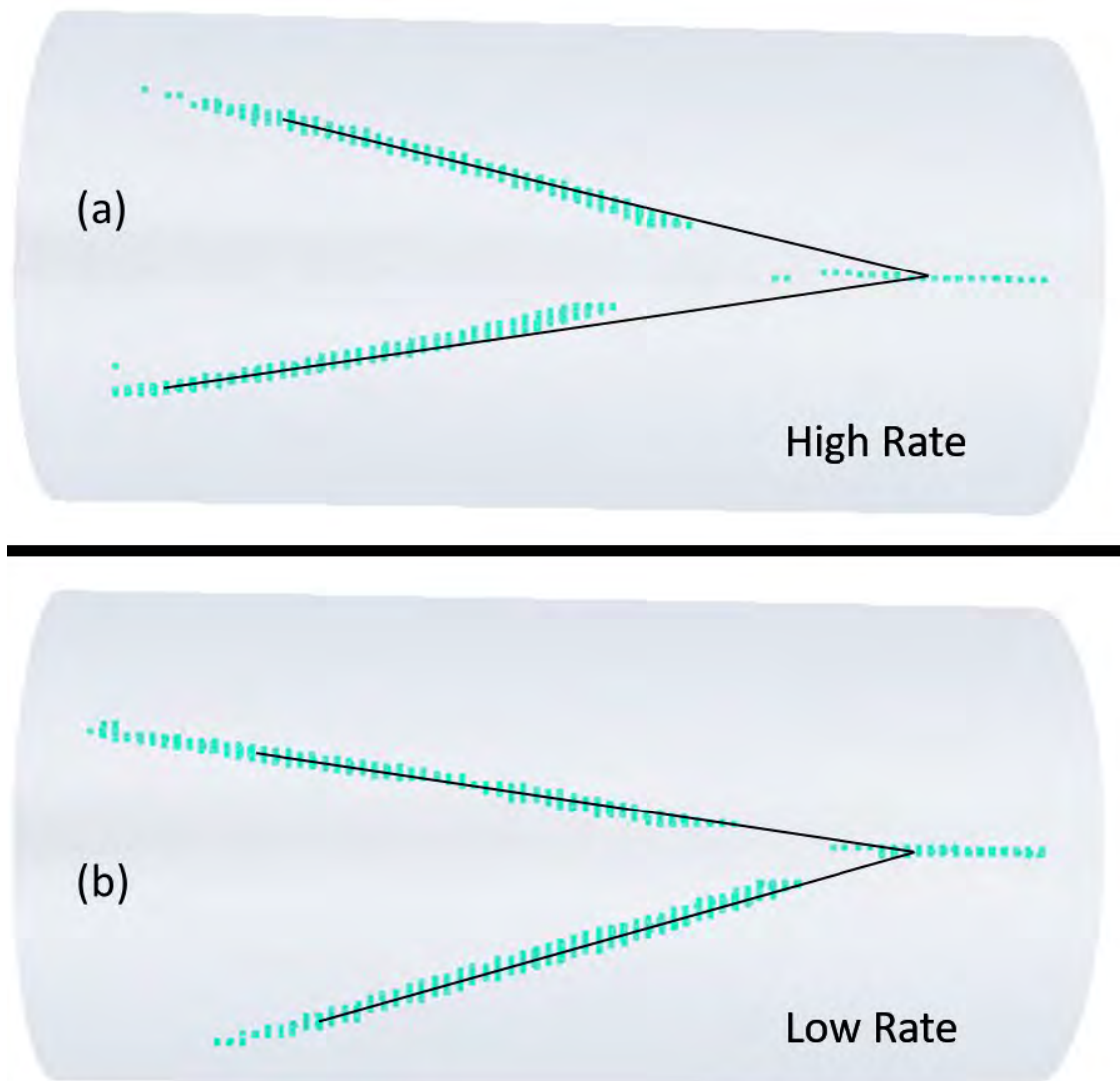


Figure 4.9 Fission events with similar vertex locations from two different runs. Shown in (a) is an event from a run where the beam rate is around 6000 pps, where space charge distortions can be observed in the curved fission fragment tracks. The event in (b) is from a run with beam rate around 2000 pps, where noticeably less space charge distortion occurs.

4.4.2 Y-Fitting

We start our discussion with a method that can be used to analyze fission data at low beam intensities where space charge can be neglected. This is achieved by fitting the event with three line tracks arranged in the form of a “Y” with the reaction vertex at the intersection point of the three lines. In the following, we call this fit a “Y-fit”. An example is shown by the red lines in Figure 4.10. The three line fit to the combined track is initialized for each event by using the RANSAC method [31, 32] to optimize the location of each of the lines. Within this method, four hits are randomly chosen, two within the beam region and two outside of the beam region. The hit that is *second* closest to the window is set as the reaction vertex while the other three define the other end of one of three lines extending from the reaction vertex. Away from the vertex, one of these lines is associated with the initial track of the beam and each of the other two points is associated with the two fission fragment tracks. The fit with the smallest χ^2 value, computed by comparing the calculated tracks in an event to all of the points in the measured tracks, will be selected as the best fit for an event.

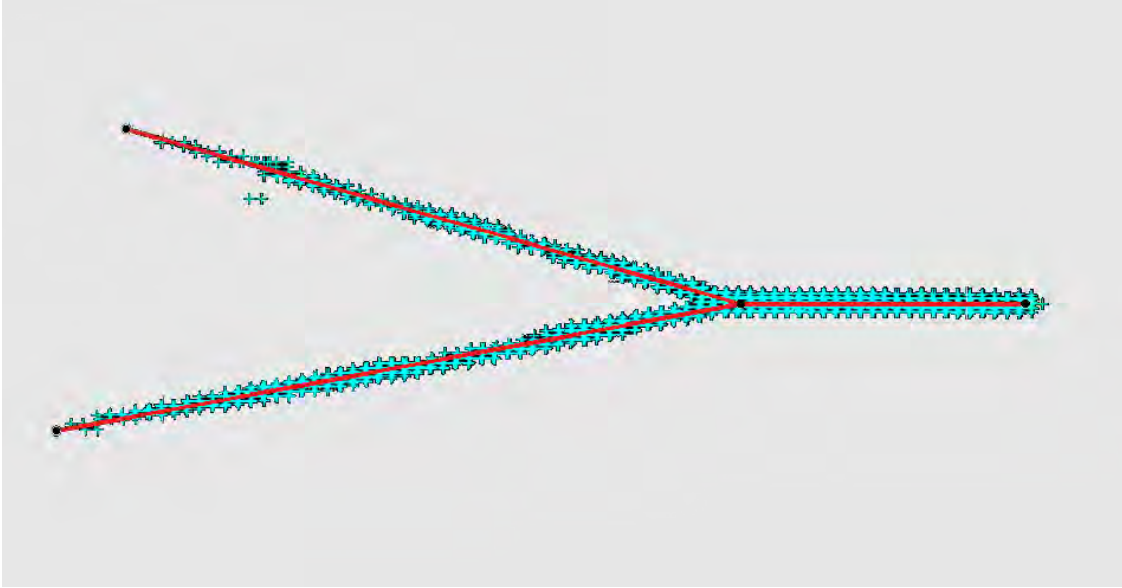


Figure 4.10 Example of the Y-fit on an event recorded during very low beam rate, and in the absence of the smart zap region.

4.4.3 ADCSum Method

The ADC Charge-Sum method has the virtue that it is not distorted by space charge effects, unlike the Y-fit method. Here, we examine the total charge collected in the pads within the beam region after deconvolution as a function of time. This method is sensitive to the sharp decrease in the beam energy loss signal that arises when the heavy projectile fuses with a helium target and the combined system undergoes fission. When this occurs, the energy loss signals in the pads that are collecting the charge ionized by the beam sharply changes, because the stopping power of heavier beam isotopes is greater than the stopping powers of the two lighter fission fragments. Also the charges ionized by the fission fragments are collected by pads in different regions than those in the low gain region that collect the charge of the beam particles. This rapid change in stopping power is spread over several time buckets located around time bucket 420, as shown in Figure 4.11. To increase the sensitivity to this change and better determine the vertex, we unfold the time dependence by differentiating the collected charge in the pre-amplifiers as a function of time. By these means, the rapid change in the stopping power can be more easily identified. This downward slope of signal and its location in time (and the corresponding depth in the TPC) shows where the fusion-fission event occurs, defining the vertex. The vertex of the fission event provided by the “ADC Charge-Sum” method is marked by a black dotted line around time bucket 420. This hit location is represented by the black solid circle in the top panel.

For comparison, the vertical red line in the bottom panel of Figure 4.11 at time bucket 396 indicates the Y-fit vertex before making space charge corrections. Clearly, this Y-fit vertex implies that the fusion fission event occurs later in time. This incorrect conclusion results from neglecting the distortion of the fission tracks caused by the ionic space charge in the beam region. In the following, we explain how these space charge effects can be corrected. When that correction is performed, one obtains the corrected track positions and corrected Y-fit vertex that are shown in the top panel as the blue crosses and solid circle, and in the bottom panel as the blue vertical line at time bucket 418. The result of the Y-fit on the corrected data is clearly in better agreement with the ADCsum result.

The ADCsum method for extracting every fission event vertex is relatively insensitive to space charge. At high beam rates, however, pileup effects make the analysis of the beam tracks more difficult and can even decrease the effective dynamic range of the electronics. In these cases, the ADCsum method does not work. Even though the ADCsum method is not an effective method to find the reaction vertices for all events, we will take advantage of its relative accuracy for suitable selected events to calibrate our corrections for space charge.

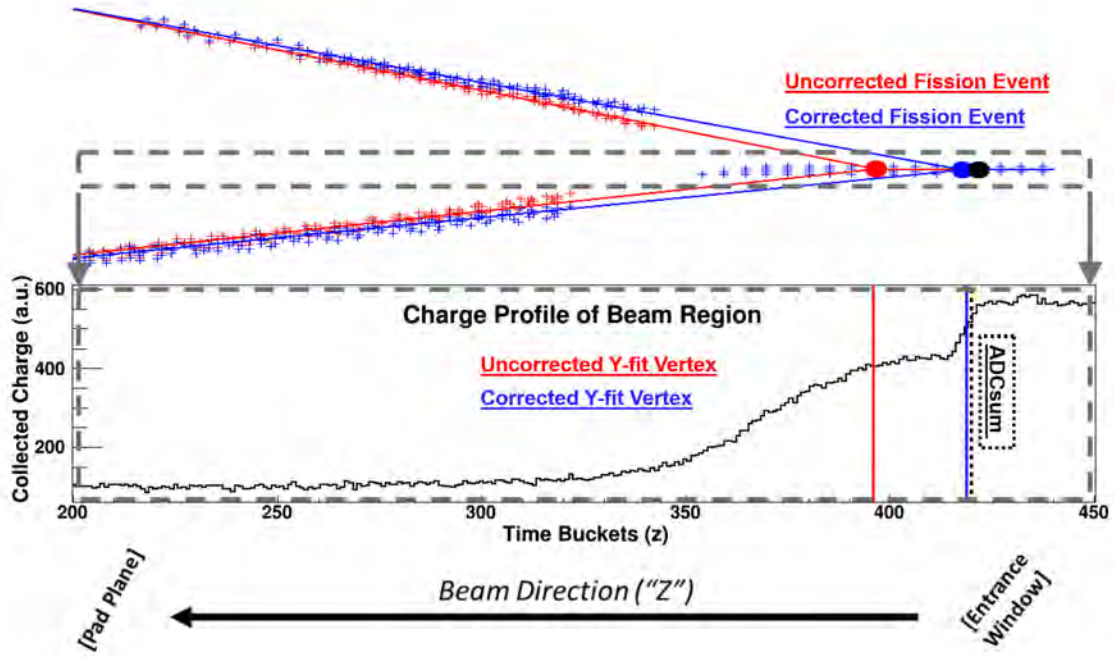


Figure 4.11 Illustration of the ADCsum method.

4.4.4 Langevin Equation

The *Langevin equation* provides a general approach that can calculate the non-equilibrium statistical transport of particles. It is an equation of motion for charged particles, both electrons and ions, in the presence of neutral atoms, and under the influence of electric (**E**) and magnetic (**B**) fields. It has the form:

$$m \frac{d\mathbf{u}}{dt} = e\mathbf{E} + e[\mathbf{u} \times \mathbf{B}] - K\mathbf{u}, \quad \tau = \frac{m}{K} \quad (4.3)$$

where m and e are the particle's mass and charge, $\mathbf{u}$ is the velocity vector, and $K\mathbf{u}$ represents a frictional force that effectively models the energy loss of the ions or electrons from interactions with the gas they are passing through. The quantity τ is a characteristic time between these collisions. When describing electrons or ions drifting in an electric field, we are referring to times much greater than τ . In other words, we can assume that the fluctuations to the drift velocity are negligible, such that $d\mathbf{u}/dt = 0$ [33]. For this fission experiment there is no magnetic field in the setup, so $\mathbf{B} = 0$. The equation of motion then leads to the simplified expressions

$$\mathbf{u} = \frac{e}{m}\tau\mathbf{E} = \mu\mathbf{E}, \quad \mu = \frac{e}{m}\tau \quad (4.4)$$

where μ is the scalar mobility [33]. This equation dictates how the drift velocities of both ions and electrons follow the drift field lines of the electric field exactly.

4.4.5 Modeling the Ion Column

Key to the correction process is the fact that the velocities of the *drifting electrons* and the *drifting ions* follow the field lines of the total electric field. The total electric field has two major contributions: 1) the drift field for the TPC that points away from the pad plane in the $\hat{z}$ direction and pulls electrons towards the pad plane, and 2) space charge fields of the positive ions that point in the $\hat{\rho}$ direction, pulling electrons towards the center line of the AT-TPC. Specifically, the TPC drift field ($\mathbf{E}_{TPC}$) is:

$$\mathbf{E}_{TPC} = |E_{TPC}| \hat{z}, \quad |E_{TPC}| = 700 \frac{\text{V}}{\text{cm}} \quad (4.5)$$

Away from the z -axis of the AT-TPC, the total electric field is given by $\mathbf{E}_{TPC}$, and the drift electrons that are created by ionization, move in the $\hat{z}$ direction with a constant drift velocity of 0.82 cm/ μ s. This drift electric field is enforced by the equipotential surfaces of the cathode and anode planes of the field cage and by the conducting rings of the field cage of the AT-TPC. Away from the space charge region, these equipotential surfaces define the electric field and largely control fluctuations in electric fields that can result from variations in the beam intensity.

The incident beams of Bi, Pb, Hg and Au isotopes each lose energy at a constant rate of about 21 MeV/cm, regardless of charge states. This amounts to a total energy loss in the TPC gas of about 2 GeV. This energy loss process creates high energy delta electrons, some with energies of 10 – 20 keV, which are scattered by the beam to small angles with respect to the beam trajectory. These delta electrons broaden slightly the location of the ionization caused by a fission fragment or a beam particle about its trajectory. The slow drift rates of the ions mean that there will be positive ions from many beam particles contributing to this ionic space charge distribution and the resulting column of helium gas ions therefore has an approximate 4-cm-diameter, which is comparable to the spread of the secondary beam.

In addition to $\mathbf{E}_{TPC}$, we must address the smaller electric field created by ions along the beam trajectory which is centered on the beam axis of the AT-TPC. Because the space charge is created from the ionization of unreacted beam particles, we can use information about the beam to model the charge distribution. Figure 4.12 illustrates the beam profile within the $\hat{x}$ and $\hat{y}$ directions, obtained by fitting straight tracks to several unreacted events from a low rate run. We use these fits to obtain the distributions of position coordinates of the beam at the TPC entrance window and the TPC pad plane. We find that the beam is focused at the entrance window and spreads out upon reaching the pad plane, with typical FWHM values of ≈ 17 mm at the entrance window and ≈ 33 mm at the pad plane. Additionally, the average beam trajectory enters the TPC with a small angle of ≈ 0.013 radians relative to the z -axis, consistent with slight shifts in the average x and y positions at the pad plane. When simulating the drift of ionized electrons using the Langevin equation, a translation of coordinates is used to account for the small angles in the beam trajectory.

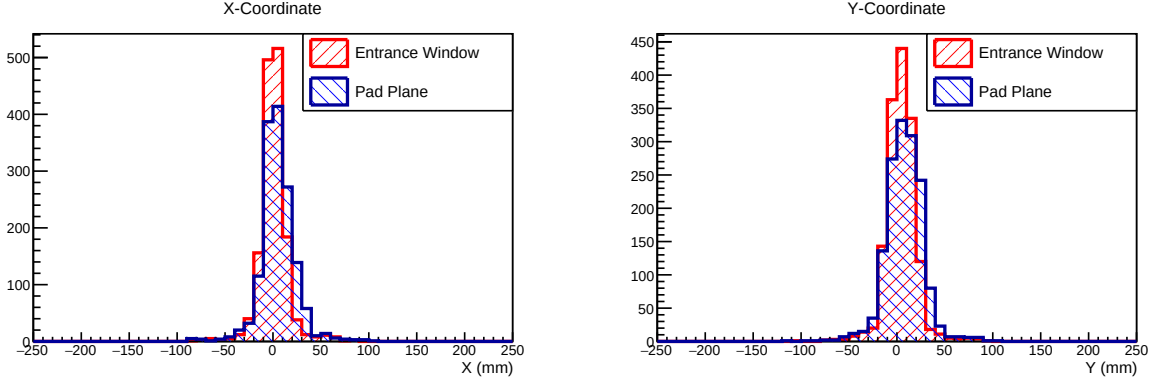


Figure 4.12 Profile of the beam dimensions at the window and at the pad plane.

For simplicity, we approximate this ionic electric field by that of a thin line charge that is centered about the z -axis of the TPC and whose charge density increases with the distance from the pad plane z . This z -dependence occurs because ions are produced with equal probability along the path of the beam and move with approximately constant velocity within the constant electric field of the TPC towards the cathode, where the entrance window is located. Because the positive ion charge density increases linearly along the z axis of the TPC, the distortions caused by the ionic electric field also increase.

Under the assumption the $\hat{z}$ -component is negligible, the induced space charge field ($\mathbf{E}_{SC}$) can be approximated by:

$$\mathbf{E}_{SC} = \frac{\lambda z}{2\pi\epsilon_0\rho L} \hat{\rho} \quad (4.6)$$

where λ is defined as the maximum line charge value of the positive ions that is achieved at $z = L$, with $L = 1$ m as the length of the active volume inside the TPC. Here, ρ is the standard cylindrical coordinate equal to the radial distance of a point from the central axis of the TPC, and ϵ_0 is the vacuum permittivity.

4.4.6 Numerical Data Corrections

To apply space charge corrections, we begin with the recorded hit locations of the collected electrons for a fission event which is compatible with the ADCsum method, such as the red points in the top panel of Figure 4.11. Initially, the extracted vertex locations from a Y-fit and the ADCsum

method will disagree significantly. An arbitrary magnitude for λ is assumed and applied to $\mathbf{E}_{SC}$. The drift of the electron hits are then simulated by numerically solving the Langevin equation with $\mathbf{E}_{SC}$ and $\mathbf{E}_{TPC}$ as inputs for the electric field. The drift is simulated in reverse with opposite field directions, so that an electron recorded at location (ρ, z_0) starts at the pad plane $(\rho, 0)$ and travels towards the entrance window until z_0 is reached.

The time scale for simulating drift is determined by the recorded hit depth z and the electron drift velocity of $0.82 \text{ cm}/\mu\text{s}$, which is around $100 \mu\text{s}$ for electrons ionized close to the entrance window. In practice, we utilize a time step of $dt = 1.0\text{e-}4 \mu\text{s}$, which provides an acceptable balance between accuracy and computational time. It is assumed that the drift of electrons along the $\hat{z}$ -direction is due only to $\mathbf{E}_{TPC}$, and the drift along the $\hat{\rho}$ -direction is due only to $\mathbf{E}_{SC}$. Starting at $(\rho, 0)$, the position coordinates of each electron hit are updated after every time step according to

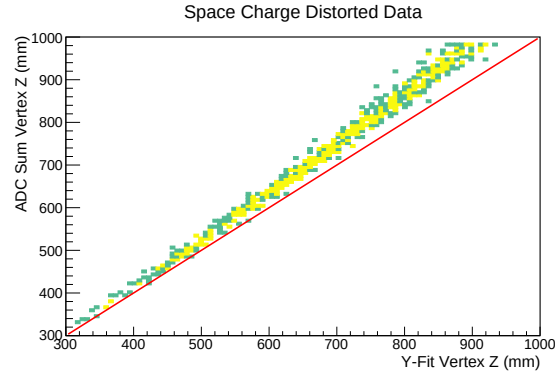
$$\rho' = \rho + dt * \mu_e * \mathbf{E}_{SC}(\rho, z) \quad , \quad z' = z + dt * \mu_e * \mathbf{E}_{TPC}, \quad (4.7)$$

until the electron has reached the recorded depth of z_0 . In this way, the original recorded hit at location (ρ, z_0) has given rise to a new corrected hit with location (ρ', z_0) .

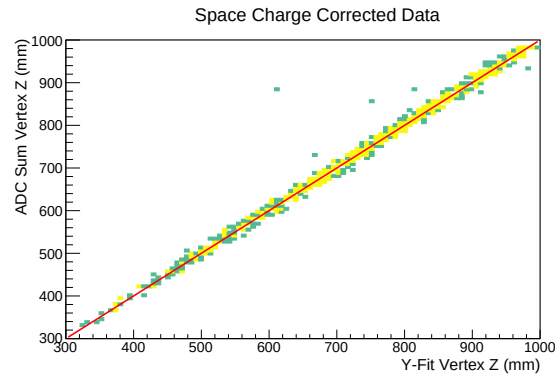
Once a set of corrected hits has been obtained for a given fission event, such as the blue points in the top panel of Figure 4.11, a new Y-fit is performed to compare vertex results with the ADCsum method. If desired agreement between the two methods has not been reached, then a new simulation is performed with a different value of λ . This iterative process allows one to empirically extract the space charge magnitude that corrects for the observed track distortions.

The upper panel of Figure 4.13 shows the disagreement between the two methods within one data run, when space charge corrections are not applied. Instead of a 45° -line (the red reference line), the distortions to the fission fragment tracks due to space charge make the vertex determined by the Y-fit method closer to the pad plane than the actual vertex, determined by the ADCsum method. The lower panel shows the agreement between the two methods after space charge corrections are applied. The χ^2 per degrees of freedom is 3.518 before corrections, and 1.012 after corrections to the time projected fission fragment tracks. Note that in addition to achieving a better agreement

between the ADCsum and Y-fit methods, the spread of differences between the corrected Y-Fit and the ADCsum tracks is also reduced, as evidenced by the tighter correlation.



(a) Before data correction.



(b) After data correction.

Figure 4.13 Difference in vertex depth determination between the Y-fit and ADCsum methods before and after correction of the data, for several fission events recorded at the same beam intensity. All events in this example have been corrected using the same magnitude for λ . The red line of '*Y-Fit Vertex Z = ADC Sum Vertex Z*' is drawn to guide the eye.

4.4.7 Correction Results

The various isotopes present in E12014 enter the TPC with similar energies, and their stopping powers vary by less than 10% across the range of the active target. For a specific isotope, the range of stopping powers across charge states varies by less than 0.1%, and significant changes to charge states inside the TPC are unlikely to begin with [34]. Given the known energy loss of the beam, each ion in the incident beam can be treated as producing approximately 4.60×10^4 positive helium ions per cm in the helium gas of the TPC. By combining this production rate, the known beam intensity, and the ionic drift velocity, one can calculate the expected magnitude of these positive

helium charges. In general, drift velocity can be defined as the product of the a mobility and an electric field, as in Eq. 4.4. If we assume that the mobility of helium ions in helium gas takes a constant value from Ref. [35] ($10.2 \text{ cm}^2\text{V}^{-1}\text{s}^{-1}$), and use $\mathbf{E}_{TPC}$ (700 V/cm), an expected magnitude of $\lambda = 3.5\text{e-}10 \text{ C/m}$ is obtained for a beam rate of 3000 pps. This amount of charge is too small to reproduce our observed space charge effects when applied to $\mathbf{E}_{SC}$. Therefore, we must extract the magnitude from the data.

To first order, the effects of space charge distortion should scale linearly with the beam rate. Fission events compatible with the ADCsum method of vertex determination are selected from runs with varying beam rate. To be compatible, these events must have a strong beam signal measurable by the beam region pads without interference from event pile up or excessive electronic noise. For high rate runs, only a small subset of events fulfill this requirement. The Langevin equation is solved numerically for these ADCsum compatible events using the Forward Euler method.

The relationship between λ and beam rate shown in Figure 4.14 is obtained by repeating this exercise over runs with varying rates. It is important to select data runs which display a stable beam rate for determining this relationship, as rate fluctuations within a run will cause a fluctuating space charge magnitude. The dependence on beam rate is approximately linear over the measured range of beam rates, as expected, and trends towards zero space charge at zero beam rate.

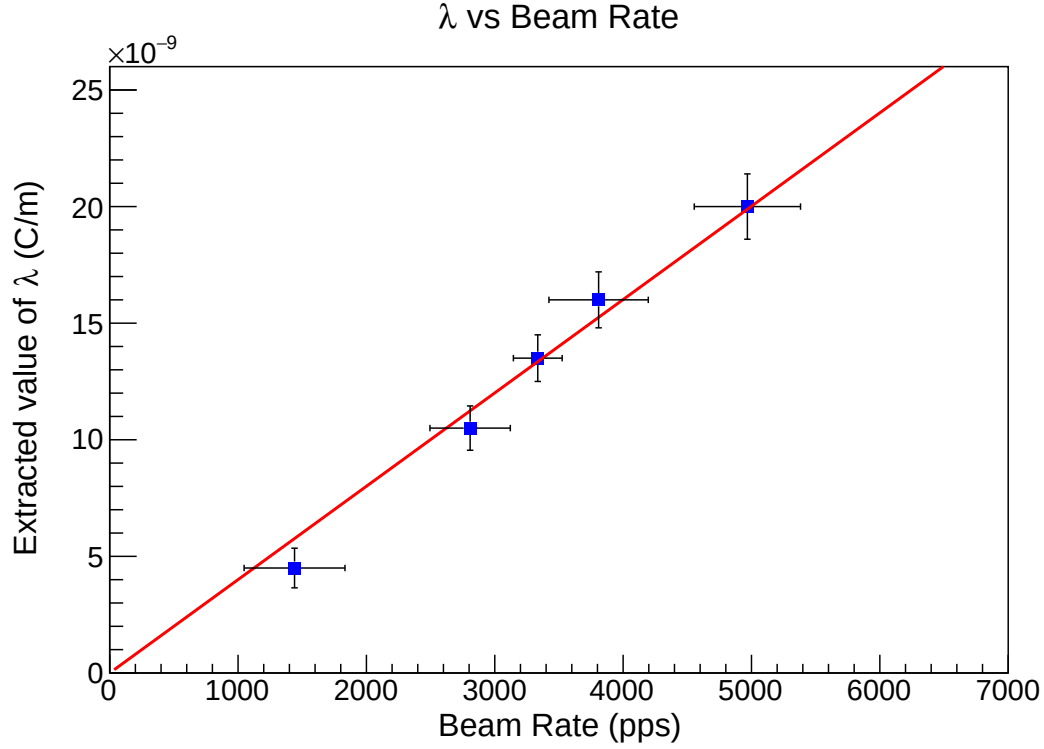


Figure 4.14 Space charge magnitude as a function of beam rate.

4.5 Summary and Outlook

Using the methods detailed in the above sections, it is possible to extract the desired physics observables in this fission experiment E12014. Uninteresting events where fission does not take place may be confidently discarded through systematic application of a sorting algorithm. Further data refinement is possible through additional filtering using machine learning models that are trained on events which were labeled using the initial sorting. The pre-fission nucleus can be identified and tagged from the radioactive cocktail beam through HEIST. Inside the AT-TPC, the active helium target provides both a mechanism for the formation of a CN, and the ability to record the tracks of product fragments in the laboratory frame. A space charge correction method specifically designed for use in active target detectors is used to correct the distorted drift paths of the electrons which record the tracks.

With the fissioning system identified and correct tracks obtained, the remaining exercise is to gather events for a particular isotope with varying vertex locations and extract the fission barrier.

This thesis will not go into detail about the specifics beyond preparing the fission tracks. Currently, barrier extraction has been performed for ^{204}At as the main thesis result for a previous graduate student [17]. Work is currently ongoing to refine and expand the analysis to other isotopes.

Part II

Charged Pion Emissions with the SpiRIT TPC

CHAPTER 5

INTRODUCTION

In this section, I introduce the concept of the Equation of State (EoS) for nuclear matter, and the lack of available constraints at very large densities and high levels of asymmetry which motivate the research of this thesis.

5.1 Background

The nucleus (core) of any atom is composed of *nucleons* of two types, protons and neutrons. Chemical properties of elements are determined by the nuclear charge Z , which is equal to the number of protons present in the nucleus. For a nucleus with proton number Z , there can be different values for the number of neutrons N . The value for N designates the isotope of that element. The values for N and Z of a nucleus govern many nuclear properties, such as its atomic number (A), its mass M , its stability against decay, and the nature of the quantum bound and unbound states that the nucleus can have. Nuclei are quantum liquid droplets of nucleons and of *nuclear matter*.

In Part I of this thesis, a semi-empirical mass formula (SEMF) for the binding energy [36] of nuclei in the Liquid Drop Model (LDM) is introduced:

$$\frac{BE(N, Z)}{A} = a_V - \frac{a_S}{A^{1/3}} - a_C \frac{Z(Z-1)}{A^{4/3}} - a_A \frac{(N-Z)^2}{A^2} . \quad (5.1)$$

Experiments have shown that the binding energy per nucleon is maximized in the deep interior of the nucleus when the density ρ in that interior region is approximately 0.16 nucleons per cubic fm (or $2.3 \times 10^{17} \text{ kg/m}^3$). This density is often referred to as the *saturation density*, or ρ_0 . The unique value for the saturation density results from a balance in the strong nuclear force being repulsive at very short ranges and attractive at larger ranges. At the surface of the nucleus, this density diminishes exponentially to zero with a decay constant of about 0.5 fm. Nuclear matter in the deep interiors of nuclei that we study in laboratory experiments has a density that is very close to ρ_0 . While the radii of nuclei increase with A , the maximum densities in their interiors do not exceed ρ_0 .

One shortcoming of the LDM is that it treats the nucleus as an *incompressible* liquid; that is,

it does not address the properties of matter at densities greater than ρ_0 . Thus a_V represents the binding energy per nucleon in the nuclear interior, which is at saturation density, and $\frac{a_S}{A^{1/3}}$ corrects for the decreased binding of nucleons in the surfaces of nuclei. The LDM does not need to consider the composition of nuclear matter and the force binding it at greater extremes in density. This changes in the deep interiors of neutron stars, for example, where nuclear matter is compressed to densities higher than ρ_0 . It also changes in the stellar surface where the density decreases gradually to zero. To describe the very neutron-rich nucleonic matter of a neutron star, proper modeling dictates a flexible approach that adapts the composition of the nuclear matter to the changes in density and takes into account how the total force changes with the matter composition.

Let us imagine a neutron star to be a giant nucleus in the limit of very large $A \rightarrow \infty$. In the interior of objects as large as neutron stars, the gravitational force binds the star. While nuclear forces are attractive at sub-saturation densities ($< \rho_0$), at the greater nuclear densities in the star the nuclear forces become repulsive and supply pressure that prevents the star from collapsing into a black hole. The electron capture process: $e^- + p \rightarrow n + \nu$ reduces the repulsive energy due to the long range Coulomb forces between protons in the star, converts most nucleons to neutrons and increases the stellar binding energy. Coulomb forces and the surface energy contribute less to the total energy of the star than the volume energy density. If we neglect the surface energy, the binding energy of a droplet of nuclear matter can be simplified to the form

$$\frac{BE(N, Z)}{A} = a_V - a_A \delta^2 \quad , \quad \delta \equiv \frac{(N - Z)}{A} \quad , \quad (5.2)$$

where δ is commonly used to represent the asymmetry between proton and neutron number, where $\delta = 1$ corresponds to pure neutron matter. Then the energy per nucleon can be represented by a “*symmetry*” term ($\delta = 0$) and an “*asymmetry*” term. This simplification becomes even more useful and interesting when we representing the factors a_V and a_S as functions of density as we discuss next.

5.2 The Equation of State of Dense Nuclear Matter

An Equation of State (EoS) is an expression that describes how matter reacts to changes in conditions, such as density, temperature, and pressure. Let us focus on the case of 0 temperature,

wherein the main dependence is on the EoS dependence of density ρ and neutron-proton asymmetry δ . While the SEMF can be used to predict the rest energies of discrete individual nuclei in their ground states, to extend this knowledge to neutron stars (whose densities vary) we also need to describe nuclear matter as a continuous function of density. Inspired by the limit of the SEMF as $A \rightarrow \infty$ (Eq. 5.2), we include this density dependence to obtain a useful representation for the EoS of nuclear matter that can also be applied to nuclei and neutron stars.

To do so, we start with an expression for the energy per nucleon in nuclear matter at a density ρ and an isospin asymmetry δ :

$$\mathcal{E}(\rho, \delta) = \mathcal{E}(\rho, \delta \equiv 0) + S(\rho)\delta^2 \quad , \quad \delta \equiv \frac{(\rho_n - \rho_p)}{\rho} \equiv \frac{(N - Z)}{A}. \quad (5.3)$$

The first term encapsulates the properties of symmetric nuclear matter with equal numbers of protons and neutrons. The second term contains the symmetry energy, $S(\rho)$, which describes how the energy per nucleon is changed when the isospin asymmetry is changed. The symmetry energy is key to understanding the properties of asymmetric nuclear matter. From this expression, the pressure can be calculated for a fixed number of particles as

$$P(\rho) = -\left(\frac{\partial \mathcal{E}}{\partial V}\right) = \rho^2 \frac{\partial \mathcal{E}}{\partial \rho}. \quad (5.4)$$

Because neutron stars are very neutron rich, the symmetry pressure, which is obtained by differentiating the symmetry energy with respect to density, plays a very critical role. The relationship between the symmetry pressure and the density in stellar matter is key to modeling the structure of neutron stars. We can probe such environments in the lab using heavy ion collisions (HICs) of neutron rich matter, discussed in Section 5.3. In the collisions that we study in this dissertation, however, we do not produce strange particles. Such studies may be possible at much higher energies. For a discussion strangeness and the complications it introduces, we refer the reader to [37, 38, 39, 40].

While constraints for the EoS of symmetric nucleonic matter $\mathcal{E}(\rho, \delta = 0)$ exist over a range of densities [41, 42, 43, 44], there are few constraints for the symmetry energy term $S(\rho)$ beyond saturation. Most existing constraints are at or below saturation [44, 45, 46], which can be seen in

Figure 5.2. Constraining the density dependence of $S(\rho)$ is an important goal of this dissertation research. It is important because it experimentally probes the symmetry energy at densities up to $1.5 \rho_0$. It provides an experimental compliment to the stellar constraints on the EoS from the radii and deformability of neutron stars, which can probe densities of about $2\rho_0$ and above [47, 48].

When considering the variation of the EoS and the symmetry energy as a function of density and asymmetry, we will find it useful to parametrize the density dependence of the symmetry energy through a Taylor expansion near saturation density using $x = (\rho - \rho_0)/3\rho_0$ [10]:

$$S(\rho) = S_0 + Lx + \frac{1}{2}K_{sym}x^2 + O(x^3), \quad (5.5)$$

where L and K_{sym} are the slope and curvature of the symmetry energy with respect to density:

$$L = 3\rho_0 \left. \frac{\partial S(\rho)}{\partial \rho} \right|_{\rho_0}, \quad K_{sym} = 9\rho_0^2 \left. \frac{\partial^2 S(\rho)}{\partial^2 \rho} \right|_{\rho_0}. \quad (5.6)$$

Transport codes can simulate $S(\rho)$ with different density functionals using different Taylor expansion parameters L and K_{sym} [9] to determine the EoS that best reproduces experimental measurements, as is shown by the curves in Figure 5.1. At densities above ρ_0 , different values for the L and K_{sym} parameters can lead to vastly different values of $S(\rho)$. In the figure, the dash-dotted curves that tend towards higher symmetry energies at high density are called *stiff*, and the solid curves that tend towards lower symmetry energies at densities greater than $2 S(\rho)$ are called *soft*.

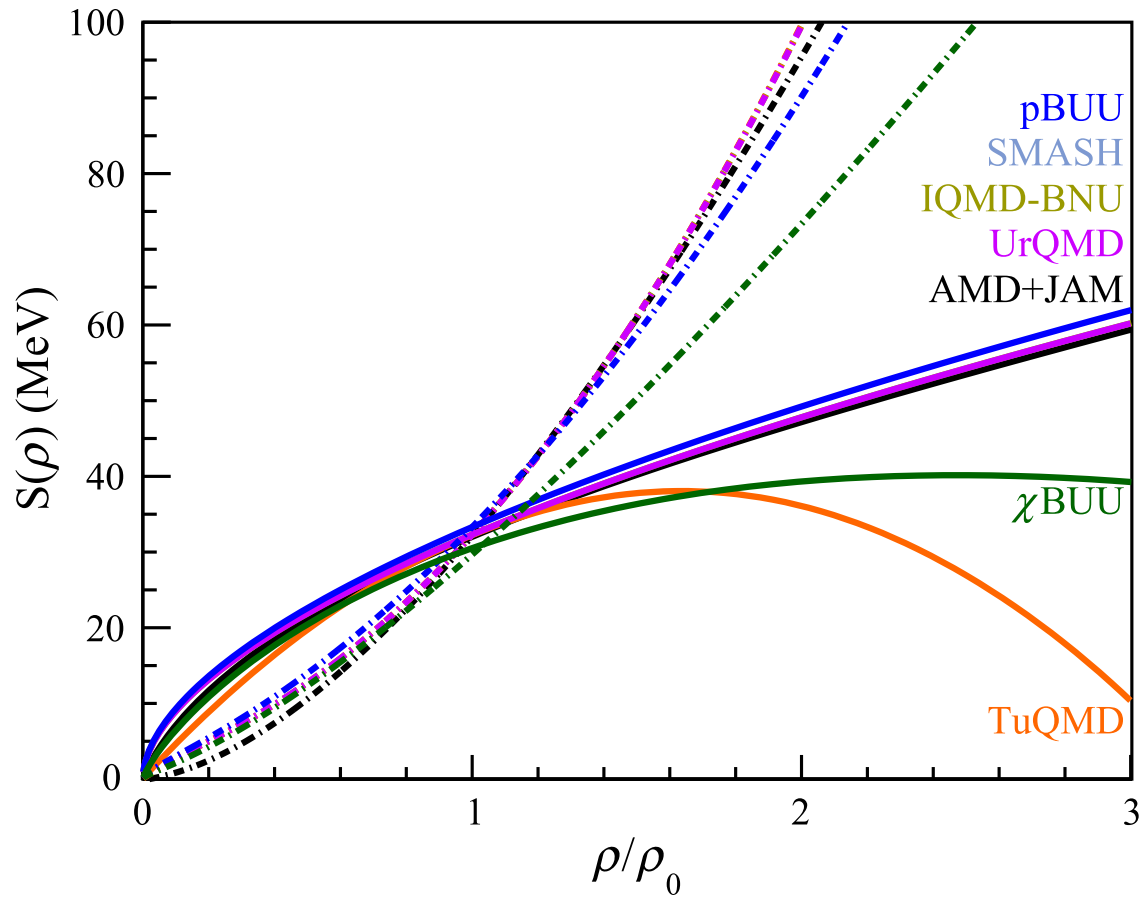


Figure 5.1 This figure illustrates various transport model calculations that use different values for the slope and curvature of $S(\rho)$. Solid lines represent softer symmetry energies, and dot-dashed lines represent stiffer. Figure from [9].

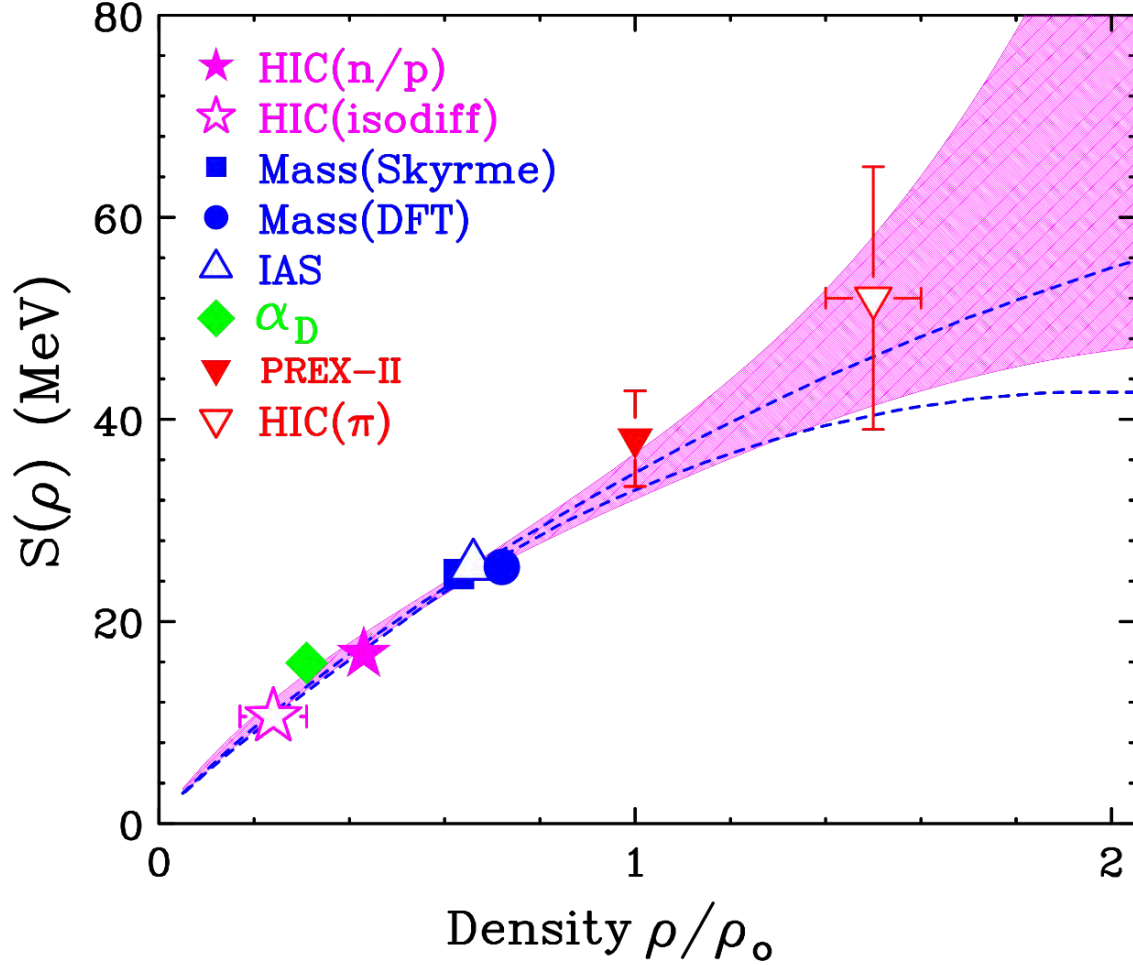


Figure 5.2 This figure illustrates the lack of constraints on the symmetry energy for densities above saturation. Figure from [10].

5.3 Heavy Ion Collisions and the Pion Observable

Figure 5.3 shows a cartoon illustration of a Heavy Ion Collision (HIC), where two nuclei collide with each other. Nucleons are classified by those that strongly participate in the collision, *participants*, and those which are less involved and just “observe” the collision, *spectators*. The distance between the colliding nuclei is called the impact parameter (b), where smaller impact parameters correspond to more central collisions and therefore higher density regions among the participants. HICs are capable of momentarily creating nuclear systems at different densities through variation of beam energies or impact parameters. Asymmetry (δ) can also be probed by using beam-target systems of differing compositions of neutrons and protons, such as those available at RIB facilities. Particularly for symmetry energy studies, the relative emissions of

particles in an isospin multiplet such as neutrons and protons are of interest, since they experience symmetry forces of opposite signs. Sensitivity to symmetry effects are enhanced even further by taking particle spectra ratios.

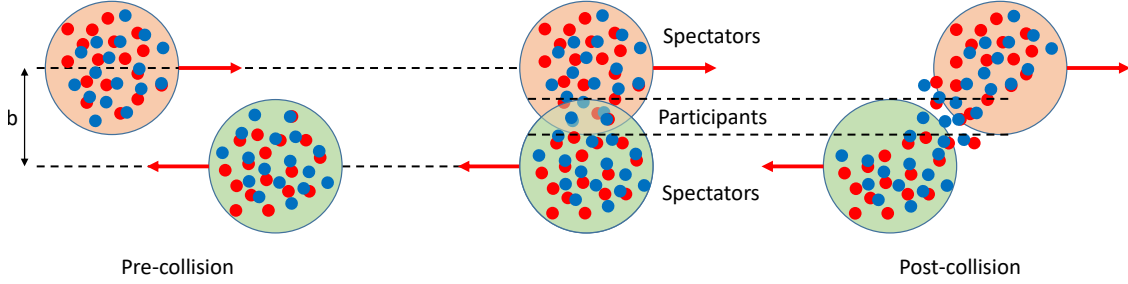


Figure 5.3 A cartoon illustration of heavy ion collisions showing the impact parameter b , and nucleons before and after colliding. Figure from [11].

It is possible to probe the symmetry energy using n/p ratios [49, 50], but such studies require measurement of high energy neutron data from which the backgrounds caused by scattering of produced neutrons off of the floor and off of the ancillary equipment around the target have been removed. Instead, we measure both negative and positive pions in central collisions and construct ratios of the π^- energy spectra divided by the π^+ spectra, for central heavy ion collisions. Despite the lower overall yields in HICs compared to neutrons and protons, charged pions are both easily detectable and identifiable in radiation detectors. Charged pion production comes mainly from resonance excitation and decay of the Δ baryons from high-density compressed regions of participant nucleons in the early stages of HICs [51, 52, 53, 54]:



where $\mathbb{N}$ and π represent a nucleon and pion. Primarily, π^- and π^+ are produced in neutron-neutron and proton-proton collisions, respectively. In the $\Delta(1232)$ resonance model for pion production, the ratio is expected from branching ratios to be [55, 56, 57]

$$\pi^-/\pi^+ = \frac{5N^2 + NZ}{5Z^2 + NZ} \approx (N/Z)^2. \quad (5.8)$$

Since the sensitivity of these measurements to the symmetry energy is related to the N/Z ratio, the sensitivity of these pion constraints to the symmetry energy will be maximized by comparing

collisions of very neutron rich systems to collisions with very neutron deficient systems.

5.4 Motivation

A previous campaign was conducted in 2016 for various Sn + Sn beam-systems of differing N/Z (1.2, 1.36, 1.56) at energies around 270 MeV/u in order to measure charged pion production from the production of Δ baryon resonances and their subsequent decays into pions [58, 12, 9]. This incident energy was chosen because it was the highest incident energy where one could study reactions induced by rare isotope tin beams. The previous experiment focused on comparing the production of negatively charged π^- to positively charged π^+ mesons because the momentum dependence of the ratio of their production is predicted to be sensitive to the density dependence of the symmetry energy. The tin isotopic beams included ^{108}Sn , ^{112}Sn , ^{124}Sn , and ^{132}Sn , and the targets were composed of ^{112}Sn and ^{124}Sn . Table 5.1 summarizes these beam-target systems. The $S\pi\text{RIT}$ TPC, which will be explained in Section 6.3, was constructed in order to measure pion and charged particle production with a special focus on the needs of these measurements. Spectral yields of pions and other produced particles, and ratios of these yields, were obtained and compared to theoretical transport model calculations [9]. This enabled constraints to be placed on the symmetry energy at nucleon densities of approximately $1.5\rho_0$ [12, 59], indicated by the $HIC(\pi)$ label in Figure 5.2.

Beam-Target System	E/A (MeV)	(N-Z)/A
$^{132}\text{Sn}+^{124}\text{Sn}$	270	0.22
$^{124}\text{Sn}+^{112}\text{Sn}$	270	0.15
$^{112}\text{Sn}+^{124}\text{Sn}$	270	0.15
$^{108}\text{Sn}+^{112}\text{Sn}$	270	0.09

Table 5.1 A list of the systems employed in the 2016 experiment.

However, the theoretical calculations under-predicted the measured pion ratio yields at low momenta, as can be seen in Figure 5.4. This raised the possibility that in HICs, non-resonant pion production could be responsible for this excess in the pion production at low momenta. For further details about the energy thresholds for pion production in nuclear collisions, see Appendix A. We decided to test this hypothesis, by repeating the measurement at higher energies where the delta

resonance contribution would be more dominant [60]. Furthermore, the statistical uncertainties in the data for the double ratio were large at higher p_T shown in Figure 5.5, offering less precise constraints than the single ratios. Our repeated measurement will offer more precise constraints on the symmetry energy, in addition to exploring the role of reduced non-resonant pion production.

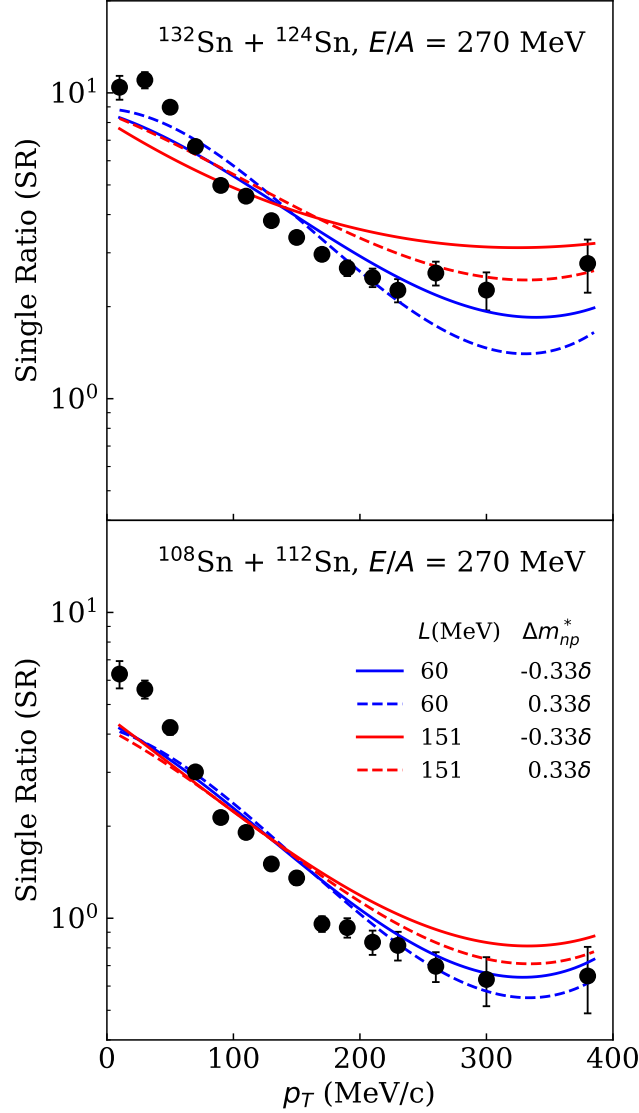


Figure 5.4 Comparison of the experimental charged pion single ratio from the 2016 measurements of $^{132}\text{Sn}+^{124}\text{Sn}$ (top) and $^{108}\text{Sn}+^{112}\text{Sn}$ (bottom) with QMD transport model calculations. Figure from [12].

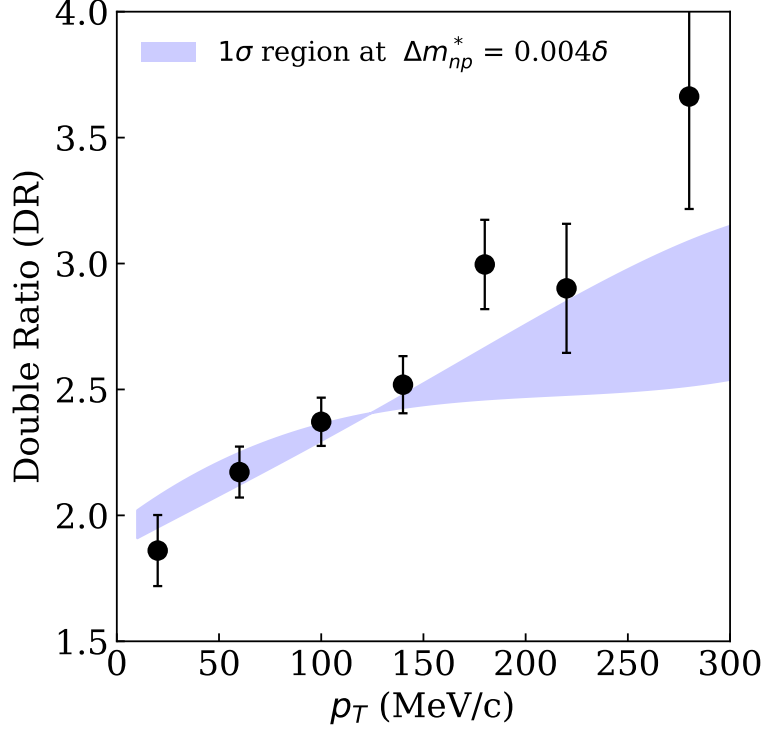


Figure 5.5 Comparison of the experimental charged pion double ratio for $^{132}\text{Sn}+^{124}\text{Sn}$ over $^{108}\text{Sn}+^{112}\text{Sn}$ with dcQMD transport model calculations. The shaded region corresponds to the 1σ confidence interval of the most probable L and Δm_{np}^* values obtained from trying to reproduce the pion single ratios above $p_T = 200$ MeV/c. Figure from [12].

Because rare isotope beam energies are less than the energies of the primary beams that are fragmented to produce secondary beams, higher energies could not be achieved with the secondary tin beams. We therefore successfully proposed in 2020 to probe the role of non-resonant pion production using xenon primary beams at the RIBF accelerator facility in Wakoshi Japan. Accelerator problems and the COVID epidemic delayed this experiment for nearly four years.

This campaign was finally conducted in 2024 with isotopic beams of ^{124}Xe and ^{136}Xe at energies around 320 MeV/u, allowing the desired measurements at higher energies and densities and the desired test of the role of non-resonant pion production. The systems of the 2024 measurements are listed in Table 5.2. In the following sections, we describe the equipment used in the experiment, the analysis procedure for the data, and the new pion production results. Preliminary comparisons to transport models will be shown at the end.

Beam-Target System	E/A (MeV)	(N-Z)/A
$^{136}\text{Xe}+^{124}\text{Sn}$	320	0.20
$^{124}\text{Xe}+^{112}\text{Sn}$	320	0.11

Table 5.2 A list of the systems employed in the 2024 experiment.

CHAPTER 6

EXPERIMENTAL SETUP

In this chapter I detail the facility and detectors used to measure pion production through heavy ion collisions (HICs).

6.1 RIKEN and RIBF

Experimental measurements of heavy ion collisions that we used to constrain the Equation of State of neutron rich matter were performed at the Rikagaku Kenkyusho (RIKEN) Nishina Center for Accelerator-Based Science [61] located near Tokyo in Wakoshi, Saitama, Japan. To perform these experiments, we joined forces with the SAMURAI Collaboration at the Radioactive Isotope Beam Facility (RIBF) to form the $S\pi$ RIT Collaboration. Its focus is on experimentally constraining the Nuclear Equation of State (EoS) at densities that are well above the saturation density ρ_0 of nuclear matter. RIBF is a multistage accelerator complex that allows the exploration of a variety of research topics. It is capable of producing unstable nuclei up to Uranium at maximum energies dictated by the capabilities of the Superconducting Ring Cyclotron (SRC) [62]. It is this capability of accelerating fast beams of heavy isotopes that makes RIKEN an excellent laboratory for experimentally probing the Equation of State (EoS) of dense nuclear matter.

6.2 SRC and the SAMURAI Dipole Magnet

Figure 6.1 shows the beamline configuration of the Superconducting Ring Cyclotron (SRC) at the RIBF facility near the Superconducting Analyzer for Multi-particles from Radioisotope beams (SAMURAI) [63], which is a multipurpose superconducting magnetic dipole spectrometer that was used for our measurements. It is used in a variety of experiments and operating parameters, given in Table 6.1, and can be varied depending on the goals of each experiment. A 3-D rendering of the SAMURAI magnet is shown in Figure 6.2.

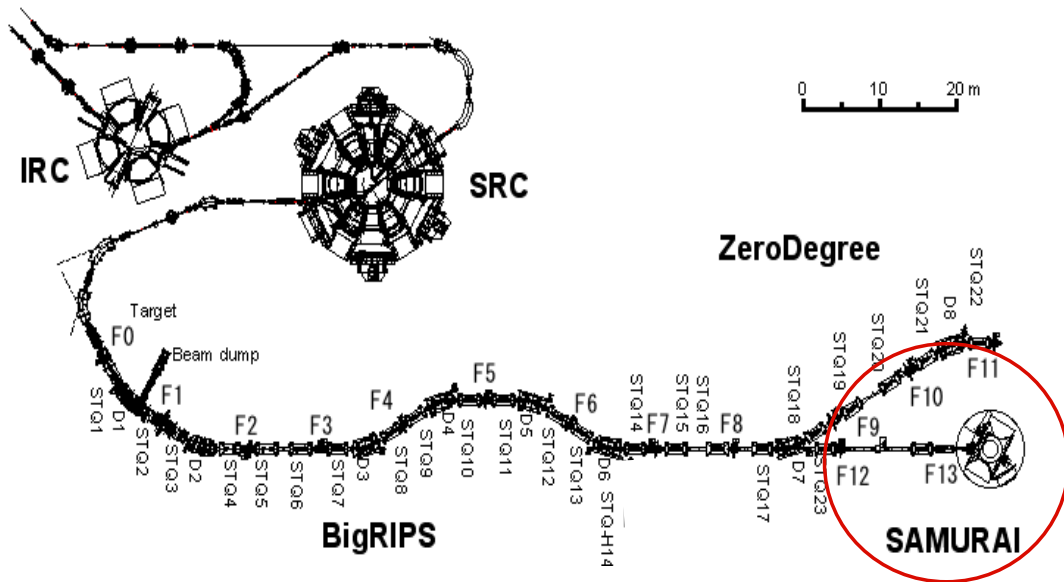


Figure 6.1 Overview of the RIBF beamline as it leads into the SAMURAI magnet, circled in red.

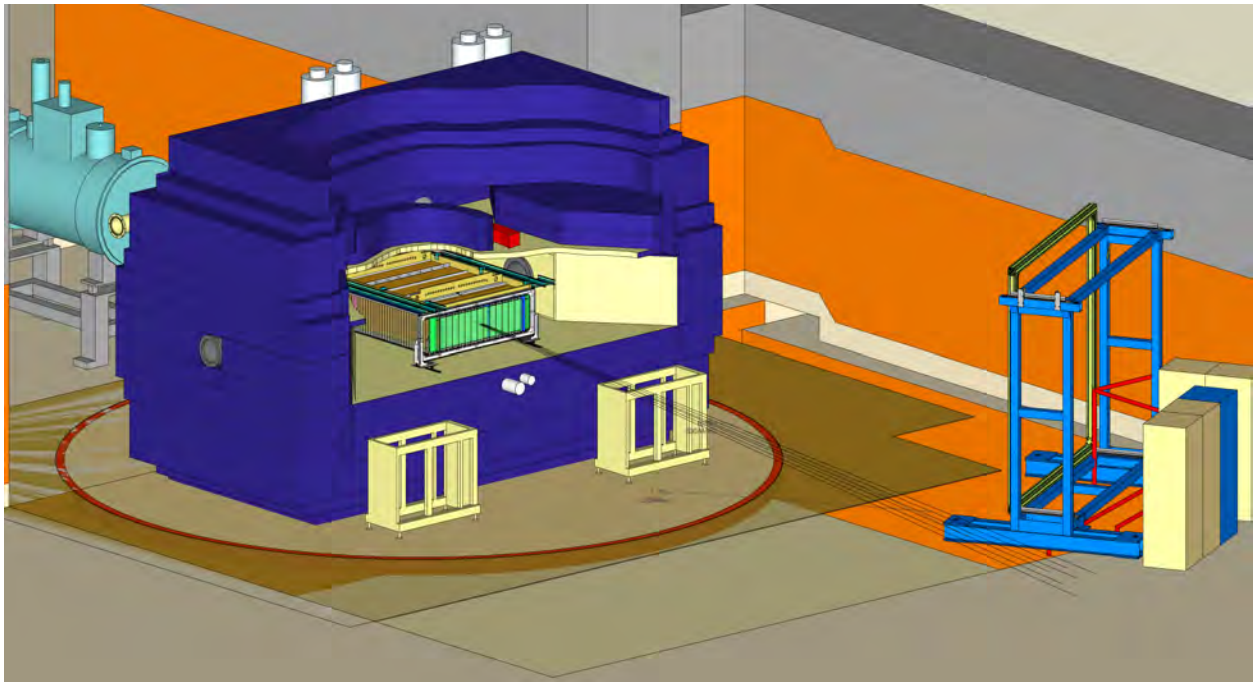


Figure 6.2 Illustration of the SAMURAI magnet (dark blue) with the $S\pi$ RIT TPC inserted. Parts of the top of SAMURAI are made transparent in this rendering to provide a better view.

Table 6.1 Operating parameters of the SAMURAI magnet.

Property	Value
Magnet type	Superconducting H-type dipole magnet
Maximum central field	3.08–3.10 T
Maximum field integral	7.0–7.05 T·m
Maximum current	563 A
Experimental Field Setting	0.5 T
Pole gap	0.8 m effective (0.88 m physical)
Pole diameter	2.0 m
Horizontal opening	3.4 m

6.3 The S π RIT TPC

The SAMURAI Pion Reconstruction and Ion-Tracker Time Projection Chamber (S π RIT TPC) shown in Figs. 6.3 and 6.4 is a gas detector that was specifically designed for use in the SAMURAI magnet as shown in Figure 6.2. It is a rectangular TPC designed to measure momentum distributions of pions and other light fragments from fixed target collisions. It has a detection volume of dimensions $86.2\text{cm} \times 51.3\text{cm} \times 134.4\text{cm}$ (width, height, length). Similar to the AT-TPC, GET electronics are used in operation [64]. Table 6.2 summarizes some of the important operating parameters for the S π RIT TPC used in this experiment. Charged particles inside the S π RIT TPC will ionize electrons from the detection gas which will be drifted by an applied electric field to the pad plane for collection and readout, illustrated in Figure 6.4. The amount of charge collected is converted into an energy deposition, dE/dx . Due to the external SAMURAI magnetic field, the charged particles will follow a curved path based on their *rigidity*

$$B\rho = \frac{p}{q} = \frac{mv}{q}, \quad (6.1)$$

where m and q are the particle mass and charge. As will be described later in Section 7.1.2, tracking software is used to analyze this curvature to determine the particle momentum, p . Combining the dE/dx information with reconstructed p allows for particle identification (PID) like the plot shown in Figure 6.5.

Table 6.2 Properties of the $S\pi$ RIT TPC, modified from [1].

Pad plane area	134.4 cm $\times$ 86.4 cm
Number of pads	12,096 (112 $\times$ 108)
Pad size	1.2 cm $\times$ 0.8 cm
Drift distance	50.5 cm
Electronics sampling frequency	40 MHz
Signal shaping time	117 ns
Electronic noise	800 q_e
Gas pressure	1 atm
Typical gas composition	90% Ar + 10% CH ₄
Gas gain	$\sim$ 1000
Electric field	132.7 V/cm
Magnetic field	0.5 T
Drift velocity	5.45 cm/ μ s
Event rate	50-60 Hz
Typical track multiplicity	60

In addition, the $S\pi$ RIT TPC was surrounded by auxiliary trigger detectors on both sides and downstream. The directions of the electric field (upward) and magnetic field (downward) are also indicated on the right side of Figure 6.4 with black arrows. In the counter-clockwise direction from the lower right of Figure 6.3: The Active Veto is a scintillator array that vetoes any beam particles that do not go through the target. The voltage step down is a series of copper rings on the bottom plate of the TPC that are designed to avoid sparking from the cathode voltage of 6,700 Volts to ground. The target mechanism (ladder) is a tray containing 5 positions upon which targets used in the experiment are mounted. Each target can be moved via external controls to be centered on the beam position during the experiment. The top plate and ribs are a reinforced Aluminum plate structure that serves as the top lid of the TPC and supports the readout electronics mounted on the external side of the plate. The pad-planes are glued to the underside of the top plate. Below the pad plane are 3 wire planes as shown in Figure 6.7. The field cage defines the drift field and also isolates the clean counter gas. With the applied 6,700 Volts and a drift distance between the cathode and ground plane wires of 50.5 cm, the strength of the TPC drift $\mathbf{E}$ -field is 132.7 V/cm. Finally, the enclosure is an Aluminum box with thin walls and windows that contain and support all the above components except the target mechanism and the active veto counter. A more detailed

description of the TPC can be found in [1].

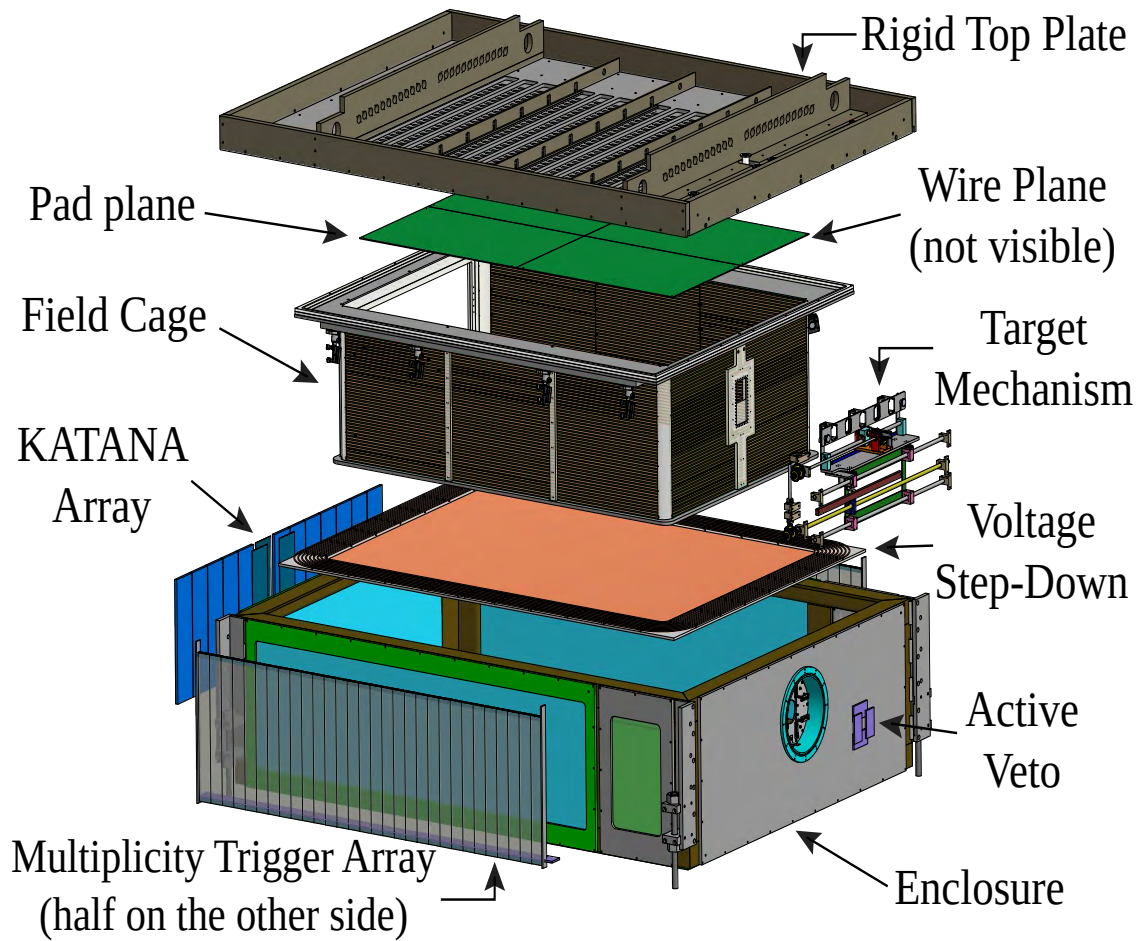


Figure 6.3 A deconstructed view of the various parts composing the $S\pi$ RIT TPC. Adapted from [13].

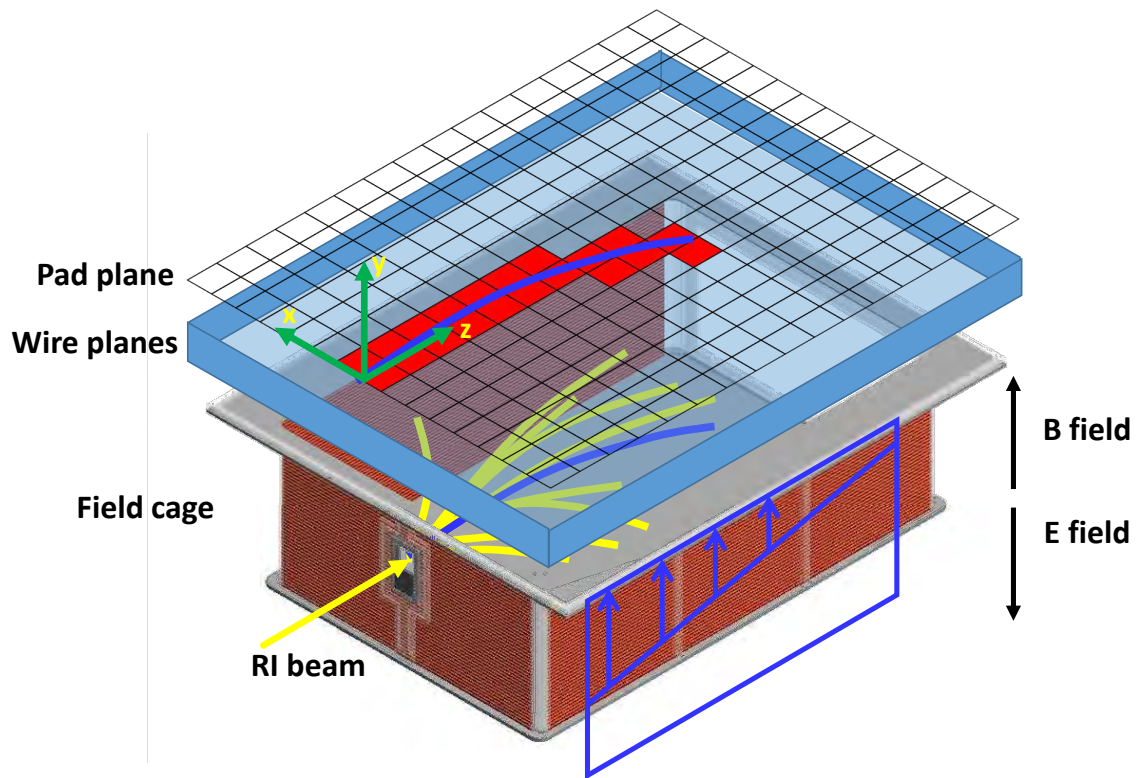


Figure 6.4 An illustration of how the ionized electrons from a charged particle track are collected by the pad plane. The coordinate system and field directions are noted by the green and black arrows, respectively. Adapted from [11]

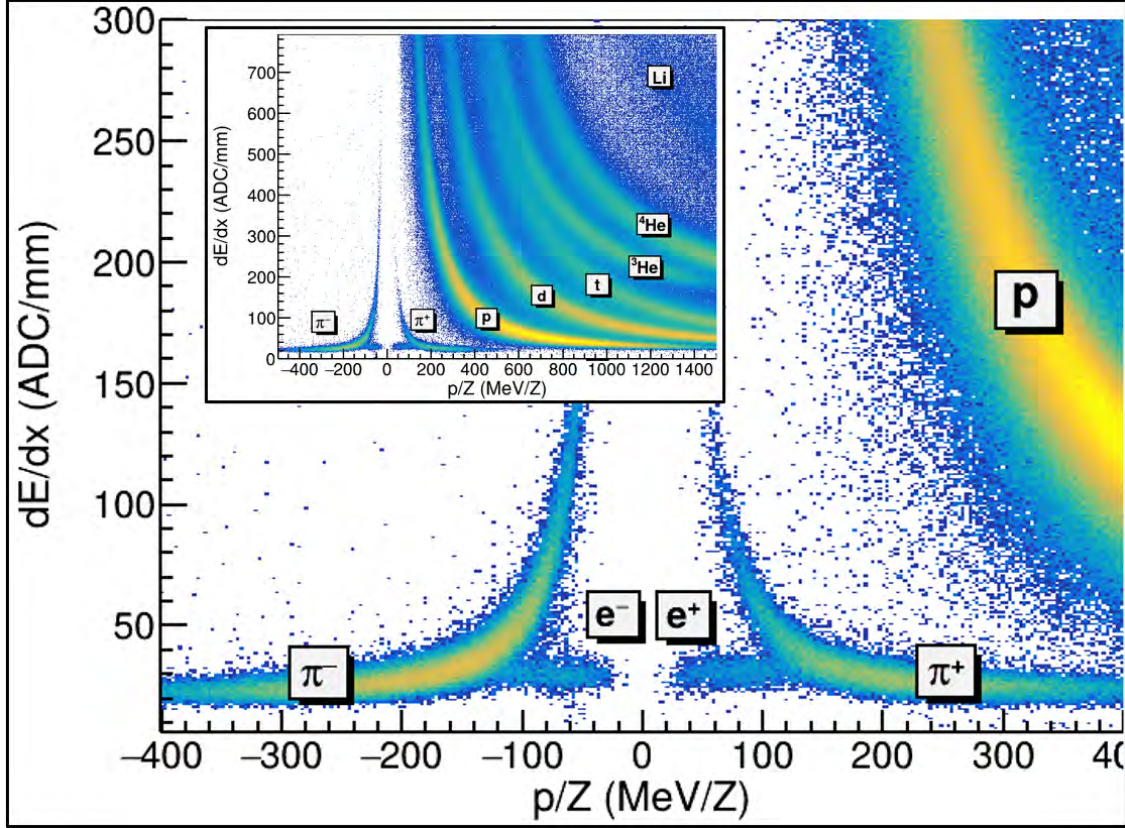


Figure 6.5 A particle identification (PID) plot from the S π RIT TPC, with energy deposition in the TPC gas (dE/dx) on the y-axis and magnetic rigidity (p/Z) on the x-axis. Particle lines are labeled. The inset figure shows a larger range of the identified charged particles, while the larger figure is zoomed in on the pion lines, which are the focus of this dissertation.

6.3.1 Kyoto and KATANA Multiplicity Trigger

The S π RIT TPC trigger is generated in a coincidence module with many inputs. It assures that the trigger signal opens the gating grid in about 220 ns, so that the electrons ionized by that event can transmit through the gating grid and their track information can be digitized. The creation of the S π RIT TPC trigger begins with the detection a beam particle on the Scintillating Beam Trigger (SBT) scintillator, which is mounted in the beamline upstream of the TPC. A copy of the SBT signal is also a delayed input to the trigger and its overlap with the other components of the trigger defines the timing of the trigger output signal. To avoid digitizing events off target, one additional component to the trigger is the signal from an active collimator that provides an output to veto the event if the beam particle will not hit the target. There is also a downstream KATANA scintillator

behind the TPC that can veto the readout if the Xe beam or a remnant of the beam with $Z > 20$ passes through the TPC. The occurrence of a KATANA signal will prevent a trigger from peripheral events from being generated. To further increase the likelihood of the event being central, the trigger also requires signals from at least 4 light charged particles produced in beam-target collisions that were detected in the Kyoto array of scintillators that cover the left and right sides of the TPC. This guarantees that there is a nuclear collision, but later analysis of the event measured with the TPC is required to determine the centrality of that collision. The locations of these two arrays during the experiment can be seen in Figure 6.6. Since the trigger decision must precede the start of readout, this urgency prevents the use of more elaborate analyses to remove every one of the less interesting events from the data stream. Nonetheless, the trigger successfully removes most of the extremely peripheral events

Of these two devices, the trigger condition enforced through the use of the Kyoto array rejects the majority of peripheral events. This Kyoto array, also called the multiplicity trigger array in Figure 6.3, consists of panel detectors located on the sides of the TPC orthogonal to the beam direction (beam left and beam right). If a central or semi-central nuclear collision occurs between a beam particle and the target in front of the TPC, the number of charged particles detected in the Kyoto array provides useful information about how central the collision was. Since we are interested central collisions, we don't record the data of events with a charged particle multiplicity in the Kyoto array that is less than 4 charged particles. To determine a relationship between the number of charged particles produced and the centrality of the collision, discussed in Section 7.4, certain *minimum bias* runs were recorded where the requirement of the Kyoto array was systematically lowered to 1. This is so that even the most peripheral events were able to satisfy trigger and be read out.

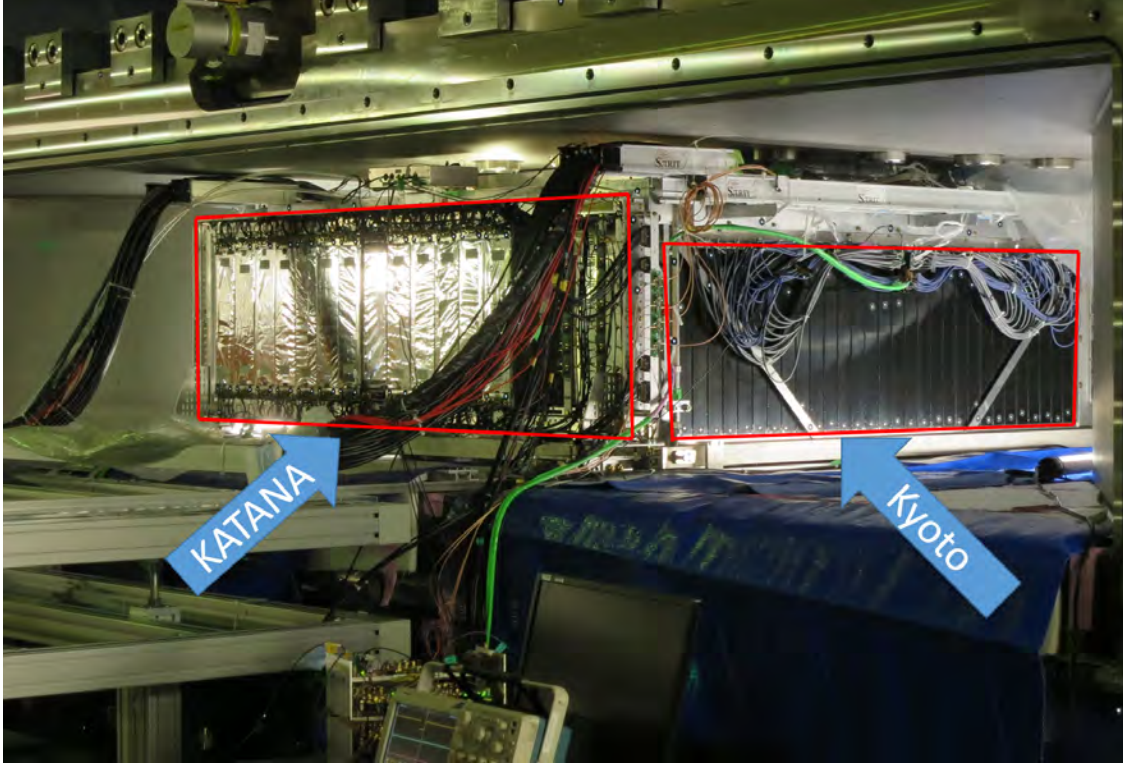


Figure 6.6 Picture of the KATANA and Kyoto trigger detectors as they are installed onto the $S\pi RIT$ TPC.

6.3.2 Gating Grid

To prevent the gas amplification on the pad plane from ionization induced by the unreacted primary beam, a gating grid (GG) is used [65]. An illustration of how the GG works is shown in 6.7; in the closed configuration (right), drifting electrons are terminated on the GG wires. The **closed** configuration is enforced by supplying alternate voltages to neighboring wires, so as to create a strong localized electric field that repels electrons from the negative wires and captures the electrons on the positive wires. When a desired event satisfies the trigger condition, the positive and negative wires on the gating grid are shorted to each other, setting all to the average voltage and **opening** the gating grid. Then the electrons will pass through the gating grid to the anode wires of the TPC where their signals are gas amplified. Garfield simulations of the field lines in the open and closed configurations are shown in Figure 6.8.

Opening and closing of the gating grid is performed by the gating grid driver (GGD), which was improved for this experiment [14] over the version employed in the 2016 campaigns. Figure

6.9 shows the circuit logic for the gating grid driver that supplies the voltages that open and close the gating grid, and allows for rapid switching between the configurations. To open the gating grid, a positive Transistor-Transistor Logic (TTL) signal is applied to the input labeled TTL-1, which closes the metal-oxide-semiconductor field-effect transistor (MOSFET) switches N1 and P1 located on the right side of Circuit A. The duration of the TTL-1 signal determines how long the switches remain closed. Once these switches are closed, the positive and negative wires in the gating grid are connected through $R_p + R_n$, which are connected in series. While not a dead short, the small series resistance of $R_p + R_n$ quickly drains the charge stored in the gating grid. When the charge decreases sufficiently, the gating grid will allow electrons from the TPC to enter the avalanche region.

To close the gating grid after event readout or rejection, we terminate the TTL-1 signal and send a TTL signal to the TTL-2 input. This closes the switches N2 and P2 in Circuit B, connecting the two sides of the gating grid to high-voltage supplies labeled HV-HIGH (V_h) and HV-LOW (V_l). The wires are then recharged within $3 \mu\text{s}$. After this period, the N2 and P2 switches reopen, and the gating grid voltages are maintained at the closed values. Once the event is recorded, the GGD closes the gating grid by restoring the voltage difference between adjacent wires, as described above. This action stops additional electrons from reaching the anode and ensures a clean, time-correlated signal for each individual event. Because ion mobilities are orders of magnitude slower than those of electrons, most ions near a closed gating grid are slowly attracted to the negatively biased wires and neutralized.

Previously, the solid-state chip resistors used often failed within a day under beam conditions. The resistor failures were not due to excessive total current or average power, which remained within acceptable limits. Instead, the issue stemmed from uneven current density distribution within the resistors, which proved to be a more critical factor affecting their reliability. For the this latest campaign, we modified the the gating grid driver by employing 5Ω high current surge resistors for R_n and R_p in the circuit.

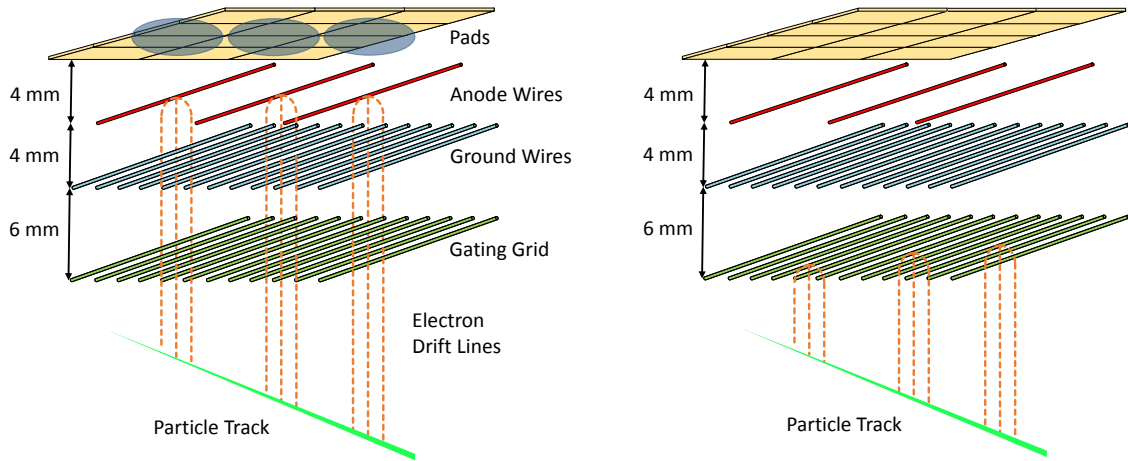


Figure 6.7 An illustration of how the GG enforces the trigger conditions from [14], to prevent electron capture on the pad plane via the closed configuration, until the trigger is met. The closed configuration (right) doesn't allow electrons to drift to the pad plane. In the open configuration (left), they can pass through.

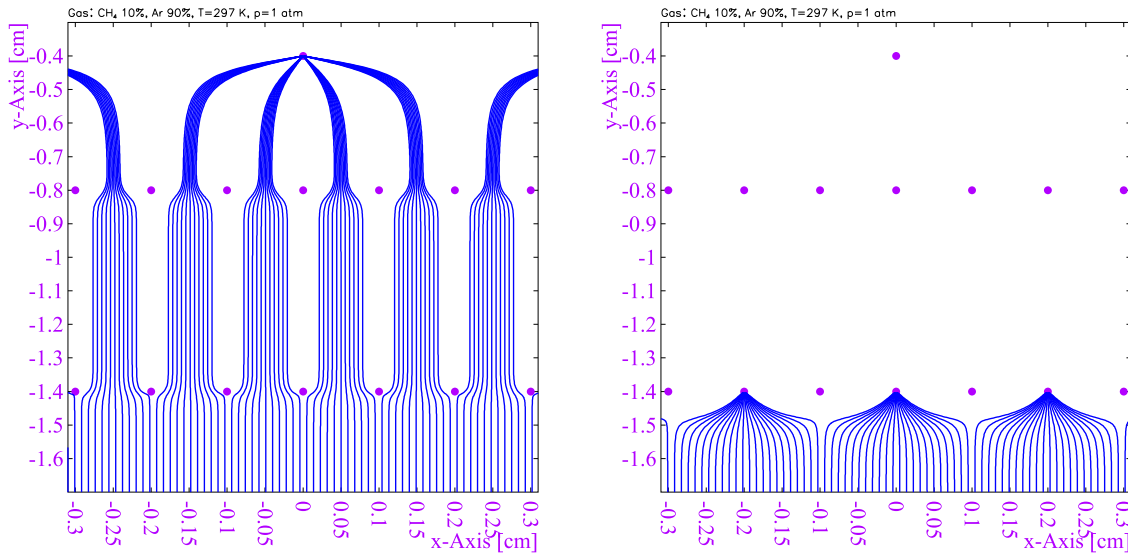


Figure 6.8 Electric field lines on the GG wires in the open (left) and closed (right) configurations from [14]. When closed, drifting electrons will terminate on the wires and not make it into the amplification region.

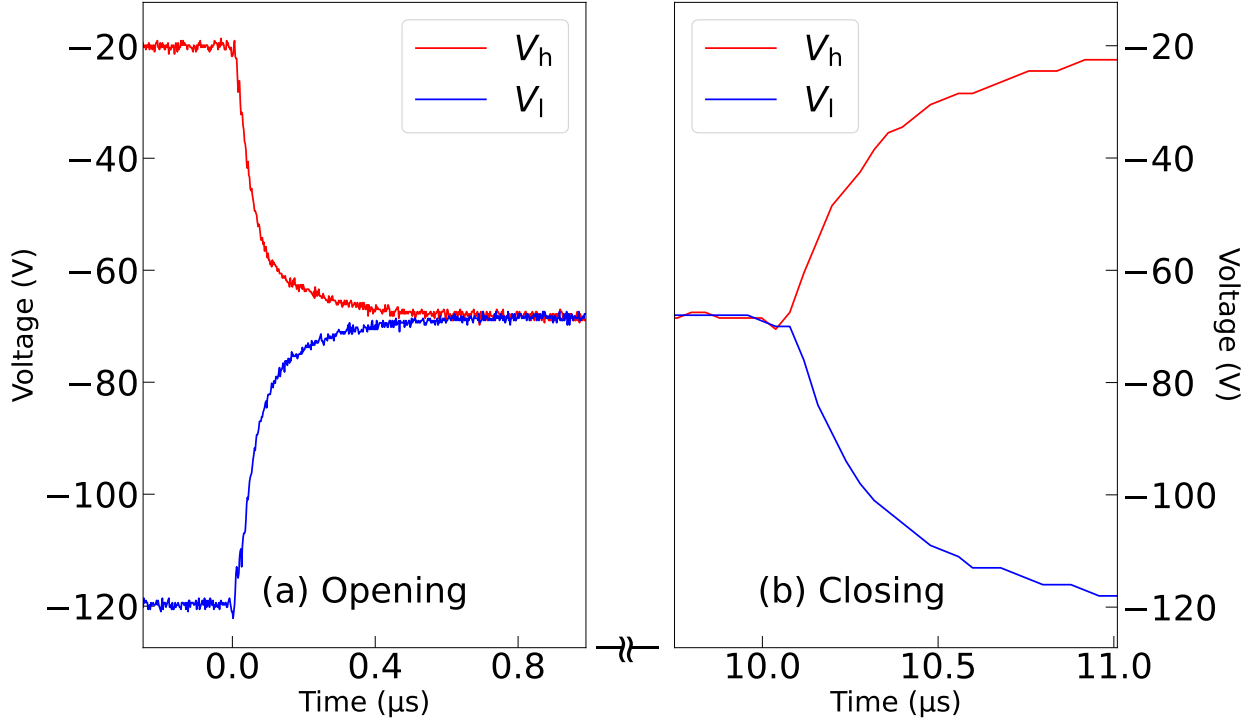


Figure 6.10 Oscilloscope trace of the V_h and V_l during the transition of the gating grid from closed to open state (left panel) and from open to closed state (right panel) for a normal event. Closed state corresponds to V_h and V_l at their nominal values of -20 V and -120 V, respectively. Open state corresponds to the state when the differences between V_h and V_l are close to zero.

These new surge resistors spread the current over a larger physical area, which eliminated localized overheating and eliminated the problem of resistors R_n and R_p failing. The GG worked perfectly after this change, with a single GGD lasting through both of the 2024 campaigns. Figure 6.10 shows an oscilloscope trace of the voltages V_h and V_l across the gating grid during the transition of the gating grid from closed to open state (left) and from open to closed state (right) for a normal event. The closed configuration corresponds to V_h and V_l at their nominal values of -20 V and -120 V, respectively. The open configuration corresponds to the state when the difference between V_h and V_l is close to zero. The gating grid opens in approximately $0.4 \mu s$. Because the readout time for the data acquisition greatly exceeds $0.4 \mu s$, we are limited to beam rates of the order of 10 kHz and there is no reason to reduce the restoration time of the wire voltages for the closed configuration. We restore those voltages for the open configuration with larger resistors and lower current values.

CHAPTER 7

ANALYSIS

In this chapter, the steps involved to process the signals collected in the $S\pi$ RIT TPC from raw data to obtain reconstructed particle tracks are discussed [13]. The analysis techniques for this Xe + Sn experiment are similar to those used to analyze a set of measurements of Sn + Sn systems performed at the RIKEN Nishina center in 2016 using. Both utilize the $S\pi$ RITROOT framework [66], based on the FAIRROOT framework [67]. Since those have been described in detail in [11, 16, 59], our discussions of certain aspects such as calibrations and corrections are basically the same. Thus, the present discussion focuses on details that are specific to the Xe + Sn system, with mention of the other aspects to provide a more complete picture for the reader.

7.1 Reconstruction of Data

This section will overview some of the important steps in converting the raw recorded data into ROOT files which are able to be analyzed, as well as some important calibrations and corrections.

7.1.1 GGD Noise subtraction

When the trigger system determines that an event should be recorded and the GG opens to allow charge to pass to the gating grid, the opening from the GGD generates noise on the pad plane. This noise has a characteristic shape that depends on the gating grid impedance. Figure 7.1 shows the gating grid noise for both the 2016 experiment (panels a-b) and the 2024 experiment (panels d-e). The improvements to the GGD removed the possibility of the resistors failing, and there was also some decrease in the GGD noise. This characteristic noise of the improved GGD is the same for all events and is routinely subtracted as described in [65, 14].

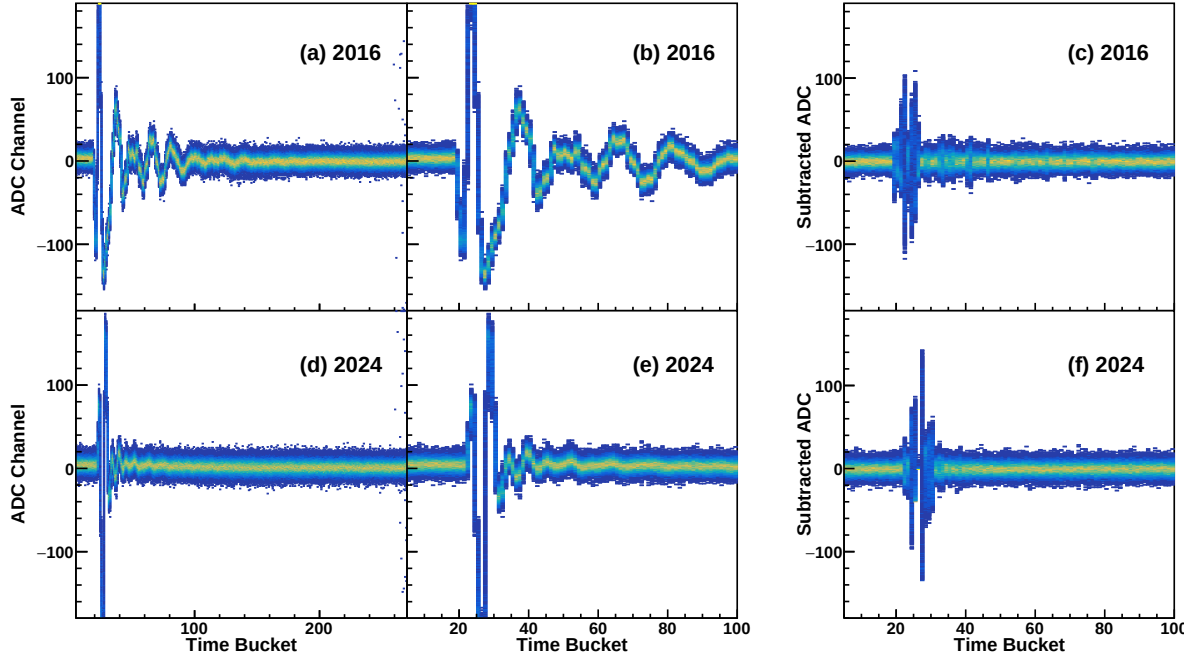


Figure 7.1 Noise spectra of a typical pad for 2016 (top) and 2024 (bottom) campaign induced when the GGD switches the GG from the closed to open configuration. Panels (a) and (d) show a compressed time scale of 270 time buckets (each time bucket corresponds to 40 ns) while panels (b) and (e) are expanded to a maximum of 100 time buckets so that oscillations of the noise signals can be seen more clearly. Panels (c) and (f) show the noise signal after subtraction of the average ADC values in each time bucket.

As discussed Section 6.3.2, a rapid opening of the grating grid is key to obtaining information from tracks that go upwards in the TPC. The GGD succeeds to open up the gating grid very quickly, taking approximately 400 ns.

7.1.2 Pulse Shape Analysis (PSA) and Tracking

The first step in data reconstruction is to convert the raw data from the TPC, which are just electronic pulses recorded by the pad plane pads, into a format that can be analyzed in ROOT. This process is referred to as decoding, and involves digitizing the pulses from each pad. Pulse Shape Analysis (PSA) is then performed to identify how much charge was collected over a pad within a given time; larger amounts of charge will result in a proportionally larger pulse. Each pulse is then classified as a *hit*, where the location of the pad on the pad plane gives information about the X- and Z-positions, and the time of the pulse within an event readout is converted into the vertical

Y-position. Because the readout is performed at fixed time intervals of 40 ns called time buckets, the starting time of the pulse and the physical drift velocity of the electrons inside the TPC can be combined to obtain the vertical position of that part of the track in the TPC.

During each bombardment, the gas purity of the P10 gas improved with time causing a slow increase in drift velocity. To obtain accurate position information, we measured the drift time experimentally using tracks which reflect from the bottom of the TPC field cage, since the distance from the bottom of the field cage to the GG is known. The time dependence of the measured drift velocity as a function of run number for the two campaigns can be seen in Figure 7.2, where gaps in the measurements correspond to specialized runs that were not performed to measure particle spectra (such as pulser calibrations). The nominal drift velocity is around 5.4 cm/ μ s for electrons within a pure P10 detection gas at a pressure of 800 Torr. The observed deviations from this value are probably caused by trace impurities in the counter gas, that is diluted over time by flowing pure P10 gas into the TPC while removing an equal volume of gas with the gas handling system. Some larger changes in the drift velocity are seen for collisions induced by the ^{136}Xe beam. These corresponded to changes in the gas flow rate.

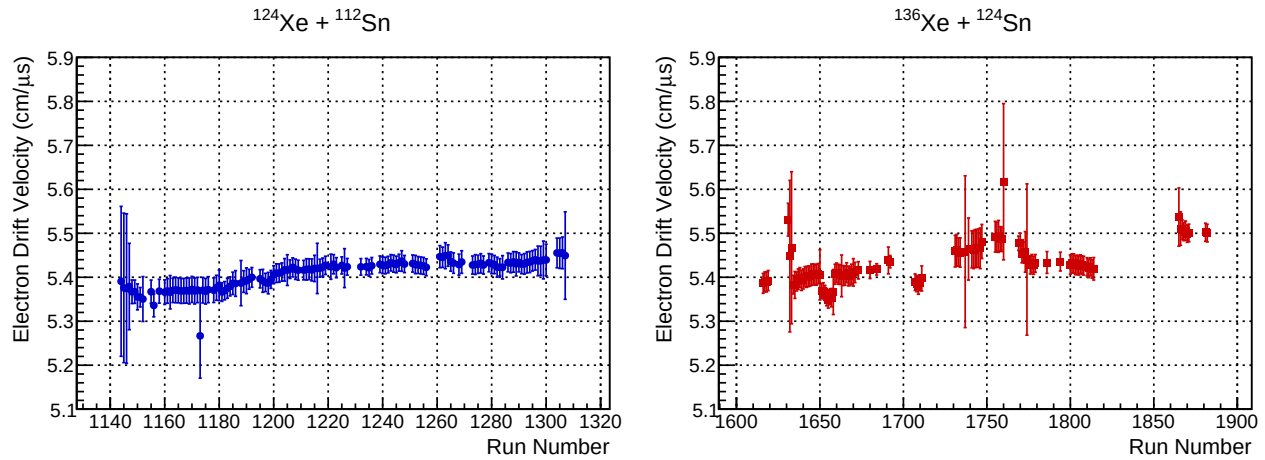


Figure 7.2 Calculated electron drift velocities as a function of run number for the two systems, using particles which reflect from the bottom of the field cage. Error bars are based on the number of events used in the calculation.

After the raw data has been converted into hits as described above, *Helix* tracking is performed [66] to identify preliminary tracks from collections of hits. A Riemann track finding algorithm

is used to map from Cartesian coordinates to a Riemann sphere, which is useful in this case because the charged particles follow helical paths under the influence of the SAMURAI **B**-field. The results of Helix tracking are used as an input for GENFIT [68], which provides an accurate reconstruction of the particle momentum by taking into account detailed information such as energy loss in the detecting gas. GENFIT tracks are then combined to find the vertex of the event using the RAVE software [69]. This allows one to distinguish between interesting events that originate from collisions of the beam with the target from uninteresting reactions of the beam with the detection gas or materials upstream in the beamline, so that latter background events can be discarded. We return to discussion of details of the analyses of the event vertex in Section 7.3.1.

7.1.3 TPC Anode Gain Calibration

To calibrate of the track ionization events measured in the TPC, we pulse the ground plane with a standard range of pulser voltages. This pulses the entire pad plane with the same charge signal, and the responses of all of the electronic channels to that charge are thereby measured. This calibrates the response of each channel consistently and thereby removes inconsistencies in calibration for specific electronic channels that are illustrated by the left panel of Figure 7.3. Here the height of the pulser signal is given by the ADC on the x-axis, and the supplied pulser voltage is shown on the y-axis. By fitting a polynomial to the output of particular channel, we obtain a calibration line. After optimizing the gain of the electronics to provide the desired dynamic range of the TPC for one channel, one then modifies the gain of the other channels as necessary so that their calibrations to the same voltage input are consistent. The result of applying this calibration can be seen in the right panel.

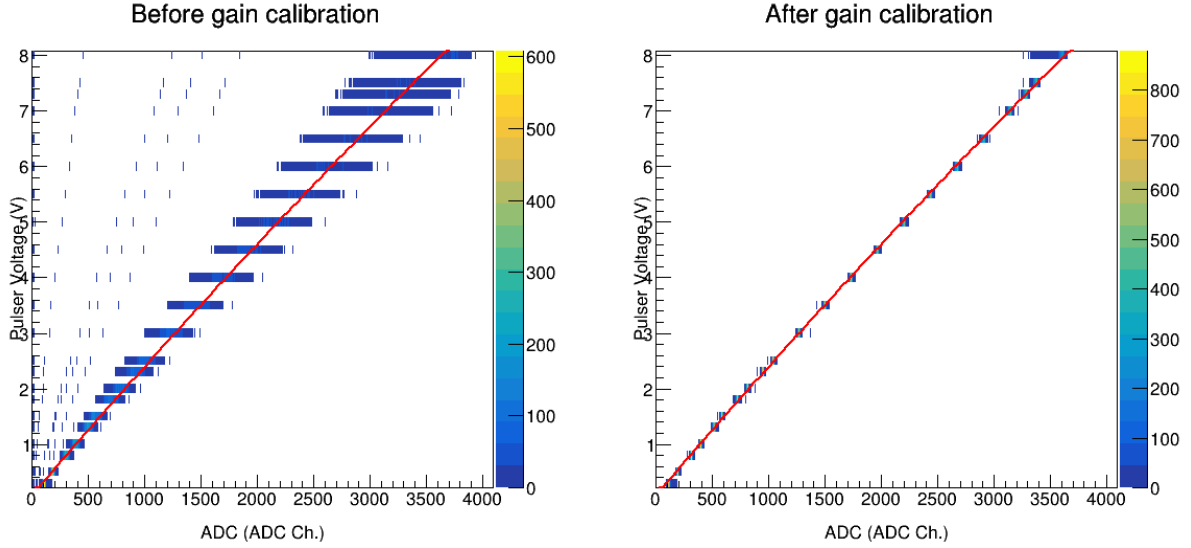


Figure 7.3 Pulser ramp signals before gain calibration (left) and after (right).

7.1.4 Beam Drift Chamber (BDC) Projections

The Beam Drift Chamber (BDC) is an ancillary detector and a key component of the Samurai detection system [70]. During our experiment, a pair of these detectors were located upstream of the TPC and outside of the SAMURAI magnet. Extremely precise information about the position and direction of the beam is provided as it passes through the BDC, to a precision of $120\ \mu\text{m}$ [71]. This information can be extrapolated to the target location at the TPC entrance to obtain the interaction point of each beam particle when it hits the target, along with values of the beam momentum and incident angles. The estimated errors on the projected values for the beam momenta, position, and incident angles are explained in more detail in [11]. At normal beam rates, we do not use pads that lie within a 15 cm radius of the beam spot on target because pads inside this radius are very often saturated and provide distorted tracking information. If the minimum distance from the extrapolated track to this beam interaction point is less than 20 mm, the BDC beam interaction point is added to the track when the momentum of the track is evaluated. If distance from the extrapolated track to the beam interaction point is larger than 20 mm, we reject the track as poorly characterized and likely nonphysical.

Table 7.1 Coordinate offsets between the BDC and TPC frame of references.

Beam-Target System	x_{offset} (mm)	y_{offset} (mm)
$^{124}\text{Xe}+^{112}\text{Sn}$	-4.62	-181.8
$^{136}\text{Xe}+^{124}\text{Sn}$	-2.70	-181.8

For the $S\pi\text{RIT}$ TPC coordinate system, the origin point ($x = 0$, $y = 0$, $z = 0$) is based on the physical edge of the pad plane just inside the entrance window. Because the BDC detectors are centered on the beamline, an offset must be applied so that their projection information is in the same frame of reference. These offsets are obtained by comparing vertex information between the TPC and BDC, and are listed in Table 7.1.

We note that our extrapolations of the beam path from the measured location in the BDC to the target ladder includes the steering by the fringe SAMURAI magnetic field. It is important to supply accurate interaction points to the GENFIT track fitting algorithm. Adding the projected point of interaction in the target provides a very accurate constraint on the track that greatly improves the accuracy of the reconstructed track momentum. In Figure 7.4, the initial result for the reconstructed track momenta is shown on the left panel. For tracks that pass criteria which are specified in Section 7.3.2, we perform refitting after including the BDC reference point as mentioned above. The result on the particle identification of including these BDC reference points is shown in the right panel. Clearly, both the separation between particle lines and the resolution of their distributions show marked improvement. In particular, the separation between the π^+ and p lines is greatly increased, which allows for better selection of pion tracks when constructing the final spectra.

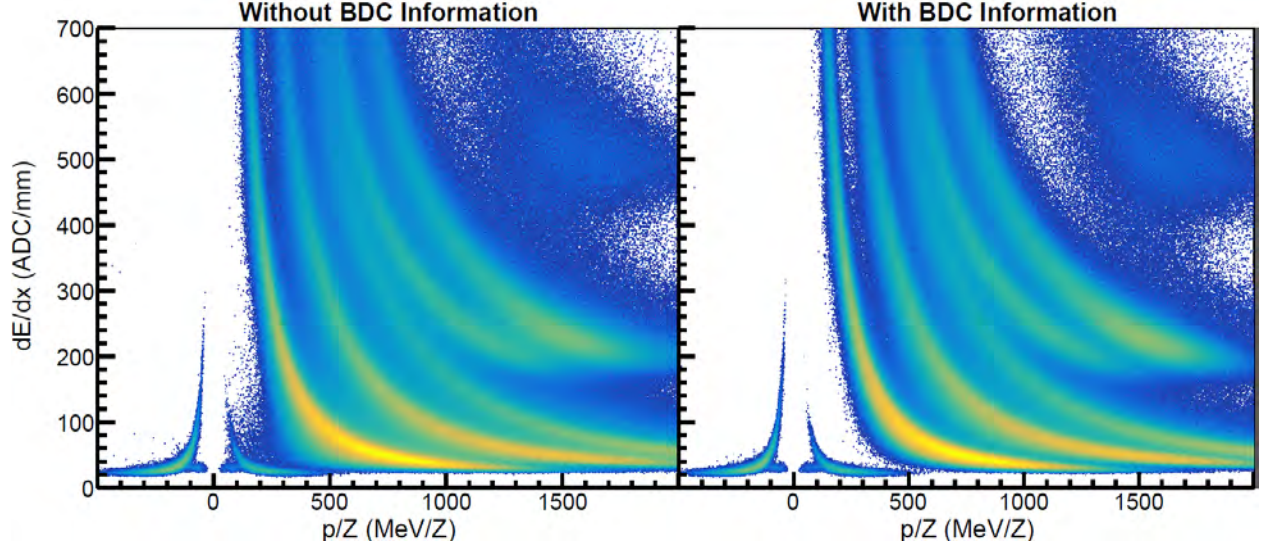


Figure 7.4 PID for charged particles produced with the ^{124}Xe beam without incorporating the BDC projection information (left), and with (right). The momentum resolution using the BDC is improved, as well as particle separation which can be seen especially between the p and π^+ lines.

7.1.5 Space Charge Corrections

Similarly to the case described in Section 4.4 of Part I of this thesis for the AT-TPC, the $S\pi\text{RIT}$ TPC also suffers from distortions to the recorded particle tracks due to the build up of positive ion space charge in the drift volume of the TPC which deflects the drifting electrons from their intended paths. While the negative e^- are quickly drifted upwards by the field cage, the positive P10 ions travel downward approximately 1000 times slower. In the $S\pi\text{RIT}$ TPC, the space charge manifests itself as a vertical sheet of positive charge along the path of the unreacted beam, as is illustrated by the red plane in Figure 7.5. This is a result of the approximate orthogonality between the direction of the beam and the applied drift $\mathbf{E}$ -field, which drifts the positive ions downward, as opposed to the AT-TPC where the space charge manifests as a cylinder because the beam and drift $\mathbf{E}$ -field are parallel.

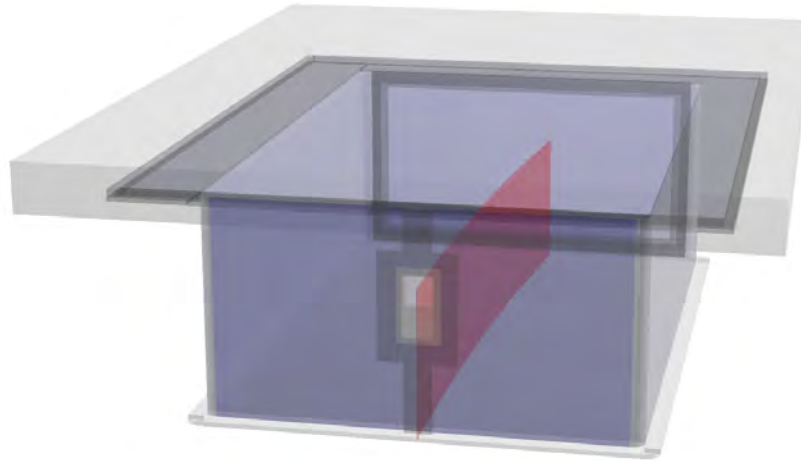


Figure 7.5 A diagram of the space charge sheet location (red) that results from the ionization of the unreacted Xe beams in the S π RIT TPC.

Shown in Figure as solid lines 7.6, are the actual paths of two fast charged particles produced in the collision. The dashed lines show the positions of the tracks deduced by digitizing the electrons ionized by these fast particles. The electrons in these dashed tracks are displaced from the actual tracks (solid lines) by the electric field of the space charge in the TPC. The largest space charge effect in the S π RIT TPC is the shift of leftward going tracks (blue) to the right, and rightward going tracks (green) to the left. The negative charges in both sets of tracks are attracted by the positive space charge sheet in the middle of the TPC. Because the ionic space charge moves downward there is also a small vertical component to the space charge field. Combining the drift field from the TPC and the space charge sheet results in a position dependent effective field represented in direction and magnitude by the length of the black arrows. The difference between the x-position of the actual event vertex and the projection of a track to the target plane is called ΔV_x , and is specifically shown for the left-going track in the figure inset. If this shift is not corrected along the track, such shifts to the measured tracks will result in incorrect momenta from the tracks and incorrect assessments of track quality cuts that will be discussed in Section 7.3.2. Determination of particle momentum relies on measurement of the track curvature inside the SAMURAI field, so an incorrect value for the momentum components will also be obtained. This space charge effect acts on all electrons, but it is somewhat greater for particles emitted in the beam right direction,

especially for the electrons in tracks that are close to the positive ion sheet.

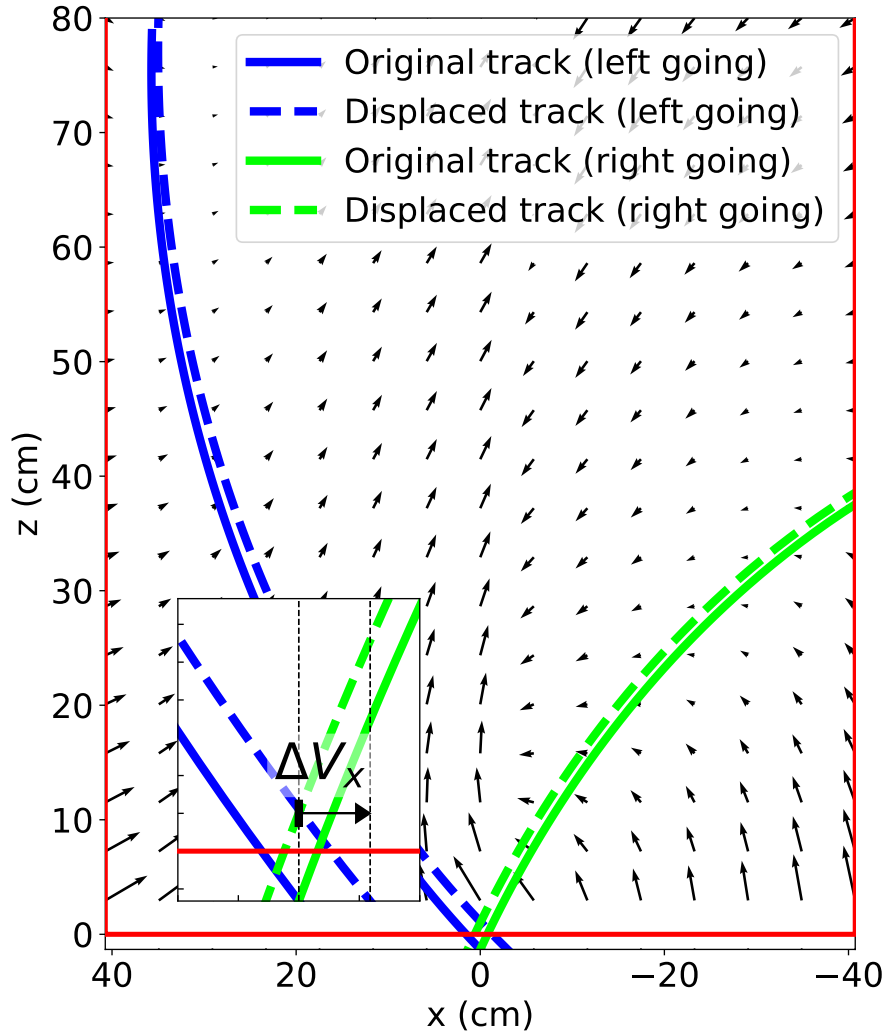


Figure 7.6 Top-down view of two tracks reconstructed inside the S π RIT TPC, with arrows representing the combined electric field from the TPC field cage and the space charge sheet. The actual tracks (solid lines) meet at the true reaction vertex, where the distorted tracks (dashed lines) are displaced. Adapted from [15].

A simpler illustration of this effect on the event vertex reconstruction can be seen in Figure 7.7, where the measured location that an individual right-going track crosses the target plane is displaced to the left of the location expected from external beam position measurements using the BDC. Thankfully, this accurate external BDC, combined with the track analysis provides us the information needed to accurately correct the track position at that location for the space charge effects and provides information to constrain the space charge displacement all along the track better

defined defined. Nonetheless, the displacement of charges by the SAMURAI **B**-field requires these shifts contained in the measured track positions to be carefully corrected using the following solution to the Langevin Equation [33]:

$$\mathbf{u} = \frac{e}{m} \tau |E| \frac{1}{1 + \omega^2 \tau^2} \left(\hat{E} + \omega \tau [\hat{E} \times \hat{B}] + \omega^2 \tau^2 (\hat{E} \cdot \hat{B}) \hat{B} \right). \quad (7.1)$$

Here $\mathbf{u}$ is the drift velocity, e and m are the electron charge and mass, τ is the mean time between collisions, and $\omega = e|B|/m$ is the electron cyclotron frequency. The method for fitting the magnitude of the space charge sheet is to find the relationship between charge density and beam intensity that results in agreement between the ΔV_x of left-going and right-going tracks. More detail about the space charge correction procedures for the S π RIT TPC can be found in [15].

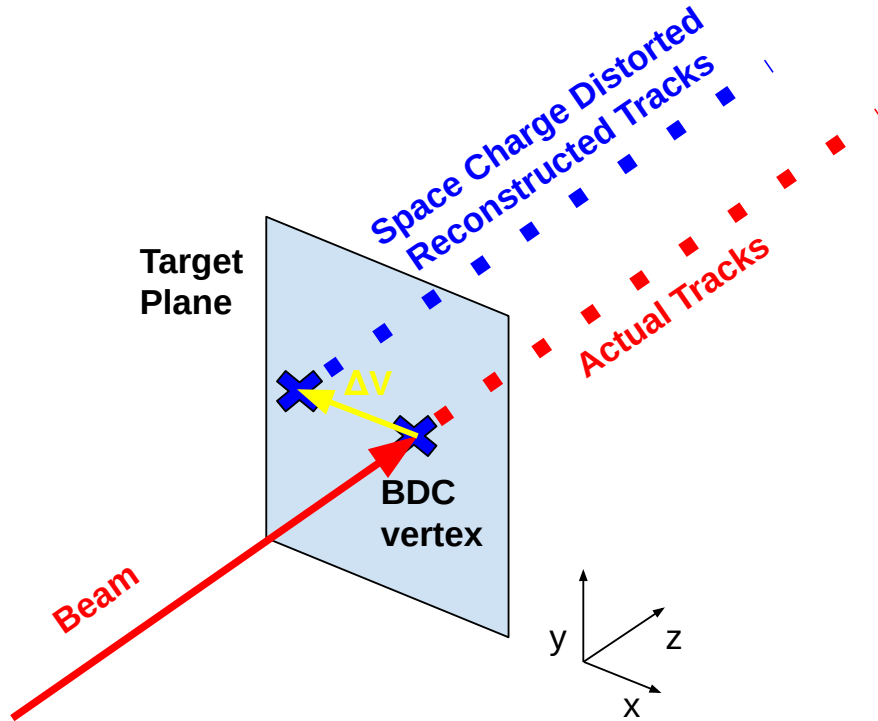


Figure 7.7 A graphic that shows the effects of space charge distortion on vertex reconstruction. The actual location of the Xe beams on the Sn targets is known thanks to information from the BDC.

7.2 Cocktail Beam Calibration

To check the momentum reconstruction calibration of the S π RIT TPC, a primary cocktail beam consisting of light charged particles was employed. The target ladder was set to an empty position

so that the beam could enter the TPC unimpeded in energy. The PID for this cocktail beam is obtained by plotting the energy loss vs. magnetic rigidity as shown in Figure 7.8, where the two most prominent peaks correspond to ${}^4\text{He}$ and d . The comparison of measured momentum to predictions from the LISE++ software [26] are seen in Table 7.2.

Table 7.2 A summary of the predicted and measured values for the cocktail beam calibration.

Particle	Energy (AMeV)	LISE++ (MeV/Z)	Measured (MeV/Z)	Difference
d	261.270	1500.07	1463.71	-2.42 %
${}^4\text{He}$	264.071	1499.49	1450.51	-3.27 %

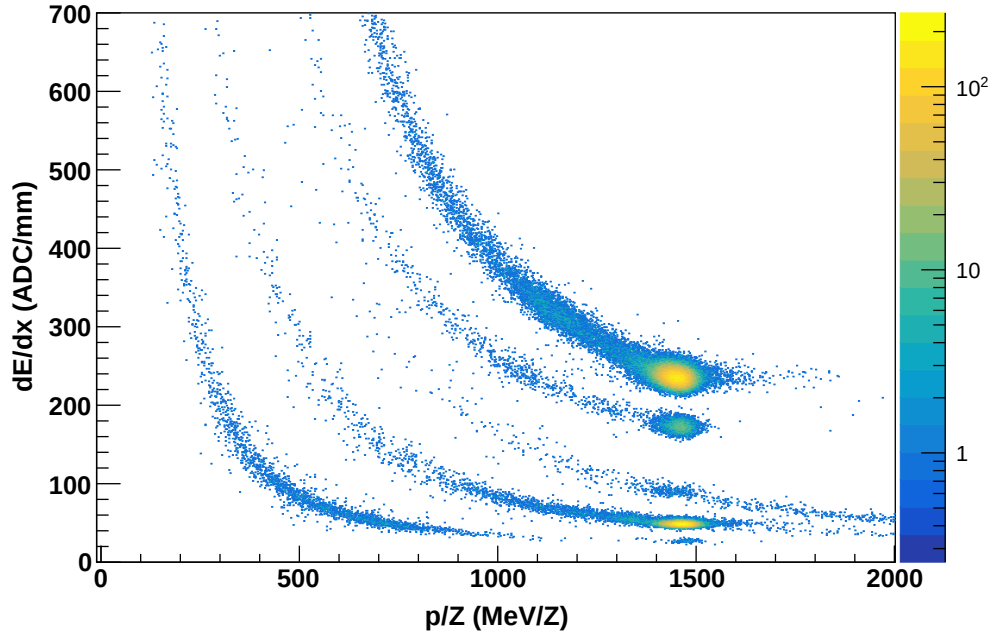


Figure 7.8 PID for fragmentation of the stable primary beam into stable hydrogen and helium isotopes producing beams with fixed rigidity that are used for checking the calibration of the TPC analyses. The three main peaks near $p/Z = 1500$ from top to bottom are ${}^4\text{He}$, ${}^3\text{He}$, and deuterium. Much smaller peaks are observed for tritium (middle) and protons (bottom).

7.3 Particle Identification (PID)

In this section, I will explain how events and particle tracks are chosen to make spectra. While the S π RIT TPC is capable of reconstructing most light charged particles that travel through its detection volume, we are only interested in those particles which come from collisions of the Xe

beams with the Sn targets. We must reject events resulting from other origins. Furthermore, the reconstruction accuracy of particle tracks within an event depends on certain track quality factors such as the overall length of the track. Only those tracks which have sufficiently accurate information resolution are selected to contribute towards particle spectra. Simulations are used to calculate the efficiency for these quality factors being satisfied and final data are correct for the efficiency.

7.3.1 Event Selection

The main criteria for which events are selected is the condition that the collision is found to have originated from the target. Figures 7.9 and 7.10 show the reconstructed vertex distributions for all events over several runs for both systems, without any selection criteria applied to the data.

In Figure 7.9, the Z component of the reconstructed vertex for each system with space charge corrections is shown. The most prominent spike near $z = 0$ comes from reactions on the target. This peak is fit for each system, and a $5\text{-}\sigma$ window centered on the fit value is used to select events as “on-target”. Entries at $z = 0$ come from interactions with the entrance window of the TPC and at $z > 0$ from reactions on the TPC target gas, and entries for $z < 0$ come from reactions on materials upstream in the beamline. Even though such collisions will produce charged particles, we are only interested in target collisions. Table 7.3 shows a summary of the reconstructed vertex from RAVE compared to the value measured from the physical position of the target ladder. When we apply space charge corrections to the track, the ion optics calculations prefer the vertex for reactions on the target to be shifted by ~ 2.5 mm upstream. This shift appears to be an unexplained artifact of the space charge task. Since this task is sensitive to the spatial distributions of ions in the region of the entrance window and the actual track lengths are more than 100 times longer than this shift, this discrepancy is overall inconsequential for construction of the pion spectra.

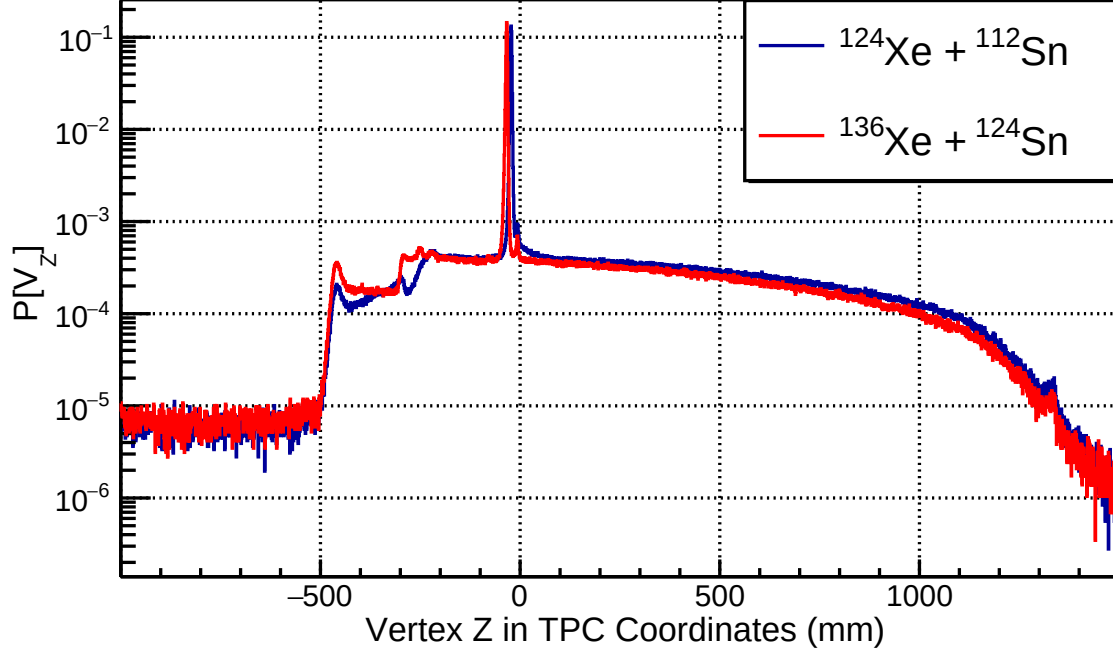


Figure 7.9 Probability distribution of the z-component of reconstructed event vertices measured for measurement to study the $^{124}\text{Xe} + ^{112}\text{Sn}$ system (blue) and the $^{136}\text{Xe} + ^{124}\text{Sn}$ system (red). The large peaks at $z \approx 19.5$ mm (blue) and at $z \approx 31$ mm (red) correspond to events when a projectile nucleus collides with a target nucleus. The tiny peaks at $z \approx 0$ mm occur when a projectile collides with a nucleus in the entrance window of the TPC. Data for $z > 0$ occur when a projectile collides with the P10 gas in the TPC that is downstream of the target. The peak at 1.3 m corresponds to reactions on the downstream window of the TPC.

Table 7.3 Summary of the expected and reconstructed values of the event vertex for each campaign.

Beam	Target	Target center (measured) (mm)	Reaction Peak (RAVE) (mm)	Reaction Peak Space Charge (RAVE) (mm)	Sigma (mm)
^{124}Xe	^{112}Sn	-19.48	-19.88	-22.45	1.78
^{136}Xe	^{124}Sn	-30.82	-30.95	-33.67	1.72

In Figure 7.10, the x- and y-components of the reconstructed RAVE vertex at the target plane location are shown before applying the on-target selection to the z-component. It is clear that a significant number of recorded events occur from other sources, compared to the central peak. As mentioned above, a $5\text{-}\sigma$ cut on the RAVE z-component is applied to determine events that have resulted from target collisions. Figure 7.11 shows the x- and y-components of the surviving events after this selection, where the vast majority of the undesired events have been removed. The target

dimensions are 30 mm x 40 mm, and the coordinate system is such that the target is centered at (0 mm, -185 mm). The physical edges of the rectangular target are now visible by the sharp cut-off in statistics. Events which have reconstructed vertices outside of the target boundaries are further excluded with cuts to the x- and y-components in addition to the $5\text{-}\sigma$ window.

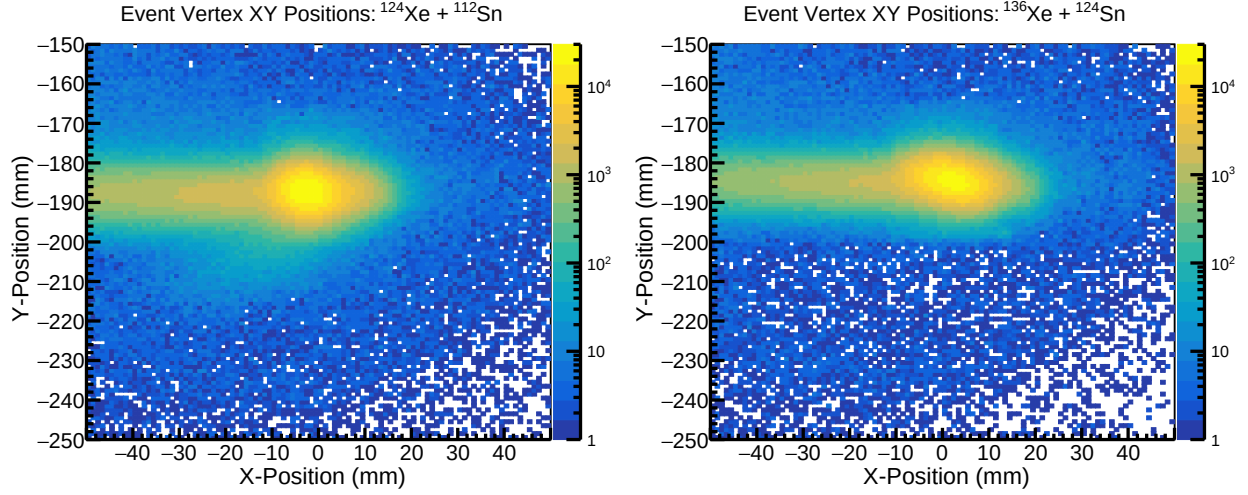


Figure 7.10 Distribution of the x- and y-components of ALL reconstructed event vertices for the $^{124}\text{Xe} + ^{112}\text{Sn}$ system (left) and $^{136}\text{Xe} + ^{124}\text{Sn}$ system (right), before applying the on-target selection to the z-component. It is clear that a significant number of events are from collisions with other materials in the beamline or the TPC detection gas.

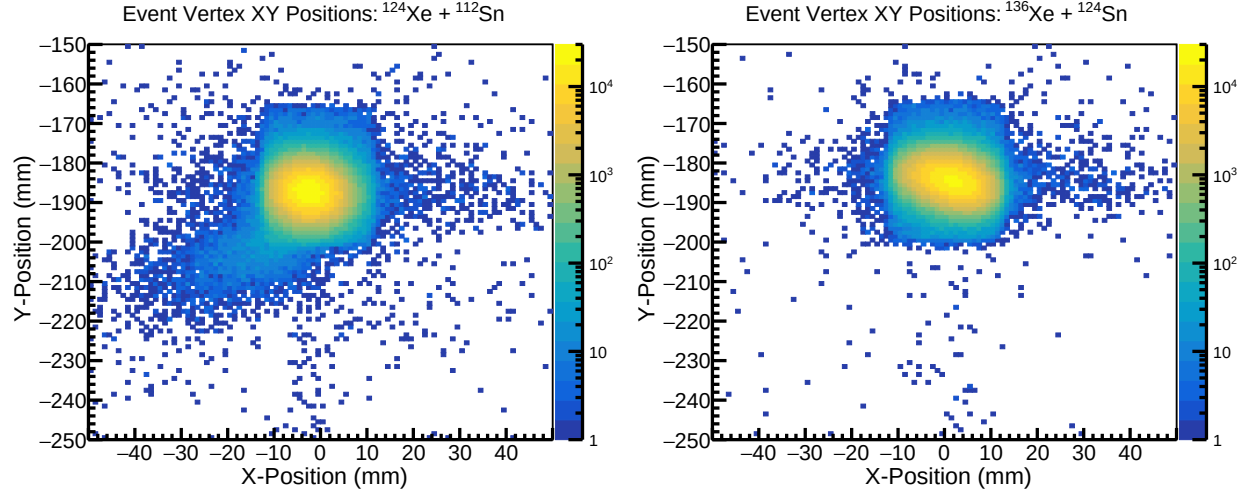


Figure 7.11 Distribution of the x- and y-components of reconstructed event vertices which have been determined to be on-target by applying a 5σ cut to the z-position around the expected target location for the $^{124}\text{Xe} + ^{112}\text{Sn}$ system (left) and $^{136}\text{Xe} + ^{124}\text{Sn}$ system (right). The target dimensions are 30 mm x 40 mm, and the coordinate system is such that the target is centered at (0 mm, -185 mm). The physical target boundaries can be seen by the apparent edges in the plots.

7.3.2 Track Selection

Charged particles produced in the collisions are emitted over a wide range of angles in the nucleus-nucleus center of mass frame. Track information is recorded in each pad the particle traverses over and recorded by the readout software. When charge deposited in a pad saturates its electronics channel, that channel goes dead for the remainder of the event. This saturation causes a characteristic distortion of the signal that is recorded in the data analysis, marking that channel as dead or distorted. Dead or distorted channels are corrected by interpolating charge deposition from the lesser charges deposited in neighboring pads in a process called de-saturation [72].

The efficiency and the resolution of the TPC data are not uniform throughout the TPC and this variation is taken into account during analysis. For example, the rectangular geometry of the $S\pi\text{RIT}$ TPC means that particles that end up in the beam left or right corners of the TPC, but do not overlap strongly with the beam trajectory will have longer tracks over more pads. This extra pad information will provide better momentum resolution. On the other hand, particles emitted at large scattering angles upward or downward in the TPC will deposit charge in fewer pads, resulting in larger uncertainties in their momenta. Thus, the acceptance and resolution of tracks measured

in the TPC must guide the analysis and ultimately be taken into account.

The directions of the angular momentum for measured events lie in a plane that is perpendicular to the beam. In this plane, the probability for collisions with different azimuthal directions of the total angular momenta are essentially the same. The effects of the asymmetry of the magnetic field on the experimental resolution are managed by sampling the measured events with respect to the orientation of their total angular momenta about the z axis. In the vertical SAMURAI magnetic field, particles with angular momenta that are closer to horizontal are measured with higher resolution.

To enable discussion of our analyses and address these issues, we describe the direction of particle emission from the target using the following definitions:

$$\phi = \tan^{-1} \left(\frac{p_y}{p_x} \right) \quad [-\pi \leq \phi \leq \pi], \quad \theta = \tan^{-1} \left(\frac{p_T}{p_z} \right) \quad [0 \leq \theta \leq \pi], \quad (7.2)$$

where ϕ is the azimuthal angle in the xy-plane and θ is the polar angle with respect to the beam. We denote the laboratory frame momentum components by $p_{x,y,z}$ and the transverse momentum component by $p_T = \sqrt{p_x^2 + p_y^2}$. Here, a polar angle of $\theta = 0^\circ$ corresponds to particles traveling parallel to the beam in the z-direction, and $\theta = 90^\circ$ would correspond to particles traveling perpendicular to the beam. A distribution of these angles for all particle species coming from on-target events is shown in Figure 7.12, where the band of $-40^\circ \leq \phi \leq 25^\circ$ represents tracks reconstructed towards beam left, and the bands $-180^\circ \leq \phi \leq -150^\circ \cup 160^\circ \leq \phi \leq 180^\circ$ represent beam right. The regions in between these bands with less statistics (or none at all) are from particles traveling up or down, where larger values of θ increasingly suffer from both short track lengths and parallelism to the magnetic field, leading to poor momentum resolution of the track.

In order to obtain precise constraints on the equation of state, it is essential to measure the momentum of the analyzed particle accurately. To achieve the required momentum resolution, only tracks from these left and right regions are used. An additional forward angle selection is made on the polar angle described in Section 8.2.1. It is useful to remember at this point that there is no significant bias in our experimental setup that prefers one azimuthal angle of the total angular momentum of the collision with respect to our magnetic field. The experiment provides a relatively

unbiased sample of events in which the total angular momentum in the x-y plane is uniformly distributed.

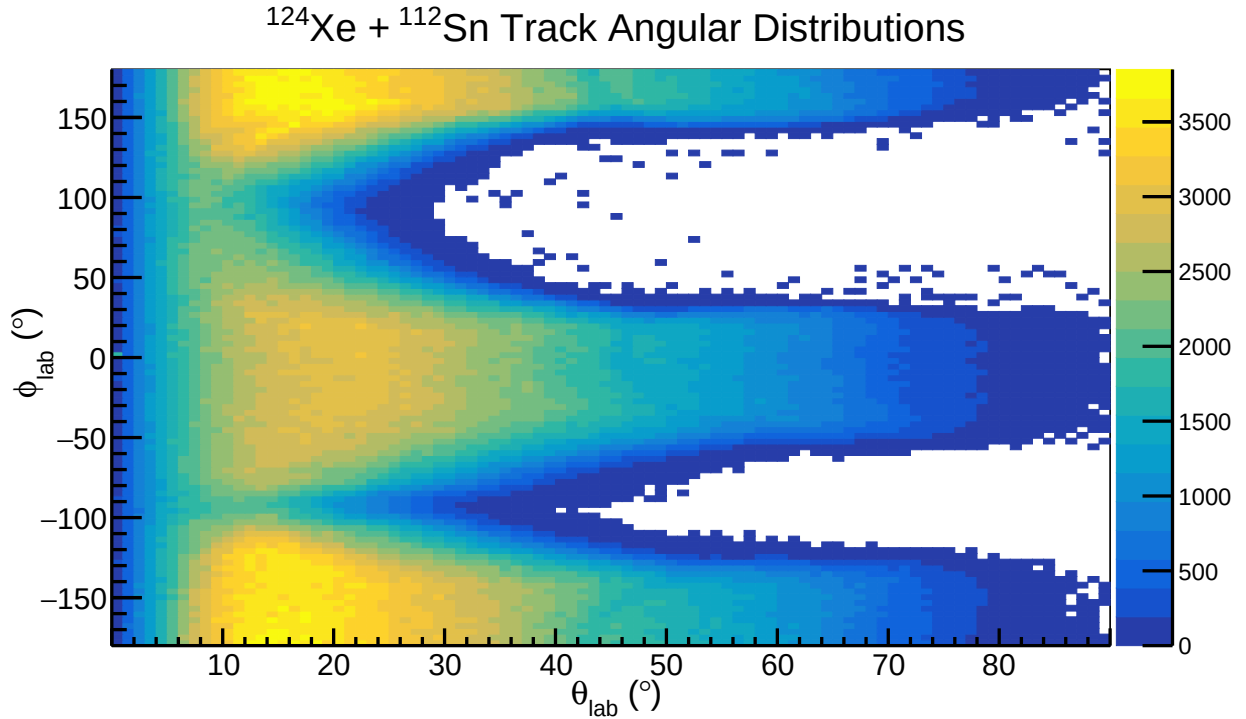


Figure 7.12 Distribution of target plane emission angles in the laboratory frame for all particles coming from collisions with the target. The long bands with greater statistics correspond to the beam left and right directions, which are regions within the S π RIT TPC that allow for more accurate momentum reconstruction of particle tracks.

Tracks of positively charged particles that scatter to beam left will subsequently bend right in the magnetic field, leading to longer track lengths. Positive particles that scatter right to the same scattering angle, then bend right and quickly exit the TPC. The short track lengths of these beam right particles reduces the efficiency and reconstruction accuracy. The negative consequences of this track length difference can also be seen in Figure 7.13, where the PID made from reconstructed left-going tracks have particle lines with slightly better resolution and separation than for the PID made from right-going tracks.

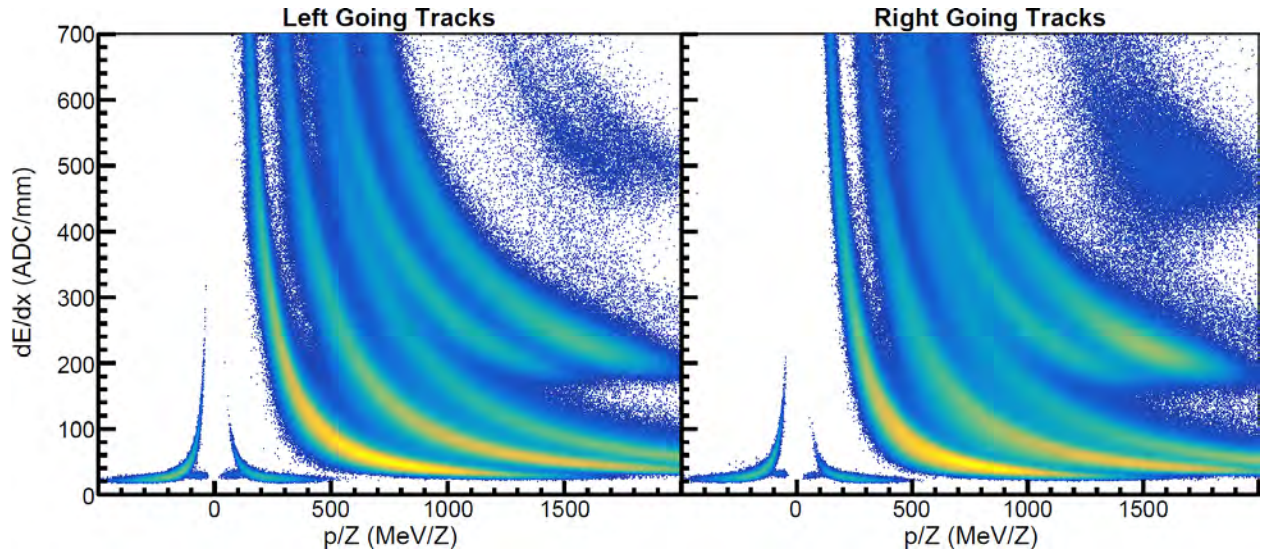


Figure 7.13 PID for charged particles produced with the ^{124}Xe beam for left going (towards beam left) and right going (towards beam right) tracks when they emerge from the target. There is slightly better resolution for leftward tracks, which can be seen most clearly between the proton and deuteron lines, and in the lithiums.

Each channel in the GET electronics used in the TPC readout employs a preamp. When the counting rates near the target are sufficiently high, electronic pile-up can saturate the preamp. After this occurs, the preamp output becomes distorted and often will cease until the preamp recovers. This prevents the acquisition of reliable data from regions that lie less than 20 cm from the target where electronic pile-up is likely. Since unreliable measurements can contaminate and render a good measurement useless, we exclude data at $r < 20$ cm and instead extrapolate the tracks from $r = 20$ cm to the target. Then, we place a threshold on the point of closest approach to the RAVE vertex for these extrapolated tracks. Figure 7.14 illustrates this point-of-closest-approach (POCA) metric. For a track to be considered as part of the collision, they must intersect a sphere of radius R centered on the reconstructed RAVE vertex location. In other words, the minimum distance d_i between the track and the vertex location must be less than the chosen R . A typical threshold is $POCA \leq 15$ mm. This metric is applied before including the BDC information in the track fitting, otherwise tracks would artificially constrained to satisfy this POCA cut.

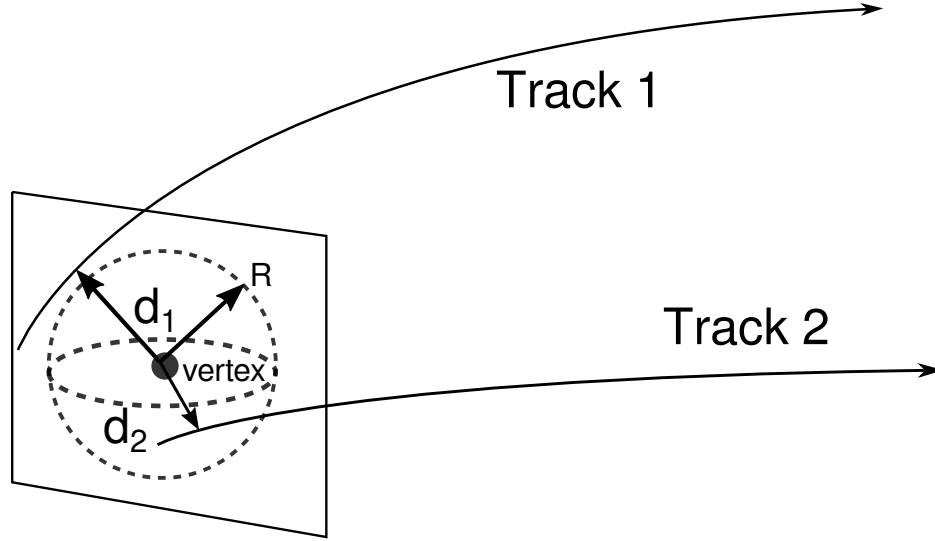


Figure 7.14 Representation of the point-of-closest-approach (POCA) metric for measuring individual track proximity to an event vertex. In this example *Track 1* fails the cut while *Track 2* passes.

Figure 7.15 illustrates the clustering metric, where the red line indicates the actual particle track position. “Clusters” refers to the groups of anode pads that have been combined to better calculate the average position of a track already found using simple pulse shape analysis. For a pad that has collected charge q_i with distance λ_i from pad center to the avalanche peak, the charge weighted average position is given by the clusters (dots). In order for the centroid to be well defined, clustering is done along the x-direction for parts of the track where $0^\circ < \theta \leq 45^\circ$, and along the z-direction for parts of the track where $45^\circ < \theta \leq 90^\circ$. The requirement of a minimum number of clusters for a track, such as 20, functions as both an effective track length cut and a track quality cut. Tracks with a larger number of clusters generally have better accuracy in their momentum reconstruction. Simulations are performed to understand the implication of this requirement (and others) about the overall detection efficiency and to confirm that the measured efficiency is acceptable. These simulations are described further in Section 8.1.3.

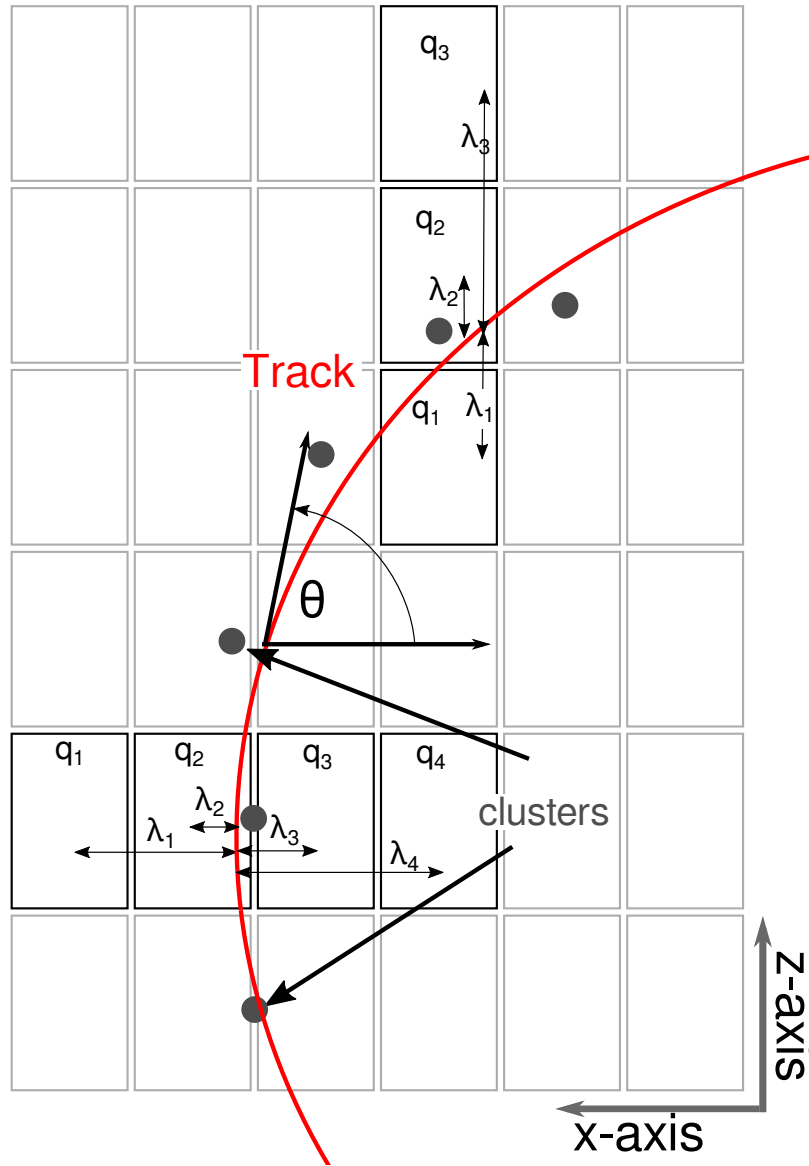


Figure 7.15 Illustration of the clustering technique for improving tracking position resolution. Figure from [16].

Table 7.4 A summary of the cut conditions used in event and track selection for particle distributions.

Event Conditions	Cut Values
$^{124}\text{Xe}+^{112}\text{Sn}$ Z-Vertex	-22.45 ± 5.34 mm
$^{136}\text{Xe}+^{124}\text{Sn}$ Z-Vertex	-33.67 ± 5.16 mm
Target width	$-15 \text{ mm} \leq x \leq 15 \text{ mm}$
Target height	$-205 \text{ mm} \leq y \leq -165 \text{ mm}$
Track Conditions	Cut Values
Number of clusters	> 20
Distance to vertex	≤ 15 mm
Left Azimuthal range	$-40^\circ < \phi < 25^\circ$
Right Azimuthal range	$-180^\circ \leq \phi \leq -150^\circ \cup 160^\circ \leq \phi \leq 180^\circ$
Forward Polar Angles	$\theta_{CoM} \leq 90^\circ$

7.4 Impact Parameter Determination

An important consideration for the centrality of a HIC is the distance between the trajectories of the two colliding nuclei in their combined center of mass system, referred to as the impact parameter and denoted by the variable b . Experimentally, we do not directly measure the impact parameter. Instead, we measure an observable, such as the multiplicity of detected charged particles, that depends on the impact parameter. Then we develop and calibrate an empirical relationship that connects the multiplicity to the impact parameter and use it to analyze events at the impact parameter values of interest. In this section, we discuss how this relationship is obtained and applied to select central collision events.

7.4.1 Charged Particle Multiplicity

In this dissertation, the charged particle multiplicity (N_C) refers to the number of particles detected in a collision. In a perfect detection system, it would equal the number of charged particles emitted in a collision. Generally, more central collisions have a larger value of N_C than more peripheral collisions. When calculating the charged particle multiplicity, N_C , each additional particle increases N_C by 1. Thus, a proton and an alpha both increase N_C by 1. Higher N_C correlates to the deposition of higher excitation energy because the emission of a particle removes both the kinetic energy of that particle and the separation energy with which that particle was bound to the

system. The normalized multiplicity distributions for each system $P[N_C]$ are shown in Figure 7.16. Peripheral events with very small N_C are rejected by the trigger system as discussed in Section 6.3.1, so there is little probability of measuring $N_C < 10$.

The total charge of the projectile and target system is $Z = 104$. If no negative pions are emitted, this sets a limit of 104 on the charged particle multiplicity. This would be achieved only rarely because the detection efficiency of the TPC is not perfect and the emission of no clusters with $Z > 1$ is highly unlikely. For very central collisions that produce the maximum multiplicity of tracks, many will have $Z > 1$ or will be moving too slowly to be properly detected and accurately reconstructed. Thus it should be expected that $P[N_C]$ drops to nearly zero after $N_C > 70$. The events that meet the trigger conditions have an average charged particle multiplicity of around 45. Events are considered for constructing the charged particle multiplicity distribution if they meet the on-target requirement of a 5σ vertex cut centered on the z-component of the RAVE value. Tracks in these events are then counted towards N_C if they satisfy $POCA \leq 15$ mm.

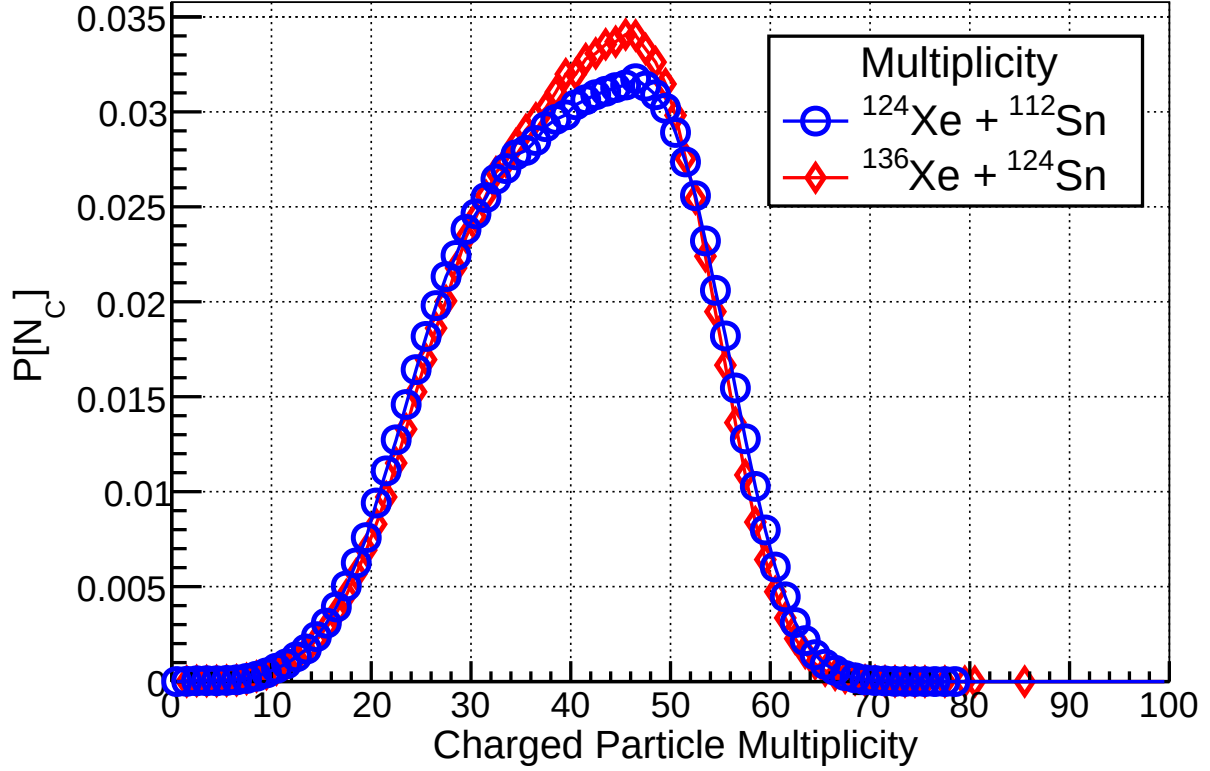


Figure 7.16 Normalized charged particle multiplicity distributions for both systems. The number of charged particles produced in a collision correspond to the centrality, where more central collisions produce more particles.

7.4.2 Impact Parameter and Central Collisions

The particle identification spectra shown in 7.13 reveal that hydrogen isotopes comprise the majority of particles that are counted in the light particle multiplicity. The comparative paucity of helium and heavier isotopes support the assumption that the multiplicity is monotonically related to the charge enclosed in the overlap region between the projectile and target nuclei. We assume that the *reduced* impact parameter ($\hat{b}$) is monotonically related to the charged particle multiplicity distribution via the geometrical relationship [73]

$$\hat{b} \equiv \frac{b(N_C)}{b_{max}} = \left\{ \sum_{N_C}^{\infty} \frac{dP(N_C)}{dN_C} \right\}^{1/2}. \quad (7.3)$$

By summing over $P[N_C]$ with Equation 7.3, we obtain the reduced impact parameter distribution shown in Figure 7.17. As expected, the number of measured charged particles increases with

decreasing impact parameter.

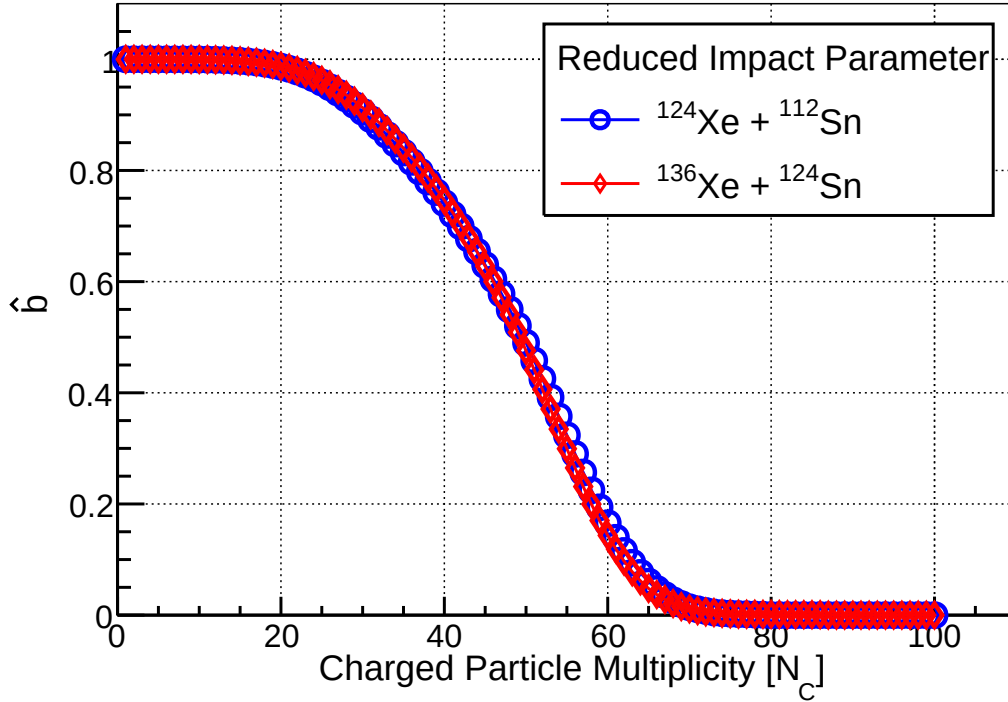


Figure 7.17 Reduced impact parameter, or the relationship between $\hat{b}$ and N_C , for both systems. Smaller values of $\hat{b}$ correspond to more central collisions.

A value of $\hat{b} = 1$ corresponds to peripheral collisions that occur at the largest possible impact parameter, or b_{max} . The full impact parameter relationship $b(N_C)$ is then obtained from scaling $\hat{b}$ by b_{max} , which is obtained in the geometrical limit where the absolute collision cross section σ is given by

$$\sigma(b_{max}) = \pi b_{max}^2 \Rightarrow b_{max} = \sqrt{\sigma/\pi}. \quad (7.4)$$

The collision cross section is calculated from experimental data using the ratio of target reactions to incident beam events and the target density:

$$\frac{\rho_{Sn} \cdot \sigma}{m_{Sn}} = \frac{\# \text{ target reactions}}{\text{incident beam}}, \quad (7.5)$$

where ρ_{Sn} and m_{Sn} are the area density of the Sn target and mass of the Sn nucleus, respectively.

The area densities of the targets employed in the 2024 campaigns were 243 mg/cm² for the

95.1% isotopically enriched ^{112}Sn , and 278.5 mg/cm² for the 97.4% isotopically enriched ^{124}Sn . Independently from the S π RIT TPC trigger system, a scintillating beam trigger (SBT) was placed upstream in the beamline that recorded each incident Xe beam particle regardless of whether a collision occurred.

Certain data runs were recorded using so-called *minimum bias* trigger settings, where the central collision high multiplicity trigger requirement in the Kyoto array was removed. This way, a more inclusive selection of events that included peripheral collisions could be obtained. Measured values of the reaction cross section and maximal impact parameters for these minimum bias runs are summarized in Table 7.5. The listed errors are obtained from the variance of measured cross section values as follows:

$$\delta b_{\text{max}} = \sqrt{\left(\frac{\partial b_{\text{max}}}{\partial \sigma_{\text{tot}}} \cdot \delta \sigma_{\text{tot}}\right)^2} = \frac{\delta \sigma_{\text{tot}}}{2\sqrt{\pi\sigma_{\text{tot}}}}. \quad (7.6)$$

An unexpected observation is that the measured total cross section for the $^{124}\text{Xe}+^{112}\text{Sn}$ system is slightly larger than for the $^{136}\text{Xe}+^{124}\text{Sn}$ system. Since the total charge of the two Xe systems is the same, and the interaction radius should scale as $A^{1/3}$, one would expect a larger cross section for the heavier system as was seen with the 2016 Sn + Sn measurements [11]. A possible explanation of this effect could be that this is an artifact of the reduced acceptance for the $^{136}\text{Xe}+^{124}\text{Sn}$ system that resulted from moving the target ladder further back from the TPC by about 1 cm. The target ladder was moved back in order to reduce a sparking issue with the field cage which was not fully understood during the campaign. Because the total cross section is calculated using the minimum bias runs which have adjusted trigger settings to record more peripheral collisions, slower moving target-like particles are less likely to be successfully detected in comparison to the faster moving beam-like particles produced in more central collisions. Further investigations about the effects of this change in experimental geometry on the total experimental efficiencies and the determination of the reaction cross section may require future simulations.

Table 7.5 Summary of the determined reaction cross sections and b_{max} .

Beam	Target	σ (barn)	$\delta\sigma$ (barn)	b_{max} (fm)	δb_{max} (fm)
^{124}Xe	^{112}Sn	1.904	0.011	7.784	0.023
^{136}Xe	^{124}Sn	1.776	0.023	7.519	0.048

Another source of uncertainty beyond the systematics arises from the fact that a collision with a given value of b has an expected but not deterministic measured result for N_C . Following Poisson statistics then, statistical uncertainty bands may be assigned via

$$\delta N_C = \sqrt{N_C}, \quad b(N_C) \in [b(N_C - \delta N_C), b(N_C + \delta N_C)]. \quad (7.7)$$

The final impact parameter relationship obtained via scaling $\hat{b}$ by b_{max} , including error bands, is shown in Figure 7.18. To make an example centrality selection such as $b \leq 3$ fm, the cut to apply on the data would be to select events with $N_C \geq 53$. This corresponds to an event weighted average of $\langle b \rangle = 2.2$ fm. It is also common to make selections on $\hat{b}$ instead, where $N_C \geq 53$ would correspond to $\hat{b} \geq 0.4$ with a weighted average of $\langle \hat{b} \rangle = 0.28$.

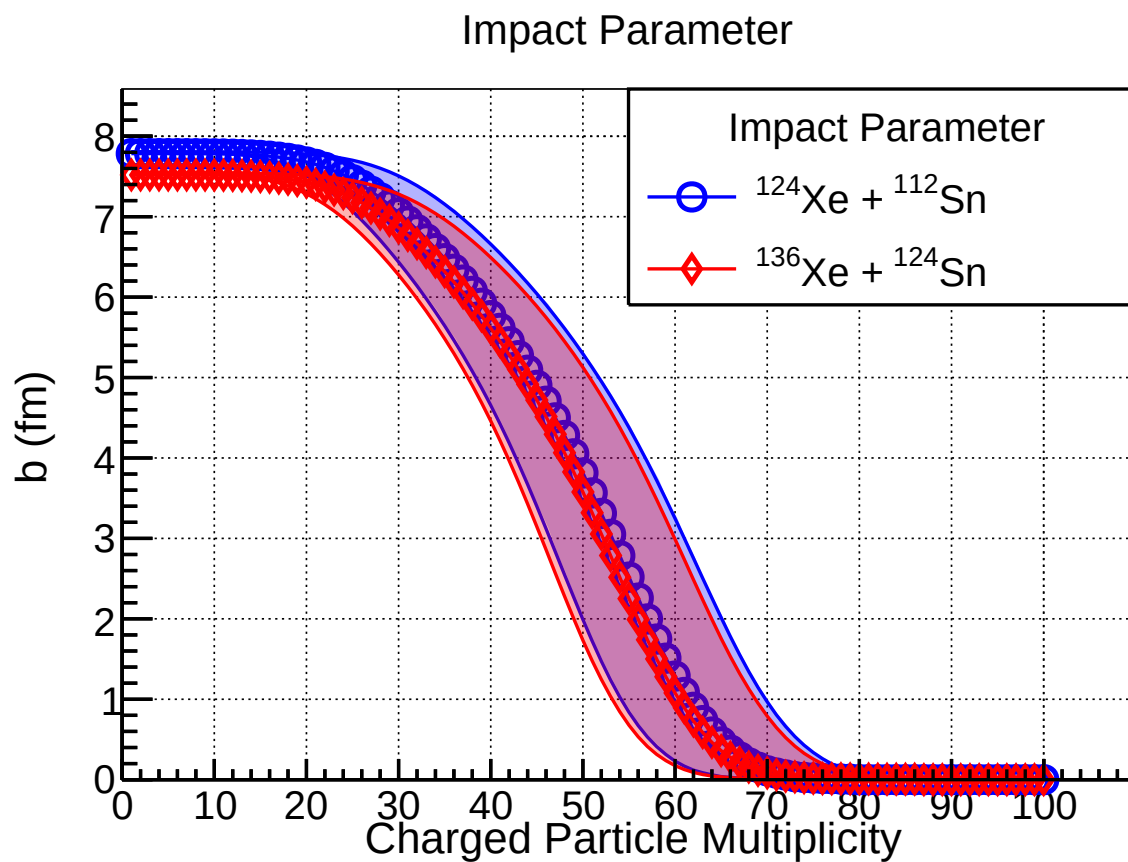


Figure 7.18 Inferred impact parameter, or the relationship between b and N_C , for both systems. The maximal b for each system is based on the maximal cross section calculated from data.

CHAPTER 8

PION SPECTRA RESULTS

This Chapter covers the final construction of the pion spectra including all counting corrections, and will provide preliminary comparisons to predictions from transport model calculations.

8.1 Pion Selection

This section details how pion tracks are selected and weighted in order to contribute towards making spectra.

8.1.1 PID Flattening

The particle lines in the PID can be represented using the Bethe-Bloch equation [33]

$$\frac{dE}{dx} = \frac{4\pi N e^4}{m_e c^2 \beta^2} z^2 \left(\ln \frac{2m_e c^2 \beta^2 \gamma^2}{I} - \beta^2 \right), \quad \beta \equiv \frac{p}{\sqrt{p^2 + m^2 c^2}}, \quad \gamma \equiv \frac{1}{1 - \beta^2} \quad (8.1)$$

where N is the number density of electrons in the TPC detection gas (P10), e is the elementary charge, z is the charge of the particle, m_e and m are the electron mass and particle mass, respectively, and I is the mean excitation energy of the detection gas. In practice, we use the following 5-parameter fit:

$$\frac{dE}{dx} = \frac{x_0}{\beta^{x_3}} [\ln(x_2 + \beta \gamma^{-x_4}) - \beta^{x_3} + x_1]. \quad (8.2)$$

By fitting Equation 8.2 to the pion lines and dividing the PID by the result, the PID is “flattened” so that π^+ and π^- are now straight lines along $y = 1$. By then taking the logarithm of this flattened PID, the spectra shown in Figure 8.1 is obtained, where the pion lines are now straight and lie along $y = 0$. In practice, the gates imposed to select the pions from this flattened PID are $-0.20 \leq y \leq 0.20$, where y is the flattened dE/dx axis.

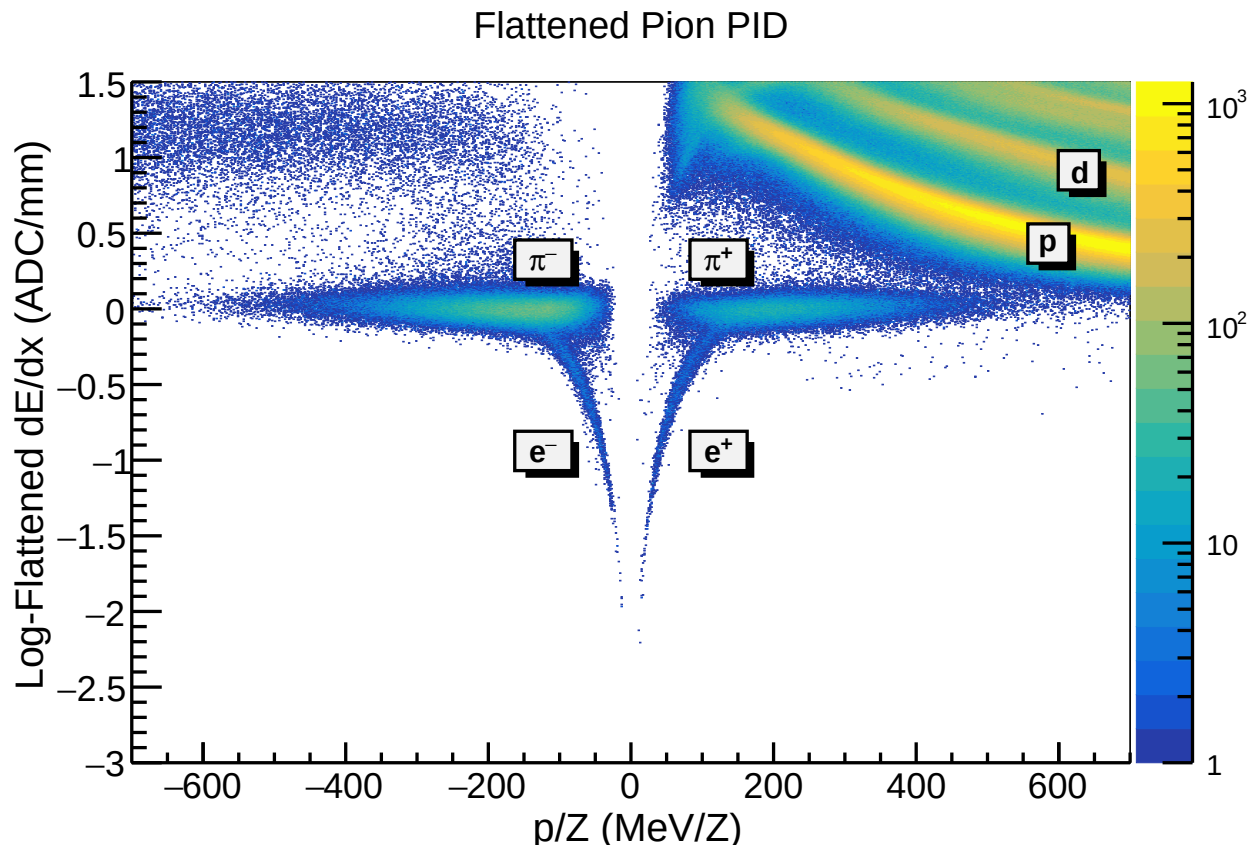


Figure 8.1 PID where the y-axis is the logarithm of “the standard dE/dx divided by a Bethe-Bloch fit of the pion line”. The result is that the pions lay along a straight line at $y = 0$, which simplifies the selection of pions.

8.1.2 Background Contamination Subtraction

The procedure for selecting pions discussed in Section 8.1.1 will inherently include contamination from various other particle lines. Specifically, the selection of π^- will include some e^- that deviate from the centroid of the line, and the selection of π^+ will include some of the p and e^+ lines that deviate from the p and e^+ lines as well. To quantify and correct for this contamination, slices of p/Z in the flattened PID are projected so that these particle lines can be fit with a convoluted Gaussian distribution that models the particle distributions reasonably well. An example of these projections for π^+ are shown for two lower momentum slices to illustrate the e^+ contamination from π^0 decay, and for two higher momentum slices to illustrate the p contamination, in Figure 8.2.

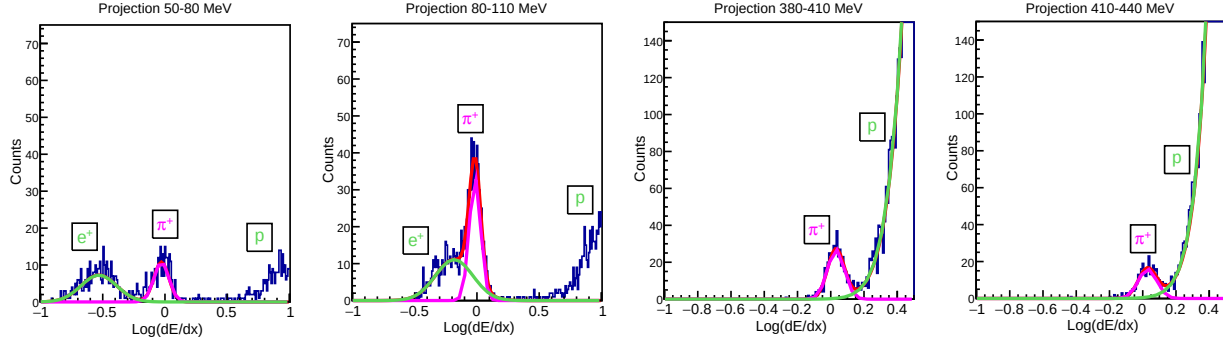


Figure 8.2 Example fits of the proton and positron contamination in the π^+ spectra. In the flattened PID, slices in p/Z are projected to the dE/dx axis and a simultaneous gaussian fit is performed.

By comparing the integral yields for the projected fits within the pion selection gate, a fraction of contamination is obtained as a function of p/Z , shown in Figure 8.3. The y-axis of this plot corresponds to the likelihood that chosen tracks are not pions, and so the weight of each track when filling the final spectra histograms can be accordingly adjusted. The peak in background contamination at low momentum comes from the e^+ line merging with the π^+ line, and then at higher momentum the fraction of p contamination slowly rises until eventually it will dominate the selection. An analogous procedure is performed for π^- , with the difference being that the only source of contamination in that case comes from the merging of the e^- line at low momentum. Overall the contamination fraction is smaller for π^- due to the increased overall yield compared to π^+ , and lack of contamination at higher momentum.

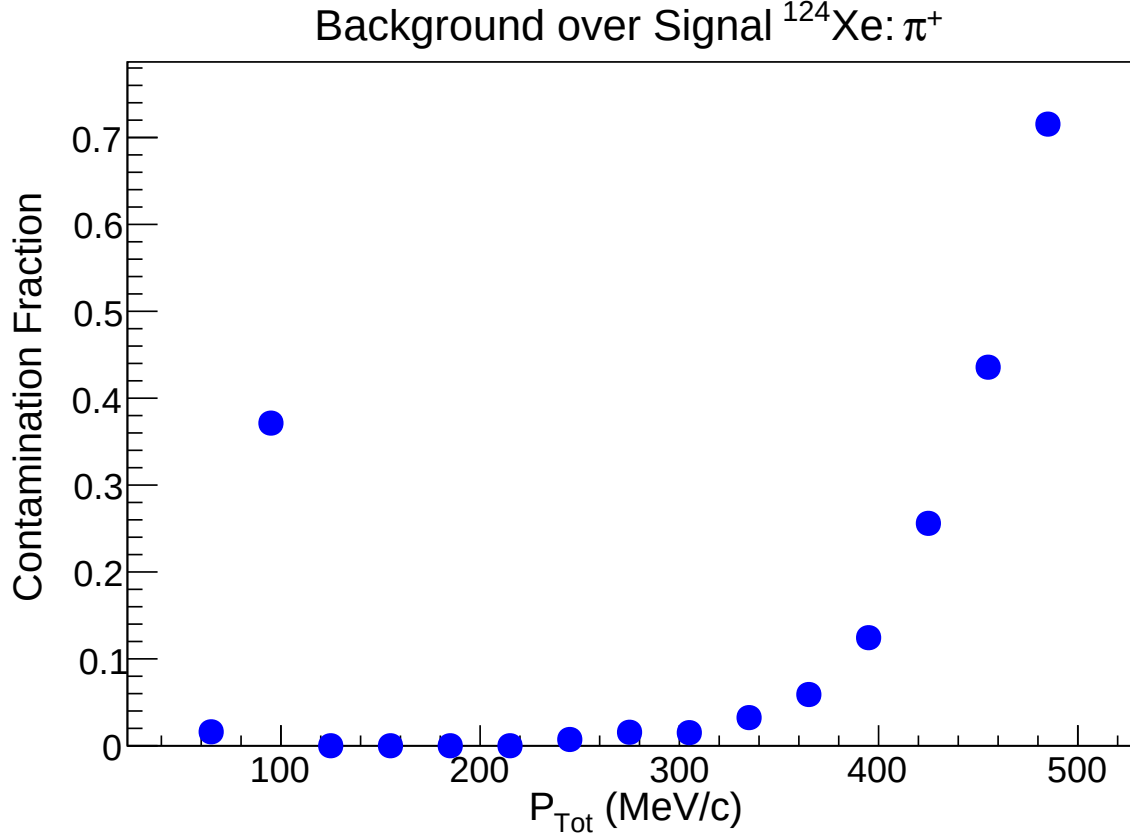


Figure 8.3 Resulting fraction of background contamination for the π^+ spectra as a function of p/Z . The peak near 100 MeV is from the positrons, and the proton contamination slowly rises after 300 MeV until it dominates the cut.

8.1.3 Embedding and Detection Efficiency

Every charged particle in the TPC is identified by finding its track and determining the track momentum. In order to correct the total pion yields for reconstruction efficiency, i.e. the fact that we cannot reconstruct 100 % of pions emitted in the collisions, and that some of the pion tracks are not cleanly resolved, we must perform simulations to estimate the pion detection efficiency. This is done via a Monte Carlo simulation process called *embedding* that takes into account the track cuts listed in Table 7.4.

In this simulation, single pions are generated at the target location with a uniform range of momentum between 0-0.7 GeV. These pions are transported through the S π RIT TPC geometry and fields using Geant4 physics libraries, and then digitized into a format that can be passed through the experimental reconstruction routine to create Monte Carlo (MC) tracks with appropriate ionization

densities and uncertainties. During the digitization process, simulation data from particles other than pions that may have been created during scattering or knockout reactions are discarded, as we are explicitly interested to study the likelihood of successful pion reconstruction. These MC pion tracks are then embedded onto real experimental data, and subjected to the same reconstruction routines that are used for the experimental data. For the embedded track to be analyzed and labeled a successfully embedded pion, it must then meet the same requirements as real data listed in Table 7.4. We can simulate 10^6 pions to obtain high statistical accuracy. The resulting PID for successfully reconstructed (embedded) MC pion tracks is shown in Figure 8.4. There are faint PID lines where the protons and deuterons would lie; these come from confusion of the simulated pion tracks with experimental data tracks during embedding, since as was mentioned above only information from simulated pions is included. Significant overlap of simulated pion tracks with other tracks in the data result in higher assigned values of dE/dx .

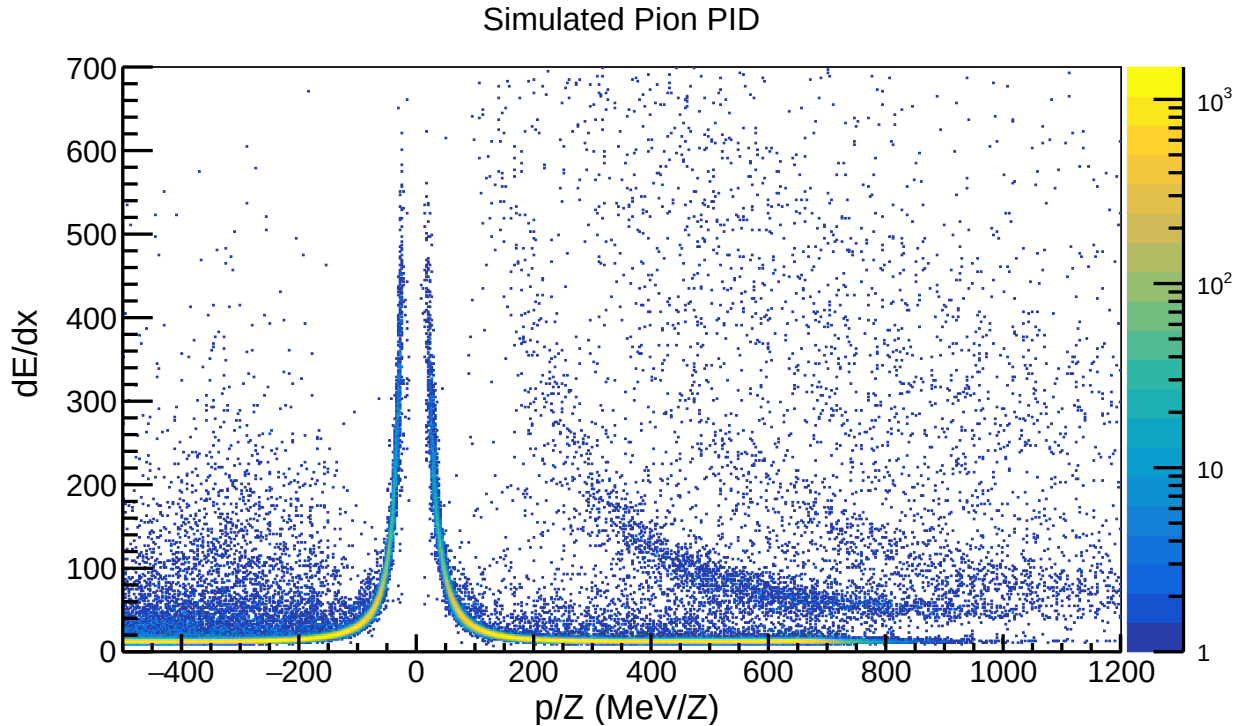


Figure 8.4 PID for the Monte Carlo simulated pions. The faint proton and deuteron lines are from simulated pions that have been confused with other particle tracks during the embedding process, which is folded into the detection efficiency.

The MC pions are then subjected to the same analysis procedures including the PID flattening

procedure as described in Section 8.1.1. This way, effects such as track overlap are taken into account in the determination of detection efficiency. A pion detection efficiency is obtained from ratio of MC pions that survive both reconstruction and selection via the flattening procedure compared to the total number simulated, shown as a function of θ and total momentum for the $^{124}\text{Xe}+^{112}\text{Sn}$ system in Figure 8.5, and for the $^{136}\text{Xe}+^{124}\text{Sn}$ system in Figure 8.6. Beam left and right pions are simulated separately when correcting for detection efficiency, since there are reconstruction differences for left and right particles as was discussed in Section 7.3.2. Because π^- are negatively charged and bend towards beam left as opposed to beam right for all the positively charged particles, their detection efficiency mirrors that of the π^+ .

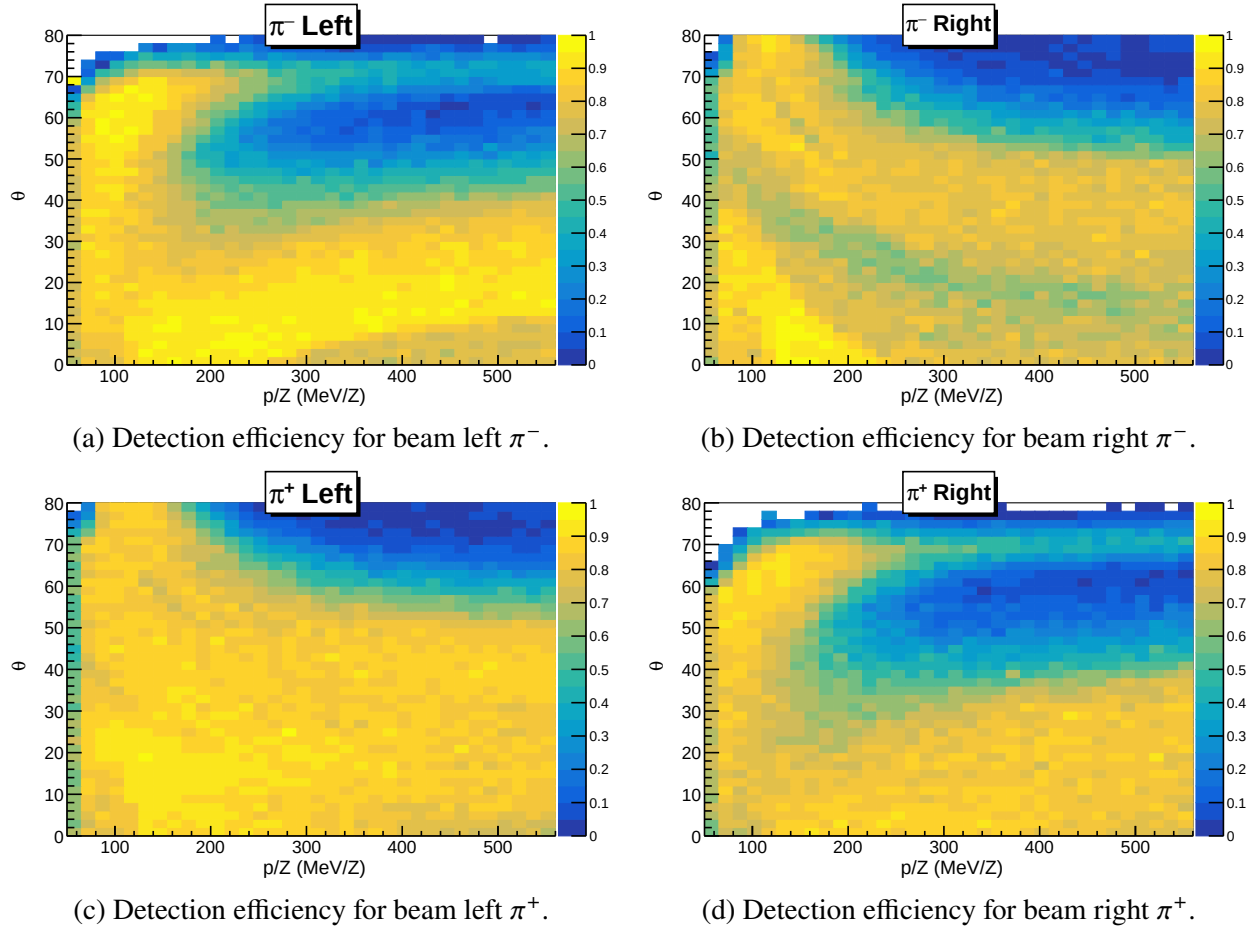


Figure 8.5 The results of simulating charged pion detection efficiencies for the $^{124}\text{Xe}+^{112}\text{Sn}$ system using the embedding method. Values near 1 on the color scale represent near 100 % detection efficiency, and vice-versa. “Left” and “Right” refer to the azimuthal cuts listed in Table 7.4.

In regards to the TPC reconstruction efficiency, the main difference between the two systems is the placement of the target ladder further back by approximately 1 cm for the $^{136}\text{Xe}+^{124}\text{Sn}$ system, as was mentioned in Section 7.3.1. This effect manifests in the embedding results as a reduction in the maximal polar angle θ for which pion tracks can be successfully reconstructed. For example, by comparing the Figure panels 8.5d and 8.6d, the cutoff for zero efficiency (white space) on the vertical axis is a few degrees less. This is also confirmation that the TPC geometry is accurately being accounted for in the simulations.

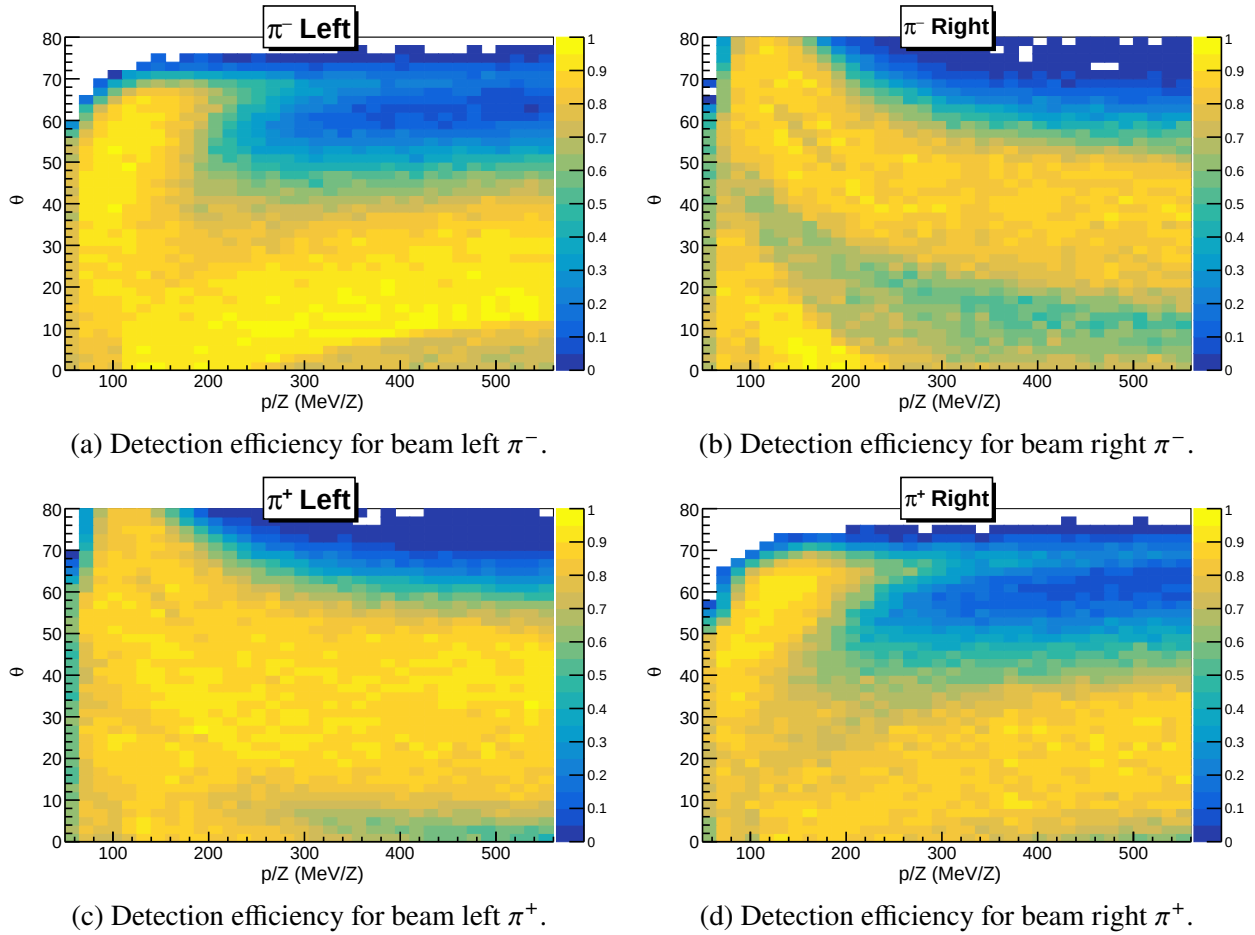


Figure 8.6 The results of simulating charged pion detection efficiencies for the $^{136}\text{Xe}+^{124}\text{Sn}$ system using the embedding method. The main difference between the two systems in regards to reconstruction efficiency is the placement of the target ladder further upstream by approximately 1 cm.

8.2 Construction of Charged Particle Spectra

This section details the procedures to construct spectra once the pion tracks have been selected and weighted as described above.

8.2.1 Transformation to CoM Frame

The beam axis is not exactly aligned with the z-direction at the target plane due to fringe effects from the SAMURAI magnetic field that start bending the beam before it enters the TPC. Using information from the BDC projection of the beam, the beam angles on target are calculated as

$$\theta_{a,\text{proj}} = \tan^{-1} \frac{p_x}{p_z}, \quad \theta_{b,\text{proj}} = \tan^{-1} \frac{p_y}{p_z}. \quad (8.3)$$

where here $p_{x,y,z}$ are the momentum components of the beam calculated at the target plane. A distribution of these beam angles for both systems is shown in Figure 8.7; they are measured in thousandths of a radian. Evidence of the tight beam tuning is reflected in the concentration of these angles. As one might expect from the vertical orientation of the SAMURAI field, there is nearly no deflection in the vertical direction, and a small deflection in the horizontal direction. The horizontal deflection is slightly smaller for the $^{136}\text{Xe} + ^{124}\text{Sn}$ system due to the larger mass of the beam.

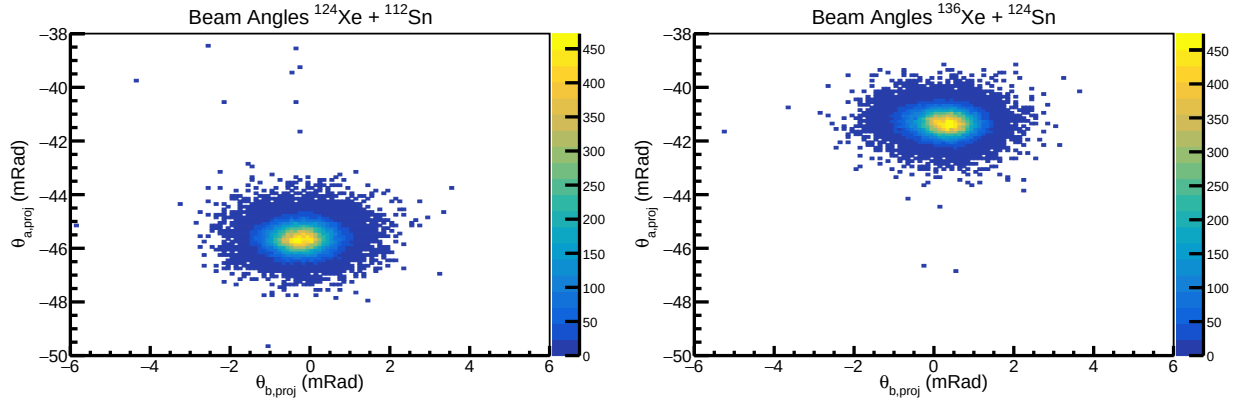


Figure 8.7 Beam angles on target as determined using the BDC information. Before boosting events into the CoM frame, a rotation using these angles is performed to properly align with the z-axis.

To facilitate an easier conversion to the center of mass (CoM) frame, these events are first rotated to align the beam momentum exactly with the z-axis by using these beam angles. Then,

they are boosted to the CoM frame using a standard Lorentz transformation [74] based on the incident energy of the beam. Once transformed to the CoM frame, tracks are subjected to an additional forward angle selection of $\theta_{CoM} \leq 90^\circ$ for our subsequent analyses of the pion events. This selection is also made in the transport model simulation predictions that will be compared to the data in Sections 8.3.2 and 8.3.3.

8.2.2 Pion Rapidities

A useful observable to consider in the context of our HICs is a measure of the longitudinal motion of an emitted particle called *rapidity*:

$$y = \frac{1}{2} \ln \left(\frac{E + p_z}{E - p_z} \right) = \tanh^{-1} \left(\frac{p_z}{E} \right) \quad (8.4)$$

where E and p_z are the particle's total energy and component of momentum along the beam direction, respectively. For pions, we can define a normalized rapidity y_0 as [75].

$$y_0 = \frac{y_{\pi CoM}}{y_{beam} - y_{CoM}} \quad (8.5)$$

where $y_{\pi CoM}$ is the pion rapidity in the CoM frame, y_{CoM} is the beam-target system rapidity in the CoM frame, and y_{beam} is the rapidity of the beam in the laboratory frame. Values of $y_0 > 0$ correspond to pions with $\theta_{CoM} < 90^\circ$, and $y_0 = 1$ corresponds to pions at beam rapidity. Figure 8.8 shows these distributions for the charged pions of both systems. Figure 8.9 shows the $p_T - y_0$ spectra for both systems with the distributions reflected about $y_0 = 0$, assuming symmetry for the collisions. Because the beam-target systems are not themselves symmetric, this is only an approximation.

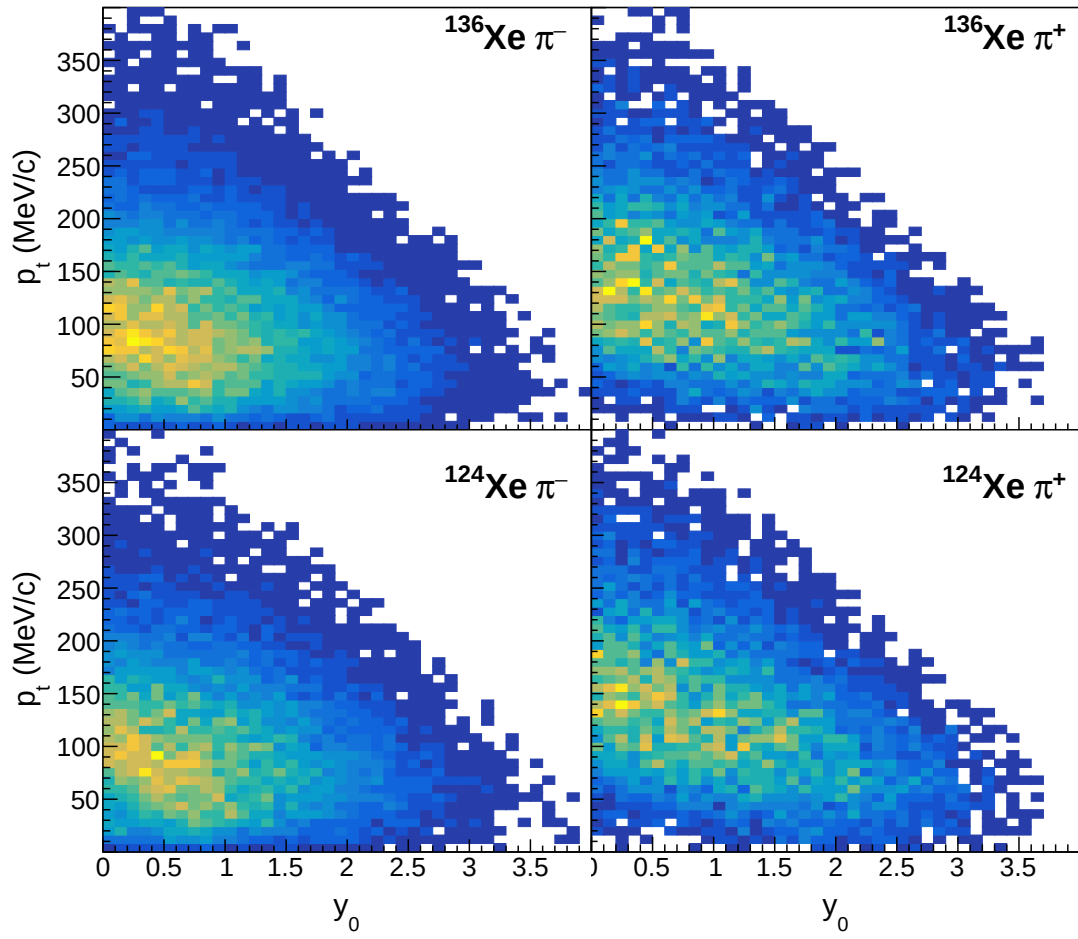


Figure 8.8 Transverse momentum rapidity spectra of π^- and π^+ .

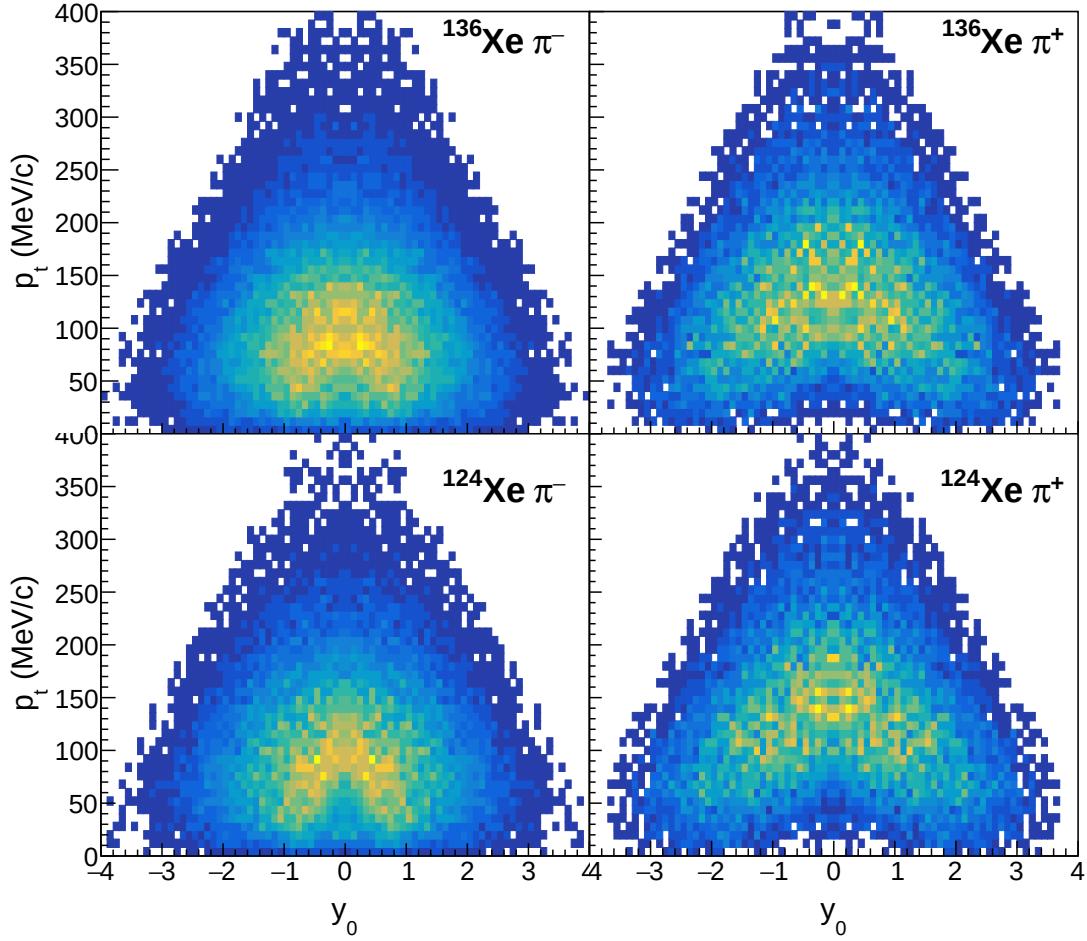


Figure 8.9 Transverse momentum rapidity spectra of π^- and π^+ . The histograms are reflected about $y_0 = 0$ assuming symmetry, since there is minimal acceptance for $y_0 < 0$ and negative values are cut out by the forward angle selection.

8.2.3 Pion Yields and Ratios

Once transformed into the CoM frame, each pion track that meets the criteria listed in Table 7.4 is filled into a histogram with a weight determined by the background fraction and detection efficiency at the reconstructed values of total momentum and polar angle. Once the spectra have been obtained in this way, a scaling correction is made for the angular selections dividing by the fraction of total angular coverage. The left and right azimuthal angle cuts cover 110° , and integration over the forward polar angles yields a factor of 2π :

$$\frac{\Omega_{\text{Selected}}}{\Omega_{\text{Total}}} = \frac{1}{4\pi} \left(\frac{110}{360} \cdot 2\pi \right) = 0.153. \quad (8.6)$$

The spectra shown in Figures 8.10 and 8.11 are the results of the completed correction process for each system, including background, efficiency, and angular coverage. The left panels show the differential π^+ and π^- yields as a function of transverse momentum, and the right panels show the Single Ratio (SR) of π^- yields to π^+ yields within a system. Constructing the SR is useful because uncertainties common to both systems will cancel out, providing a more accurate measurement. Even further, the Double Ratio (DR) can be constructed, which is the SR of one system divided by the other. A typical convention is to use the more asymmetric system as the numerator. This DR is shown in Figure 8.12.

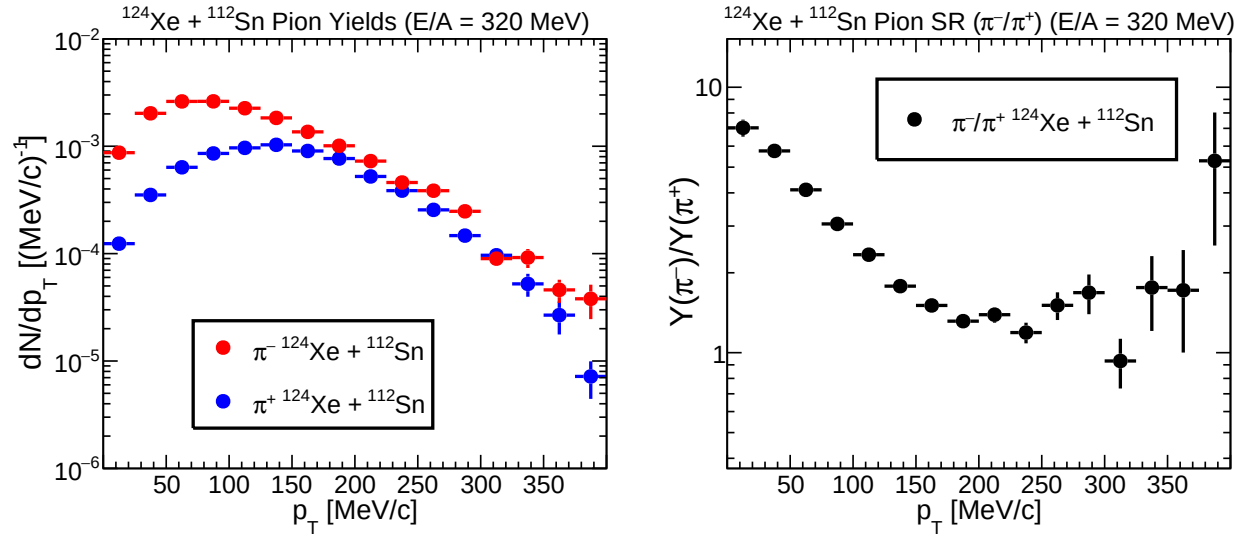


Figure 8.10 (Left) Corrected single particle pion transverse momentum spectra for the $^{124}\text{Xe} + ^{112}\text{Sn}$ system. (Right) Spectral single ratio of π^-/π^+ . These data have been corrected for detection efficiency, background subtraction, and angular coverages.

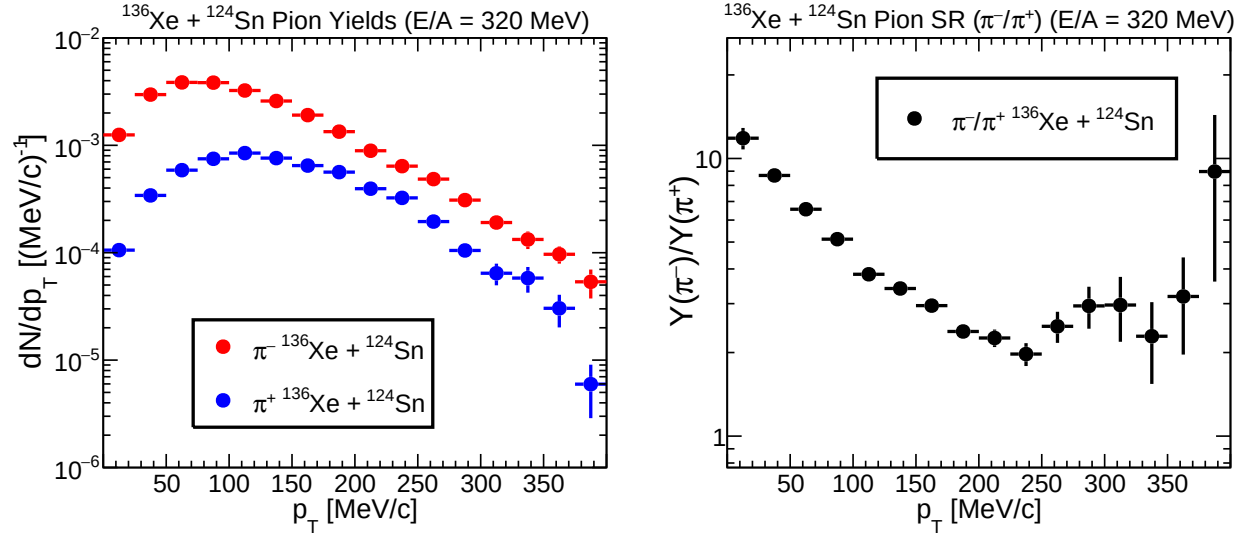


Figure 8.11 (Left) Corrected single particle pion transverse momentum spectra for the $^{136}\text{Xe} + ^{124}\text{Sn}$ system. (Right) Spectral single ratio of π^-/π^+ . These data have been corrected for detection efficiency, background subtraction, and angular coverages.

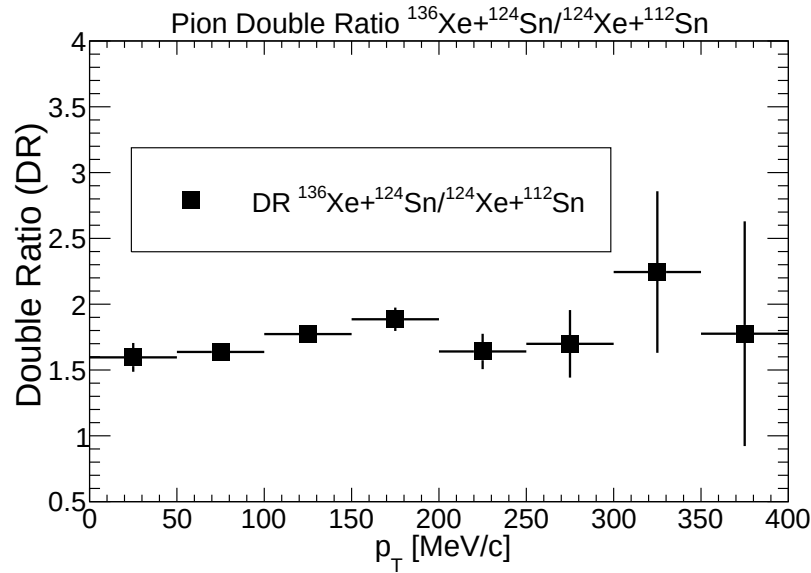


Figure 8.12 Double ratio of the $^{136}\text{Xe} + ^{124}\text{Sn}$ system SR over the $^{124}\text{Xe} + ^{112}\text{Sn}$ system SR.

An interesting systematic study is the trend of the pion yields as a function of total charged particle multiplicity. Figures 8.13 and 8.14 show this trend for the individual pions yields in the left panel, and SR in the right. The individual yields can be seen to grow non-linearly with N_C , and the ratios show an increase for the largest values of N_C . To check if this effect is physical or systematic, the cross system pion yield ratios were also studied. These are the yield ratios of just π^- from one

system compared to the other, and the same for π^+ , which are seen to be stable as a function of N_C as shown in Figure 8.15. The DR as a function of N_C is also seen to be stable, shown in Figure 8.16. This suggests that the individual yield and SR increases result from a physical effect.

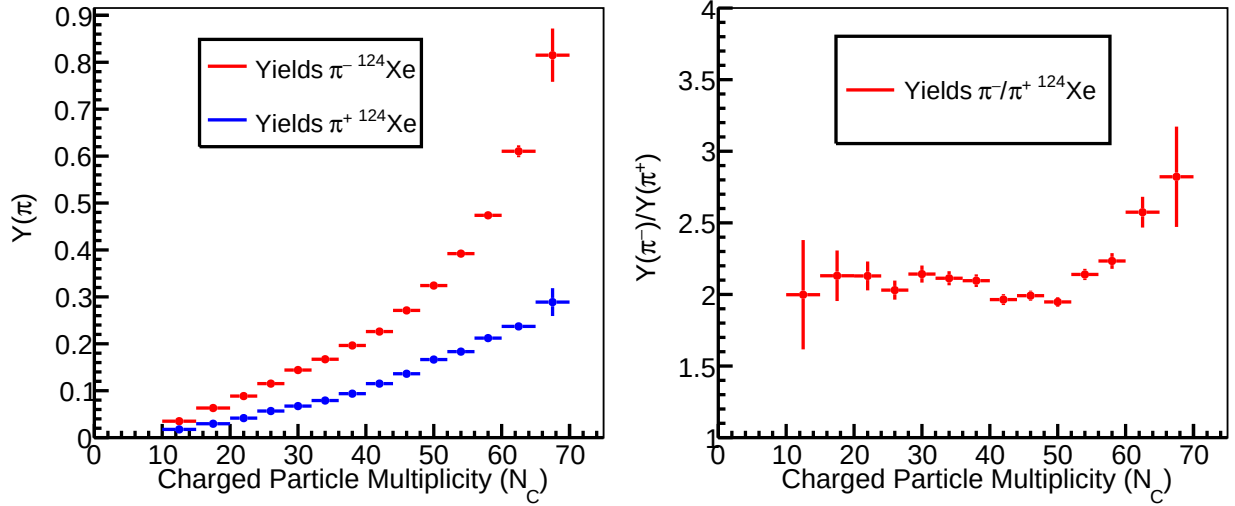


Figure 8.13 (Left) Corrected pion total yields for the $^{124}\text{Xe} + ^{112}\text{Sn}$ system as a function of N_C . (Right) Corrected pion total ratio as a function of N_C .

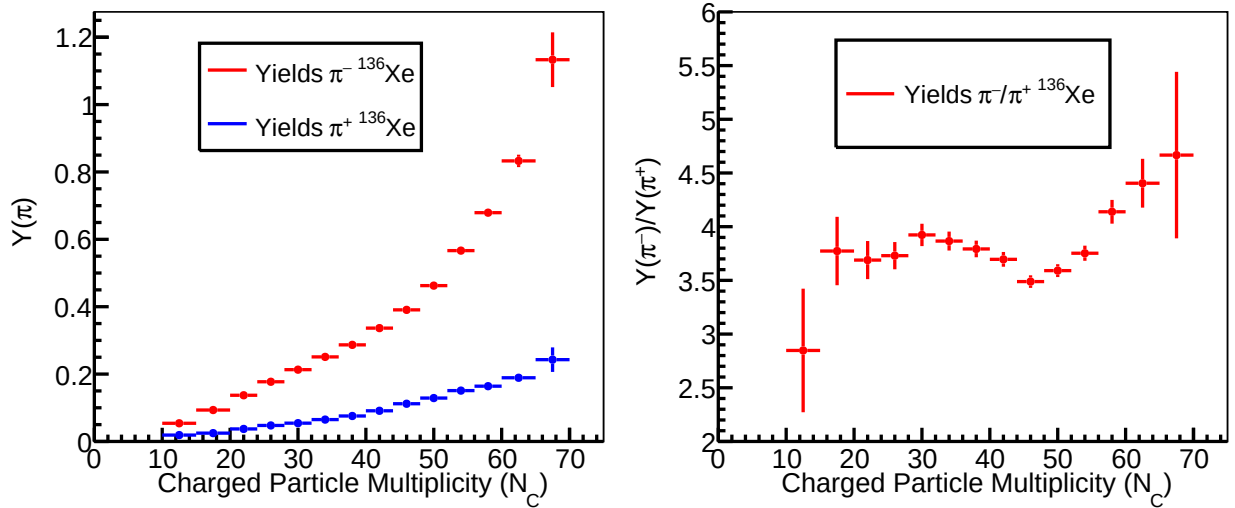


Figure 8.14 (Left) Corrected pion total yields for the $^{136}\text{Xe} + ^{124}\text{Sn}$ system as a function of N_C . (Right) Corrected pion total ratio as a function of N_C .

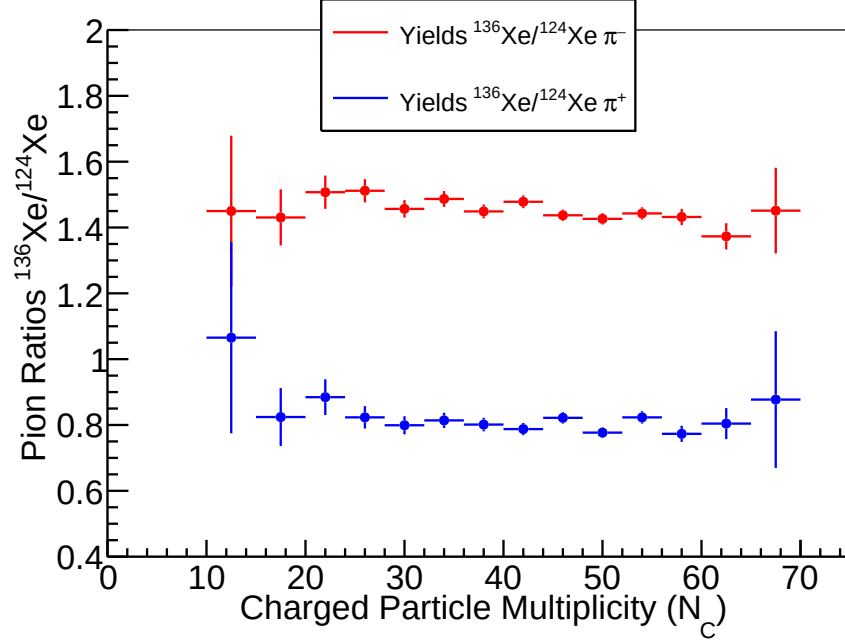


Figure 8.15 Cross system pion ratios as a function of N_C .

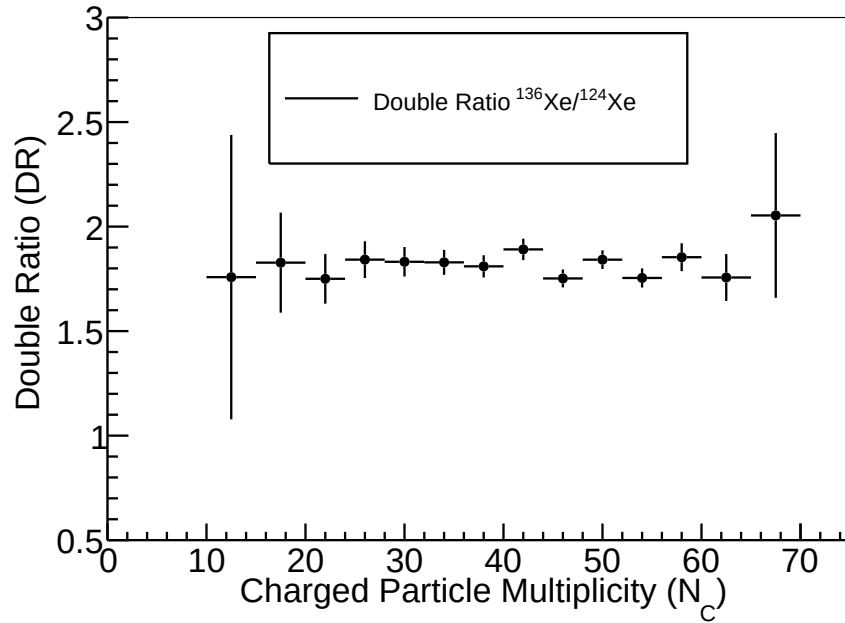


Figure 8.16 The pion double ratio as a function of N_C .

8.3 Comparison to Theory

In this section, I compare the particle spectra, single ratios, and double ratios to the provided transport model predictions from AMD + sJAM [2] and dcQMD [3]. Details about the transport

models and some of their relevant parameter settings are provided.

8.3.1 EoS Parameters

Listed in Tables 8.2 and 8.1 are some of the parameter setting for the transport models that are used to compare with the preliminary pion results of the 2024 Xe campaigns. Here, ρ_0 is the nuclear saturation density, K is the Symmetry Energy curvature, m^*/m_N is the isoscalar effective nucleon mass, S_0 is the Symmetry Energy at ρ_0 , L is the slope of the density dependence of the Symmetry Energy at ρ_0 , and $\Delta m_{np}^*/(m_N\delta)$ is the neutron-proton effective mass splitting. More details about the transport models will be given in the following Sections 8.3.2 and 8.3.3.

Table 8.1 Some nuclear matter properties for the SLy4 interaction of the AMD + sJAM transport model. Modified from [2].

Parameters	AMD(SLy4) + sJAM
ρ_0 [fm ⁻³]	0.160
K [MeV]	230
m^*/m_N	0.70
S_0 [MeV]	32.0
L [MeV]	46
$\Delta m_{np}^*/(m_N\delta)$	-0.18
in n-rich matter	$m_n^* < m_p^*$

Table 8.2 Some nuclear matter properties for the dcQMD transport model. Modified from [3].

Parameters	dcQMD
ρ_0 [fm ⁻³]	0.160
K_0 [MeV]	230
m^*/m_N	0.67
$S(0.625\rho_0)$ [MeV]	25.5
L [MeV]	58
$\Delta m_{np}^*/(m_N\delta)$	-0.07
in n-rich matter	$m_n^* < m_p^*$

The predictions provided to compare with our experimental results were calculated by members of the Transport Model Evaluation Project (TMEP) [76, 77, 78, 79, 80], who have also provided extensive theoretical calculations for the results of the 2016 measurements. TMEP was originally created in order to compare a collection of commonly used transport codes using the same inputs

and initial conditions, to observe the different evolutions of various observables. Some preliminary simulations have been performed by two particular TMEP transport models for fixed EoS parameter settings *without* having seen our results first. That is, these results are predictions which have been tuned for parameter values based purely on other experimental data, and not from a fit to the current 2024 measurements. Constraints on the symmetry energy will be obtained by determining a range of parameter values which best reconstruct the measurements. These involved simulations have not yet been performed.

8.3.2 1D Spectra Comparison AMD(SLy4) + sJAM

The AMD + sJAM framework is an improvement on the AMD + JAM model [52, 81, 53], which underestimated the pion ratio. Initially, anti-symmetrized molecular dynamics (AMD) [82] was used to solve the nucleon dynamics of the collisions, and was combined with a hadronic cascade model (JAM) [83] to handle the reactions related to pions and Δ resonances. The extension from JAM to sJAM was made by incorporating momentum-dependent mean field potentials for the Δ resonance in the collision term of the transport model, which were not present before. Several interaction models are available for the framework; currently, the SLy4 framework has been employed. In general, the SLy4 interaction of the AMD + sJAM framework tends to slightly overestimate the pion ratios, which is consistent with observations made for the 2016 Sn + Sn data in [2]. AMD + sJAM is able to reproduce the data much better than AMD + JAM as a result of the momentum-dependent mean field potentials.

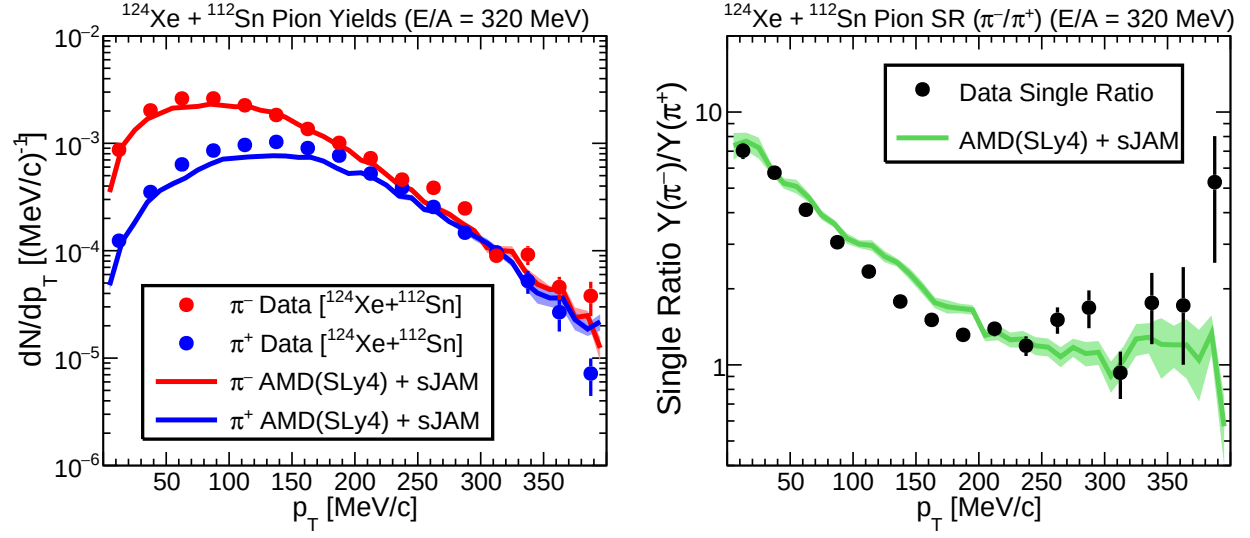


Figure 8.17 (Left) Comparison of single particle pion transverse momentum spectra for the $^{124}\text{Xe} + ^{112}\text{Sn}$ system to calculations from AMD. (Right) Spectral single ratio comparison of π^-/π^+ . Data and calculations both have an impact parameter selection of $b < 3 fm$.

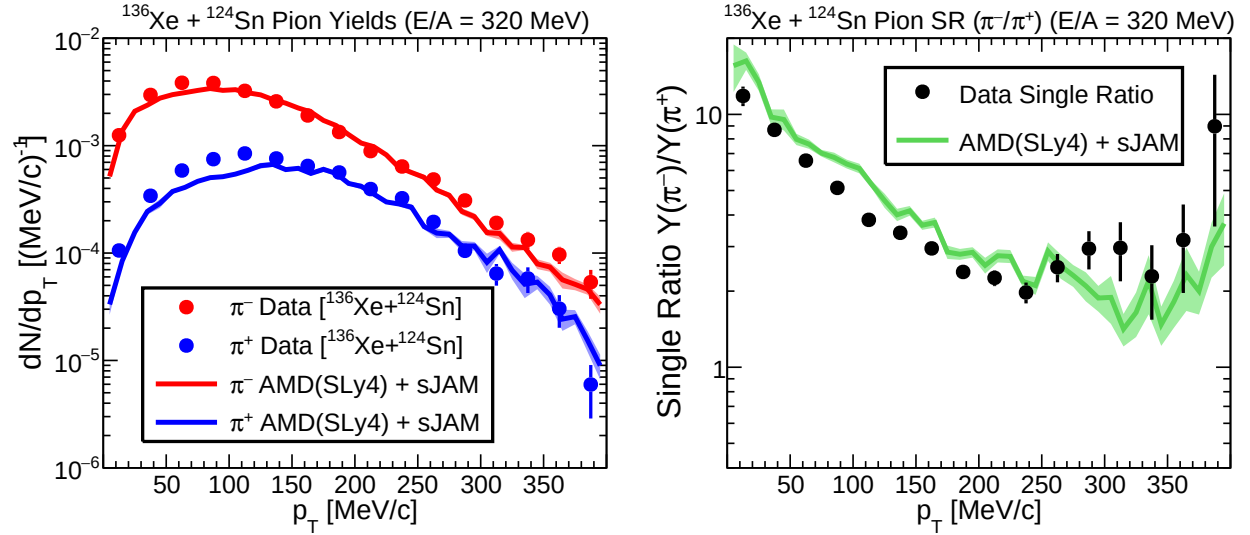


Figure 8.18 (Left) Comparison of single particle pion transverse momentum spectra for the $^{136}\text{Xe} + ^{124}\text{Sn}$ system to calculations from AMD. (Right) Spectral single ratio comparison of π^-/π^+ . Data and calculations both have an impact parameter selection of $b < 3 fm$.

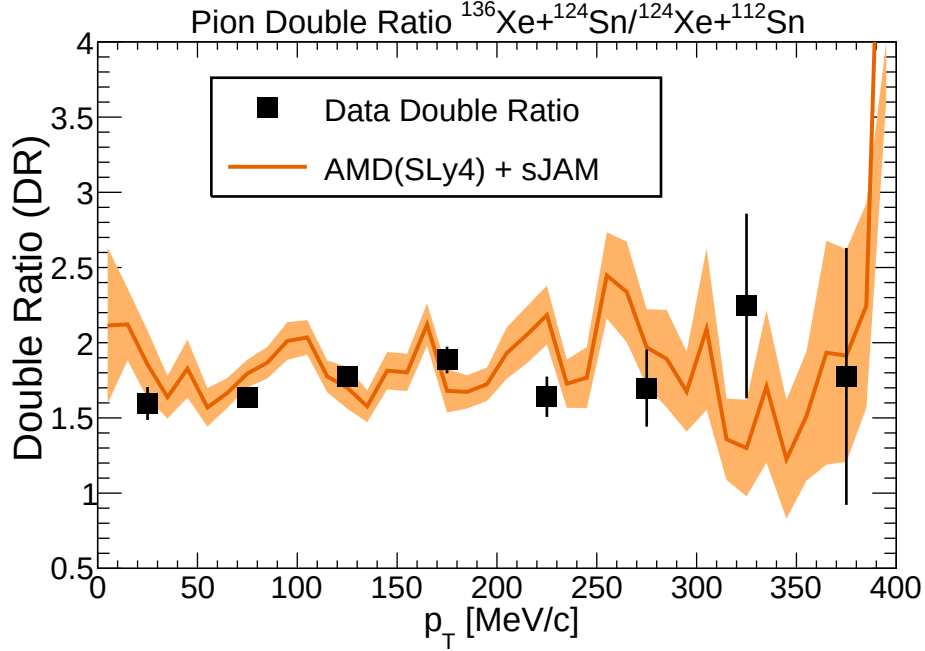


Figure 8.19 Double ratio comparison to calculations from AMD.

8.3.3 1D Spectra Comparison dcQMD

The Dan Cozma quantum molecular dynamics (dcQMD) model was developed based on the Tübingen QMD (tuQMD) model [84]. While several adjustments have been made, the main improvement of dcQMD over tuQMD relevant to our experimental campaign is the inclusion of isospin-dependent potentials, notably for the $\Delta(1232)$ resonance. Details about the most recent iteration of the dcQMD transport model can be found in [3]. An interesting note is that despite the higher incident Xe beam energy of 320 AMeV, there is still an under prediction of the SR at low p_T values from the dcQMD model. It can be seen most notably for the less asymmetric system in Figure 8.20, which mirrors the observations seen in comparisons with the 2016 measurements in [12]. This would suggest that either 320 AMeV is not a high enough incident energy to significantly reduce pion production coming from non-resonance, or that the discrepancy with the dcQMD model is from other factors.

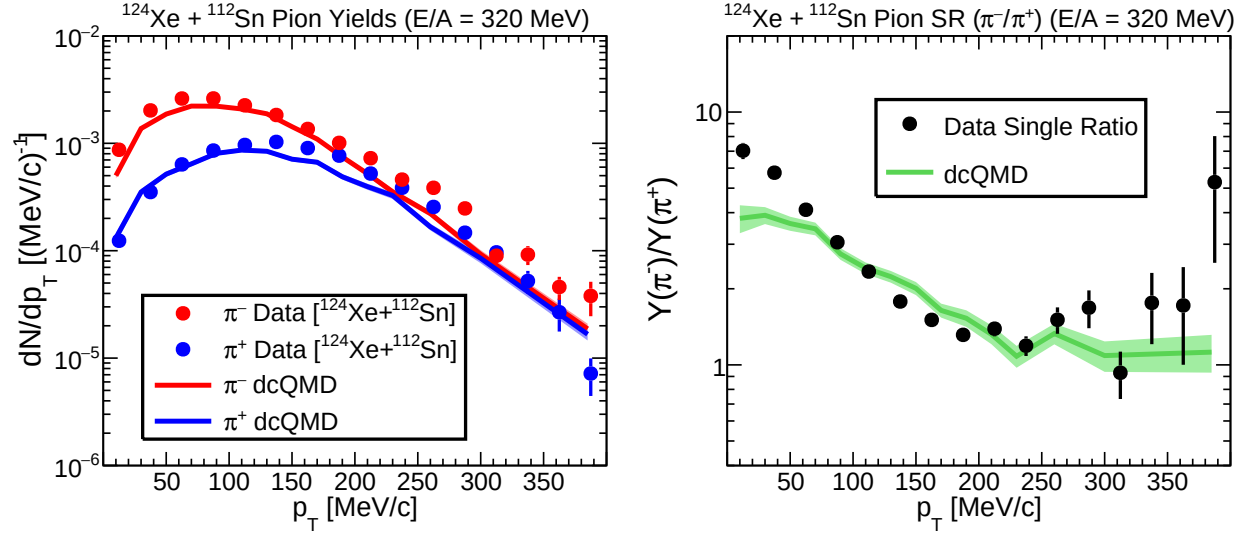


Figure 8.20 (Left) Comparison of single particle pion transverse momentum spectra for the $^{124}\text{Xe} + ^{112}\text{Sn}$ system to calculations from dcQMD. (Right) Spectral single ratio comparison of π^-/π^+ . Data and calculations both have an impact parameter selection of $b < 3 fm$.

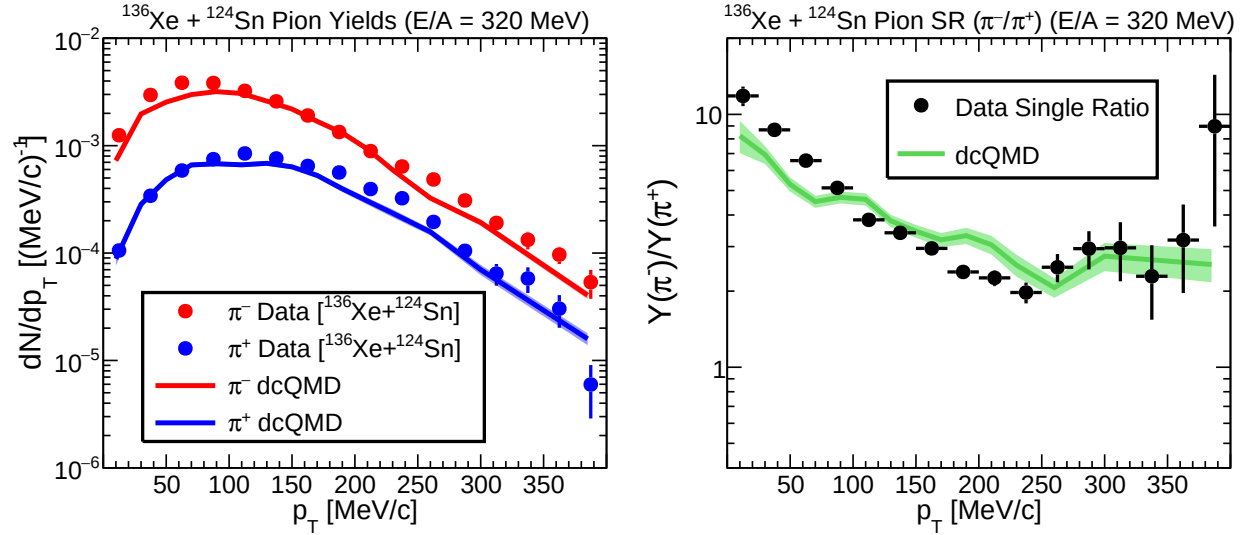


Figure 8.21 (Left) Comparison of single particle pion transverse momentum spectra for the $^{136}\text{Xe} + ^{124}\text{Sn}$ system to calculations from dcQMD. (Right) Spectral single ratio comparison of π^-/π^+ . Data and calculations both have an impact parameter selection of $b < 3 fm$.

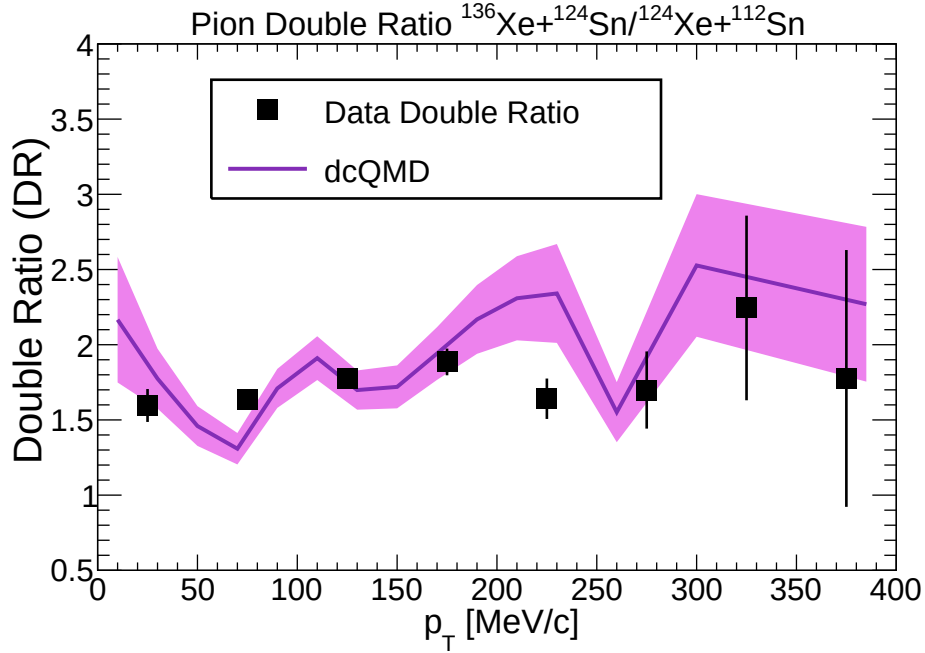


Figure 8.22 Double ratio comparison to calculations from dcQMD.

8.3.4 Total Yield Comparison

Figure 8.23 shows the total pion yields for the 2024 data with respect to the 2016 data, including the dcQMD and AMD(SLy4) + sJAM predictions. There is roughly a factor of 2 more pion production compared to 2016 (red x's), which suggests that the incident beam energy of 320 AMeV available with the Xe beams results in more energy available to occupy the $\Delta(1232)$ resonances. Despite the higher incident beam energy, the total yield ratios lie along the same $(N/Z)^{3.4}$ power law fit as the 2016 data, as shown in Figure 8.24. This suggests that overall, the densities reached in collisions during the 2024 experiment were not much greater.

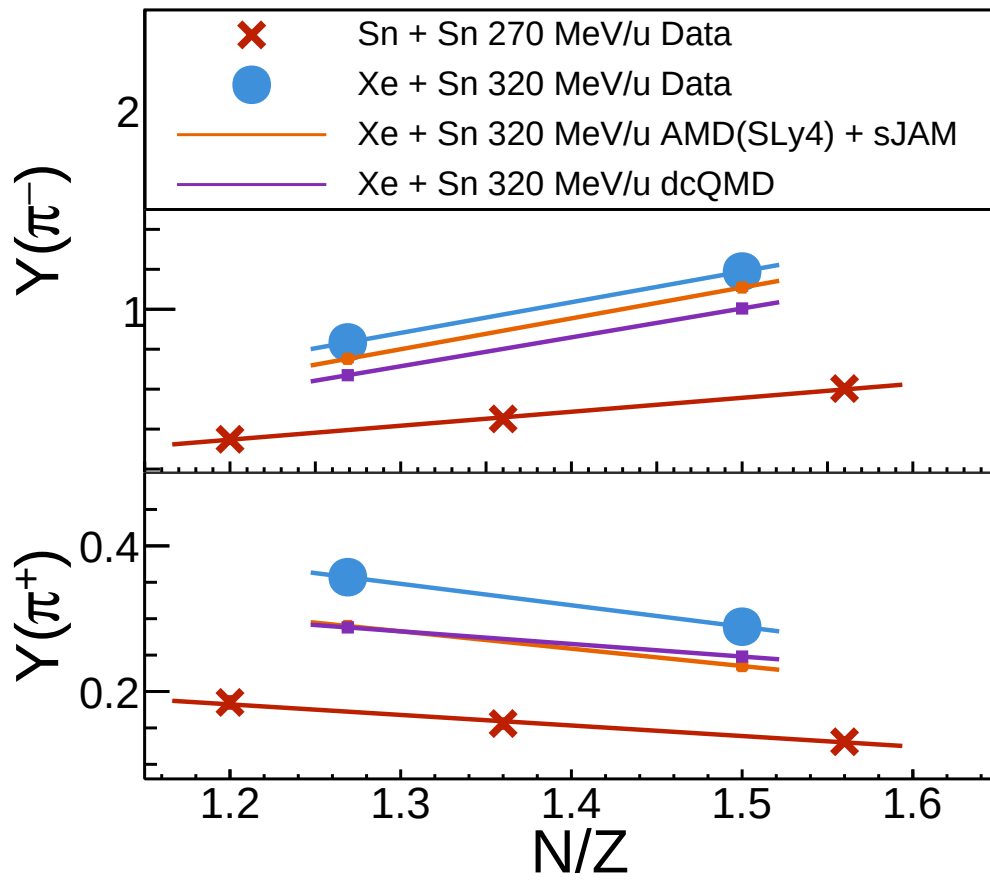


Figure 8.23 Total integrated yields of π^- and π^+ as a function of the beam-target system asymmetry values. Results of the previous Sn + Sn measurements are included with a red 'X'. The total yields for the Xe + Sn data are roughly a factor of 2 larger than the Sn + Sn data due to the increase in beam energy.

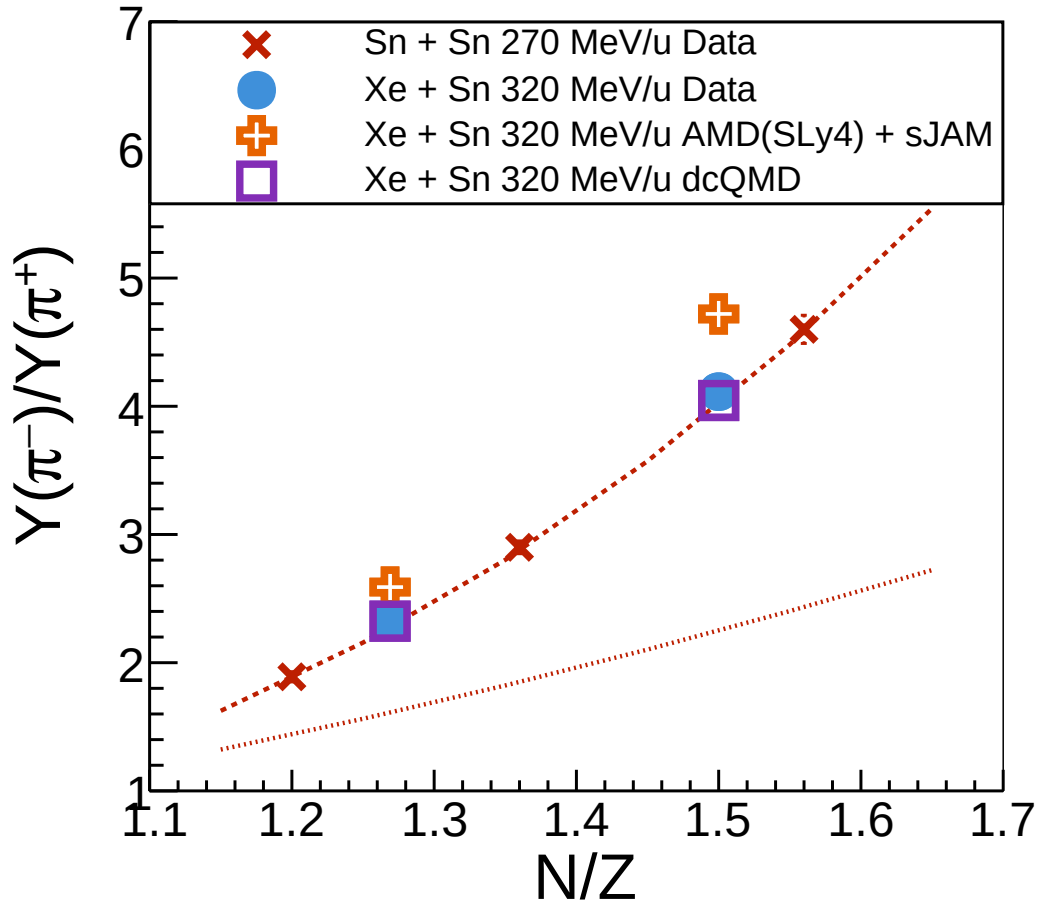


Figure 8.24 Total integrated yield ratios of π^- and π^+ as a function of the beam-target system asymmetry values. Results of the previous Sn + Sn measurements are included with a red ‘X’. Assuming a simple Δ -resonance model, the pion ratios would be expected to scale as $(N/Z)^2$, represented by the dotted red line. The data for both the Xe + Sn and Sn + Sn campaigns lie on the dashed red line of $(N/Z)^{3.6}$, indicating the role of more complicated factors.

Shown in Table 8.3, the two transport model DR predictions agree quite well with the data, to within one error bar of the experimental uncertainty. This shows the improvement in transport models since the initial comparisons of the 2016 data [9], where none of the provided calculations were able to reproduce the measured results. This highlights the improvements of TMEP models.

Table 8.3 Summary of the total DR values from data and the provided transport model predictions.

Source	Total DR	Error
S π RIT Data	1.76	0.06
dcQMD	1.74	0.03
AMD(SLy4)+sJAM	1.82	0.03

8.4 Summary and Outlook

Preliminary charged pion emission yields have been obtained from analysis of heavy ion collisions between $^{124}\text{Xe} + ^{112}\text{Sn}$ and $^{136}\text{Xe} + ^{124}\text{Sn}$ beam-target systems at $E/A = 320$ MeV from an experimental campaign which took place at RIKEN's RIBF in 2024. These measurements compliment a previous campaign from 2016 of various Sn + Sn systems at $E/A = 270$ MeV. The higher energies available using the Xe beams produce higher overall charged pion yields, while also addressing concerns about non-resonant pion production for which the threshold.

Corrections to data from the Xe + Sn experiments are still being finalized, and therefore the results are not yet ready for official publication. Namely, the space charge correction method requires more refinement to be satisfactory, and the embedding efficiencies need to be studied more carefully in order to be properly understood. The parameters and cuts for selecting the pions may also be improved in the future.

Preliminary comparisons to theory are promising, although these simulations also need to be adjusted. There is strong qualitative agreement for the single particle spectra, but experiment is quantitatively larger than both the AMD and QMD calculations. The agreement between the double ratios, which significantly cancel out systematics in the experimental data, show excellent agreement.

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APPENDIX A

CALCULATION OF PION PRODUCTION THRESHOLDS

From relativistic two-body kinematics for the collision of an accelerated nucleon in the beam with a stationary nucleon in the target, the minimum threshold for charged pion production is ≈ 290 MeV [85, 74], which we will now derive. Production of a pion from a nucleon-nucleon collision can be written as

$$\mathbb{N} + \mathbb{N} \rightarrow \mathbb{N} + \mathbb{N} + \pi \quad (\text{A.1})$$

where $\mathbb{N}$ and π represent a nucleon and pion, respectively. The invariant *Mandelstam* variable s [86] is related to the square of the four-momentum for the nucleon-nucleon system (p_{CoM}^2). For a collision between a beam nucleon (A) and fixed target nucleon (B), we have the following definition:

$$s \equiv E_{CoM}^2 \equiv -p_{CoM}^2 = (E_A + E_B)^2 - (\mathbf{p}_A + \mathbf{p}_B)^2 = m_A^2 + m_B^2 - 2\mathbf{p}_A \mathbf{p}_B + 2E_A E_B. \quad (\text{A.2})$$

With the fixed target nucleon at rest in the lab frame, we have $\mathbf{p}_B = 0$ and $E_B = m_B$ so that

$$s = m_A^2 + m_B^2 + 2E_A m_B = 2m_N^2 + 2E_{beam} m_N. \quad (\text{A.3})$$

The threshold condition is that the pion and two nucleons are produced at rest relative to each other in the center-of-mass (CoM) frame. In other words, the invariant CoM energy ($\sqrt{s}$) needs to be at least the sum of the rest masses of the two nucleons and the pion:

$$\sqrt{s} = 2m_N + m_\pi. \quad (\text{A.4})$$

Solving equations A.3 and A.4 for the beam kinetic energy (T) gives

$$T_{threshold} \equiv E_{beam} - m_N = 2m_\pi + \frac{m_\pi^2}{2m_N}. \quad (\text{A.5})$$

Substituting values of $m_N = 939$ MeV and $m_\pi = 140$ MeV for the nucleon and charged pion ($\pi^\pm$) mass gives $T_{threshold} \approx 290$ MeV. Using $m_\pi = 135$ MeV for the neutral pion (π^0) mass gives $T_{threshold} \approx 280$ MeV. However, in the particular case of a HIC (nucleus-nucleus collisions instead of nucleon-nucleon), there is a known lowering effect on this threshold. This is a result of the superposition of internal nucleon momentum with the overall nucleus-nucleus relative momentum [87].

APPENDIX B

RUNS ANALYZED FOR PION SPECTRA

Table B.1 A list of the run numbers used to create the pion spectra for each beam-target system.

System	Total Runs	Run numbers
$^{124}\text{Xe}+^{112}\text{Sn}$	69	1151, 1153, 1155, 1158, 1160, 1161, 1163, 1164, 1165, 1166, 1167, 1168, 1169, 1170, 1171, 1172, 1174, 1175, 1176, 1179, 1181, 1182, 1183, 1184, 1185, 1189, 1190, 1191, 1192, 1195, 1241, 1242, 1243, 1244, 1245, 1246, 1247, 1248, 1251, 1252, 1253, 1254, 1255, 1256, 1273, 1274, 1275, 1276, 1277, 1279, 1280, 1281, 1282, 1283, 1284, 1286, 1287, 1288, 1289, 1290, 1291, 1292, 1293, 1294, 1295, 1296, 1297, 1298, 1299
$^{136}\text{Xe}+^{124}\text{Sn}$	69	1634, 1635, 1636, 1637, 1638, 1639, 1640, 1641, 1642, 1643, 1644, 1645, 1646, 1647, 1648, 1649, 1650, 1651, 1652, 1653, 1654, 1655, 1656, 1657, 1659, 1660, 1661, 1662, 1663, 1664, 1665, 1668, 1669, 1697, 1706, 1707, 1708, 1709, 1731, 1732, 1733, 1734, 1741, 1742, 1743, 1744, 1745, 1746, 1775, 1776, 1777, 1778, 1779, 1794, 1800, 1801, 1802, 1803, 1804, 1805, 1806, 1807, 1808, 1809, 1810, 1811, 1812, 1813, 1814