

LAMBDA FEMTOSCOPY IN $\sqrt{s_{NN}} = 2.76$ TeV Pb-Pb COLLISIONS AT ALICE

DISSERTATION

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By

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ABSTRACT

In this thesis, we present lambda baryon femtoscopy measured by the ALICE collaboration in $\sqrt{s_{NN}} = 2.76$ TeV Pb-Pb collision. Correlation functions are shown in three centrality ranges for $\Lambda\bar{\Lambda}$ and two for the combined $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ results. Femtoscopy is capable of measuring the femtometer-scale spatial and temporal aspects of particle collisions by looking at relative momentum correlations among pairs of particles. In addition, femtoscopy can extract information about the final state interactions of the strong nuclear forces that occur between those particles. The femtoscopic radii and scattering parameters (scattering length and effective range of interaction) of $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$ are presented.

Following recent measurements of $p\bar{\Lambda}$ and $\bar{p}\bar{\Lambda}$, we have expanded the standard femtoscopic fitting process to account for the correlated effects of particles that decay into Λ . We have estimated the λ parameters that describe the relative contribution of each residual correlation. Residual correlation effects are included in the fitting process by estimating their correlation strength and then smearing the relative momentum of their contribution according to the kinematics of their decay.

To my loving family (yes, even the ones that meow at me in the middle of the night).

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Publications

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Fields of Study

Major Field: Physics

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Chapter 1

INTRODUCTION

One of the major goals of high-energy physics is to study quantum chromodynamics (QCD), the theory of the strong nuclear force which describes the interactions between quarks and gluons. Quarks and gluons are the building blocks of atomic matter; every proton has two up quarks and a down quark, and every neutron has two down quarks and an up quark. Gluons mediate the interactions between those quarks. Despite their prominence, these particles are never observed in isolation under standard terrestrial conditions. They are always found together in the form of hadrons, which are particles such as protons and pions that have two or three quarks bound together. This phenomenon—the fact the quarks are not normally observed in isolation—is called confinement. However, at sufficiently high density and temperature, it is possible to have a medium that is too hot for hadrons. Instead, the quarks and gluons exist freely, and their collective behavior exhibits fluidlike properties [1]. This medium is called the quark-gluon plasma (QGP). By studying the QGP, as well as the conditions that exist after the QGP dissipates, physicists can get more insight into the strong nuclear force.

Modern particle accelerators create conditions under which the QGP can exist for a short time [2]. This is primarily done via the collision of heavy ions such as ^{197}Au and ^{208}Pb . At energies seen by the Large Hadron Collider (LHC), the ions travel at nearly the speed of light ($c = 3 * 10^8 \frac{\text{m}}{\text{s}}$), and they have center of mass energies per nucleon on the order of tera electron Volts (TeV). The collisions of these ions can produce a hot, dense medium from which thousands of particles are emitted [3]. This medium (the QGP and subsequent hadronic phase) has several names, such as “interaction region” or “particle-emitting source”, though it is colloquially called a fireball. The collision, the medium, and the subsequently emitted and measured particles are collectively known as an event.

Figure 1.1 shows a cartoon of a heavy ion collision. It is suspected that the QGP exists for only a few moments after the collision of these ions, on time scales of about a few fm/ c (femtometers per speed of light, with $1 \text{ fm}/c = 10^{-15} \text{ m}/c \sim 10^{-23}$ seconds). Over time the plasma expands and cools, and eventually the quarks and gluons combine to form hadrons, a

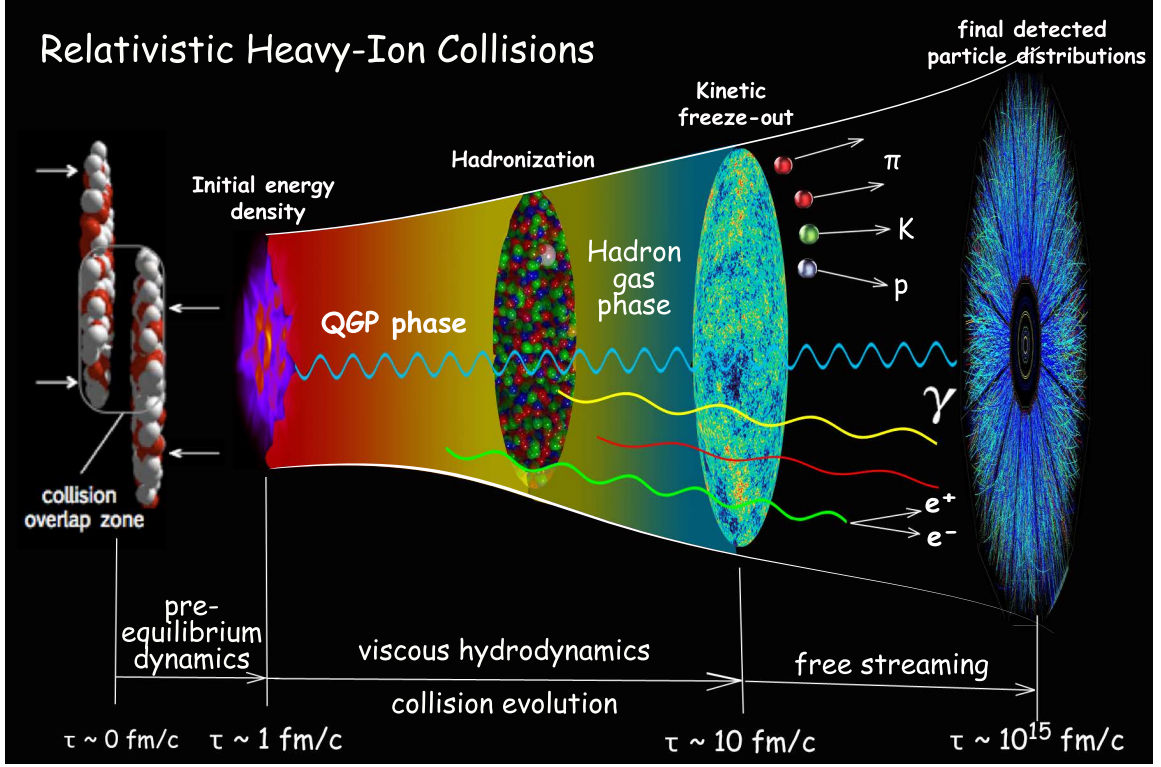


Figure 1.1: Evolution of a heavy ion collision [4]. Two ions collide and leave behind a medium of high energy density. After about 1 fm/c ($\sim 10^{-23}$ s), the medium thermalizes as a plasma of deconfined quarks and gluons. After a few fm/c, the quarks and gluons cool down and hadronize into pions, kaons, and many other particles. These hadrons scatter off each other for a few more fm/c, until they become so dilute that they free-stream to a detector.

process called hadronization. In the subsequent hadronic phase, the hadrons scatter off each other in an elaborate billiards game, with particles near the edge of the expanding fireball having the possibility to free-stream to the detector. Eventually, the source becomes dilute enough that the inter-particle scatterings die off. This occurs on time scales on the order of 10 fm/c [5]. While people often speak of a singular kinetic freezeout time (the time when the rescatterings cease), in actuality particles are emitted at varying rates throughout the lifetime of the fireball.

Heavy-ion physics programs such as ALICE (A Large Ion Collider Experiment) seek to study the evolution of these fireballs. The field tries to answer questions such as “When does hadronization occur?”, “What equations of state describe the evolution of the interaction region?”, “What interactions are occurring between particles during this evolution?”, and “What is the size and shape of the medium at kinetic freezeout?”. A number of models have

been developed to probe these questions and others [6, 7, 8, 9], and experimental results help refine these models.

Generally speaking, collision models have three main components: initial conditions describing the shape and energy density of the source; a theory describing how the system evolves over time; and the final result of collision, namely a description of the particles that are emitted including their momenta, emission time, etc. Collider experiments have been designed to simplify assumptions about the initial conditions of fireballs and to measure their asymptotic state. ^{197}Au and ^{208}Pb ions are used since these particles have a spherical symmetry to their nuclear structure that gives the fireballs a relatively simple initial geometry. Particle collider experiments provide insight into the final state of collisions by measuring the final momenta of particles produced in millions of events.

Beyond simply measuring asymptotic particle momenta, collider experiments can use an assortment of analysis techniques to probe various facets of the fireball evolution. These techniques include flow analysis [10, 11], the study of particle spectra [3], and an interferometry technique known as femtoscopy [12, 13]. In the following sections, we'll give an overview of the LHC and the ALICE detector. We'll then give a short description of a few of the ongoing physics programs. Finally, we'll move on to the main topic of this thesis: femtoscopy.

1.1 Particle collider experiments

1.1.1 The Large Hadron Collider

ALICE is one of several experiments at the Large Hadron Collider (LHC). The LHC itself is part of the European Organization for Nuclear Research (CERN, from the French *Conseil Européen pour la Recherche Nucléaire*). CERN is near Geneva, Switzerland, and the 27 km ring of the LHC lies partly in France and partly in Switzerland.

The LHC is literally built upon previous successful experiments. In order to get protons and ions traveling at very nearly the speed of light, they are sent through a series of increasingly powerful accelerators, each of which was a state-of-the-art machine for particle physics in its day. Figure 1.2 shows a diagram of the LHC and the accelerators that feed into it. Protons and ions are injected into the LHC ring in bunches of up to 10^{11} particles, with bunches traveling in both clockwise and counterclockwise directions around the ring [15]. The bunches are spaced a few nanoseconds apart. There are thousands of bunches per beam in proton-proton mode, while Pb-Pb mode has hundreds. Those bunches are steered around the ring with magnets as strong as 8 Tesla. The two beams are only allowed to intersect at four points, the four major LHC experiments: ALICE, LHC-b, CMS, ATLAS. At these intersection points, there can be millions or even hundreds of millions of proton-proton collisions per second. Ions have collision rates in the thousands.

CERN's Accelerator Complex

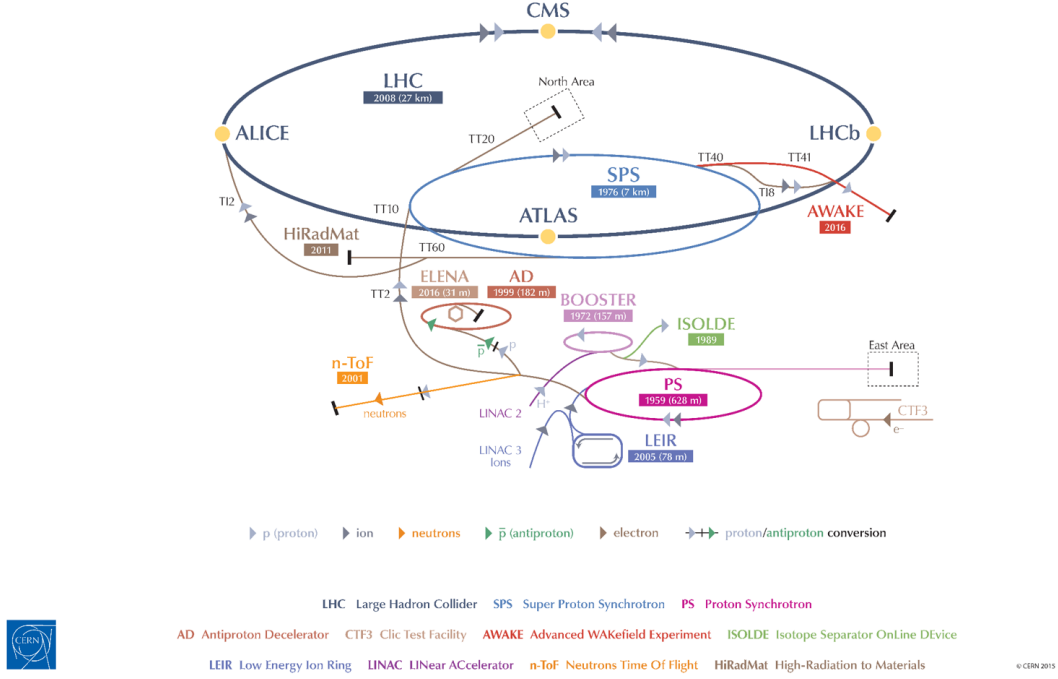


Figure 1.2: A schematic of the Large Hadron Collider [14]. Proton and/or Pb ion beams go through a series of accelerators before being injected into the LHC. Beams travel in both directions around the LHC, and the two beams are only allowed to collide at the four experiments: ALICE, ATLAS, LHC-b, and CMS.

1.1.2 The ALICE detector

Figure 1.3 shows a cartoon of the ALICE detector. The main barrel of the detector is 14.1 m long and 15.8 m in diameter. The outer layer (shown in red in the cartoon), is a 0.5 T solenoidal magnet. The two beams are steered such that most collision events occur within 10 cm of the longitudinal center of the detector (i.e. along the beam axis), and within about a cm of the radial center.

The ALICE detector is composed of many subdetectors [16], several of which are designed to monitor the trajectories of charged particles that pass through them. These reconstructed trajectories are called tracks. The innermost detector is the Inner Tracking System (ITS), which provides full azimuthal coverage with a radial range of 3.9 cm to 43 cm. The ITS consists of six layers of silicon detectors: two Silicon Pixel Detectors, two Silicon Drift Detectors, and two Silicon Strip Detectors. The ITS serves multiple purposes.

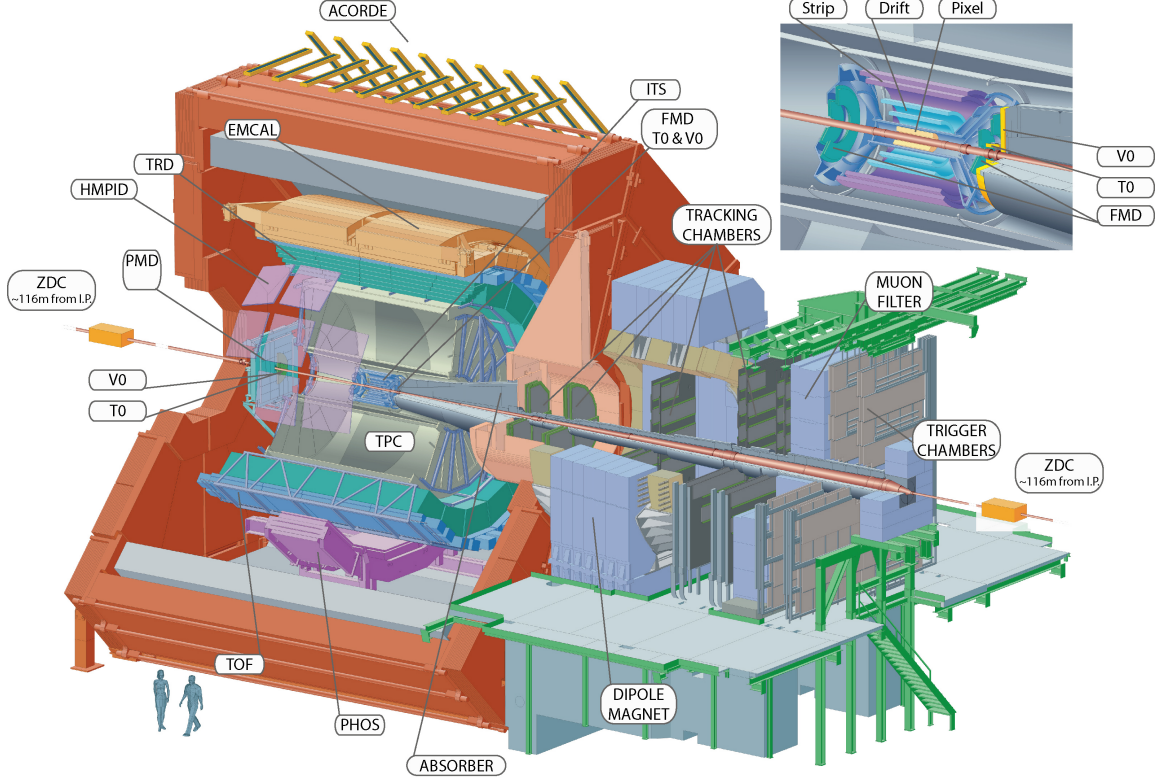


Figure 1.3: A cartoon of the ALICE detector [16].

The spatial resolution of the various layers ranges from tens to hundreds of μm . With hundreds (p-p) or thousands (Pb-Pb) of particles produced in each collision, the ITS is able to reconstruct the position of the collision with a resolution better than $100 \mu\text{m}$. The collision location is called the primary vertex, as nearly all the reconstructed particles trace back to it. The ITS also provides position tracking and particle identification (PID) for low-momentum particles via $\frac{dE}{dx}$ measurement [17]. Finally, the tracking resolution of the ITS allows it to identify secondary vertices, the locations where hyperons (baryons containing a strange quark) and various heavy mesons decay into charged particles.

The next detector is the Time Projection Chamber (TPC) [18], which spans $r = 85$ cm to $r = 250$ cm. The TPC is a 5 m long barrel full of a 90:10:5 mixture of Ne, CO_2 , and N_2 gas. The barrel is divided in half longitudinally by a 100 kV electrode, and there are readout electronics on either end. When a charged particle passes through the TPC, it ionizes the gas in its wake. The electric field then drives the electron clouds toward the readout electronics. These clouds quickly reach their max drift velocity of $2.7 \text{ cm}/\mu\text{s}$ (27 km/s). The position of the track can thus be determined by the length of time it takes for the ionized gas to reach the readout chamber. The 0.5 T magnetic field causes the track's

trajectory to bend, and the particle's momentum can be determined by the curvature of the track. The TPC therefore is able to determine the 3D position of the track throughout the TPC with mm precision, the momentum of the charged particle, and its dE/dx value. This is the main tracking system of the ALICE detector, and it can provide PID information for particles with momentum up to ~ 1 GeV/ c .

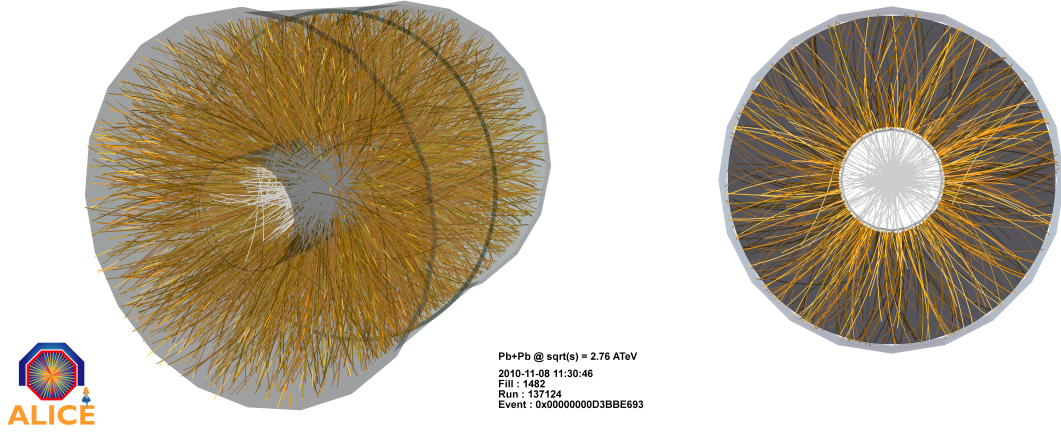


Figure 1.4: An example of a Pb-Pb collision measured at the ALICE detector. Charged particles traveling through the detector leave behind tracks that the detector can measure, each of which is depicted by a colored line. The white region in the center of the detector is the Inner Tracking System, while the surrounding grey region is the Time Projection Chamber. Tracks measured by the ITS are grey, while tracks found by the TPC are yellow. Combined tracking information from both detectors is used to improve the momentum resolution of each track.

Figure 1.4 shows a Pb-Pb collision recorded in the first run of the LHC. The tracking by the ITS is shown in white, while tracking from the TPC is yellow. The TPC can provide quality tracking information for ion events with thousands of charged particles, but it has the drawback of being a slow detector compared to the ITS. The maximum drift time of a track is about $90 \mu s$, so if two or more events occur within that window their tracking information can overlap. This is called event pile-up. The TPC can isolate and record central Pb-Pb events (events with head-on collisions) at a rate of 300 Hz and p-p events at a rate of 1 kHz.

The Time-Of-Flight (TOF) detector is a 20 cm thick shell around the TPC. In conjunction with the T0 detector which determines the timing of the initial collision, the TOF

can measure the time it takes a particle to reach it. With this information and momentum information from other tracking detectors, the mass of the particle can be determined. The TOF can successfully distinguish between particles with momentum as low as 0.5 GeV/ c up to 3-4 GeV/ c , depending on the particle, and it is essential for PID of particles above the 1 GeV/ c limit of the TPC.

The T0, VZERO, and Zero Degree Calorimeter (ZDC) detectors provide event triggering for the ALICE detector. Event-triggering detectors monitor activity in the beam collision area, and they inform the detector when to record data. The ALICE detector therefore does not need to record information for every interaction, including interactions from the occasional cosmic rays striking the detector or the beam striking ambient gas. Instead, it can use event triggers to selectively record events that appear relevant to the designated research goals of the ALICE collaboration. For example, there may be an interest in central collisions, which produce the most particles. The primary triggers used by ALICE with Pb-Pb events are minimum bias, semi-central, and central, each of which will be described in the following section.

1.1.3 Event characterization

It is often useful to classify a heavy-ion collision by how central/peripheral (head-on/glancing) the collision is. Figures 1.6, 1.7, and 1.5 show several ways of quantifying this [19, 10]. One is by impact parameter b , which is the distance from the center of one nucleus to the center of the other. Another is by the total number of participants N_{part} — the nucleons from either ion that interact with nucleons from the other ion. With 208 nucleons in a Pb nucleus, the maximum number of participants is 416. A third way of classifying them is by multiplicity (the number of particles produced in each collision), as central collisions tend to produce more particles than peripheral collisions.

While any of these definitions might be useful from a theoretical standpoint, they have varying degrees of usefulness and accessibility in an experimental setting. Impact parameter, with a scale of femtometers, is inaccessible by experimental measurement. The number of participants can be estimated using detectors such as the ZDC, which measure the number of spectator nucleons (nucleons that are not participants) in a collision. And the number of particles is easily quantified; in practice, experiments often use the number of charged particles in an event as a proxy for the total number of particles in the event. At ALICE, typically only particles found at mid-pseudorapidity $|\eta| < 0.8$ are used in these estimates, where pseudorapidity is defined as

$$\eta = \frac{1}{2} \ln \left(\frac{|p| + p_z}{|p| - p_z} \right), \quad (1.1)$$

and p and p_z are a particle's three-vector momentum and momentum along the beam axis,

respectively.

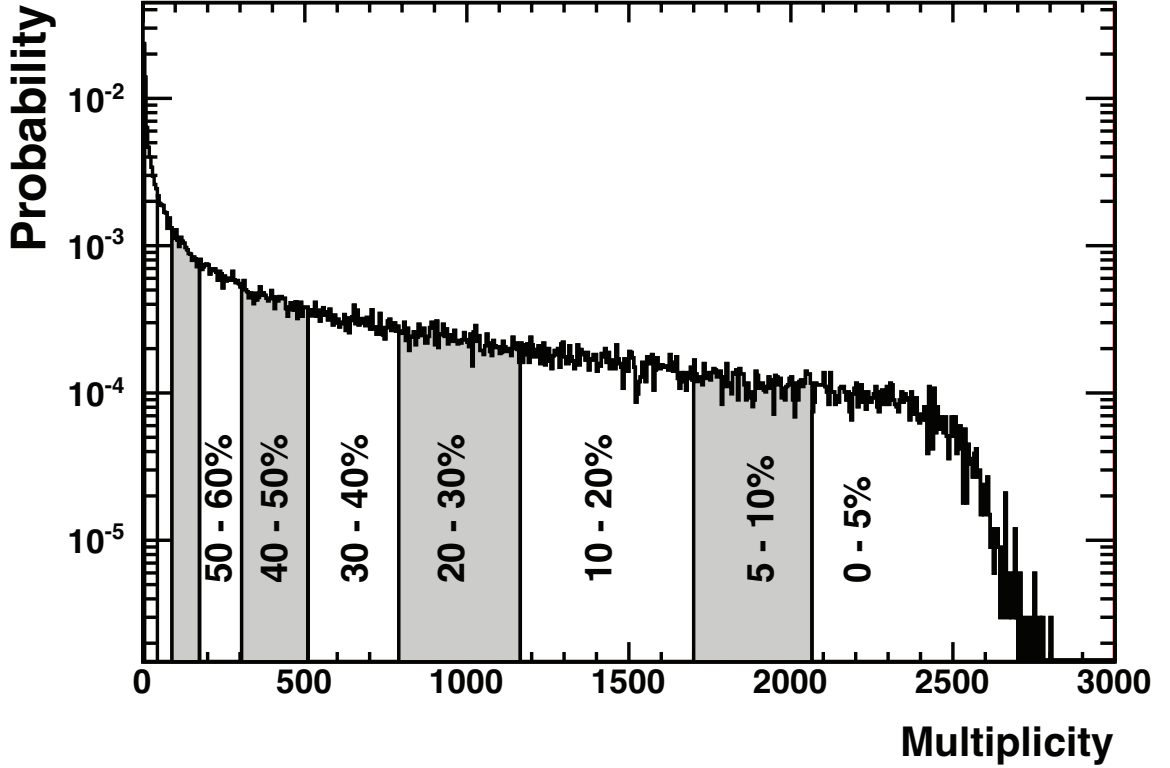


Figure 1.5: Uncorrected charged-particle multiplicity measured by the ALICE TPC at mid-pseudorapidity ($|\eta| < 0.8$) in $\sqrt{s_{NN}} = 2.76$ TeV Pb-Pb collisions [10]. Centrality classes are defined in percentile bins with 0-5% corresponding to the most central collisions with the largest multiplicities and 95-100% corresponding to peripheral collisions with the smallest multiplicities.

To bridge these various ways of classifying collisions, we define centrality classes. There are a number of ways to calibrate centrality, the most common of which is by ranking events in terms of their multiplicity. Central collisions (0-10%) have the highest multiplicity, while peripheral (90% and above) collisions have the smallest, as seen in Figure 1.5 [10]. Centrality can also be mapped onto theoretical calculations of impact parameter (Figure 1.6) and N_{part} (Figure 1.7) [19].

Multiplicity estimates from the TPC, along with timing information and other signals from the ZDC and VZERO, provide event triggering and centrality measurements. The minimum bias trigger records an even sampling of the centrality distribution. The central trigger records only 0-10% central events, and the semi-central trigger records events in

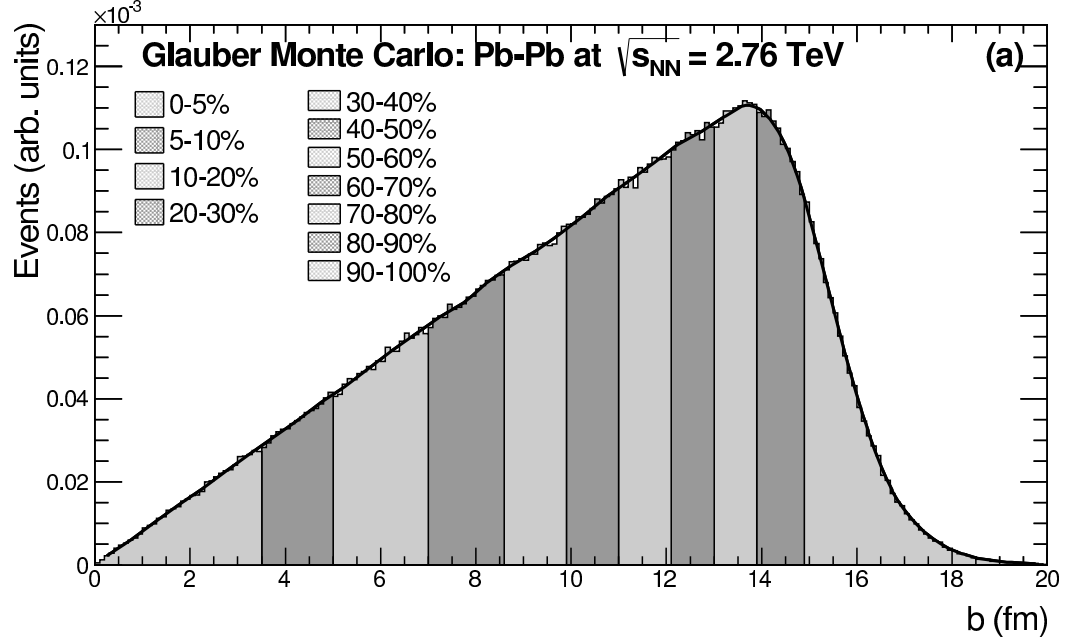


Figure 1.6: Distribution of $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collision impact parameters b from a Glauber Monte Carlo simulation [19]. Impact parameter is defined as the distance between the center of one nucleus to the center of the other at the time of the collision. Bands indicate the impact parameters associated with different centrality classes.

the 10-50% centrality range.

1.2 Heavy-ion analyses

1.2.1 Flow

When two ions collide, the topology of their overlap region has a pronounced effect on the momentum distribution of the particles produced in the collision. Figure 1.8 shows a cartoon of a non-central heavy-ion collision. At LHC energies, the colliding ions are Lorentz length-contracted by a factor of ~ 1500 , effectively making them flat disks or pancakes with non-uniform density (more nucleons in the center than toward the edge). When these two pancakes collide with a non-zero impact parameter, the overlap region of their participant nucleons is roughly almond-shaped. As shown in the figure, the almond-shaped fireball is taller than it is wide. This creates a larger pressure gradient along the axis of the impact parameter (the short axis) than along the long axis. This drives a faster expansion along the short axis as the fireball evolves, a phenomenon called anisotropic flow. In this way, the

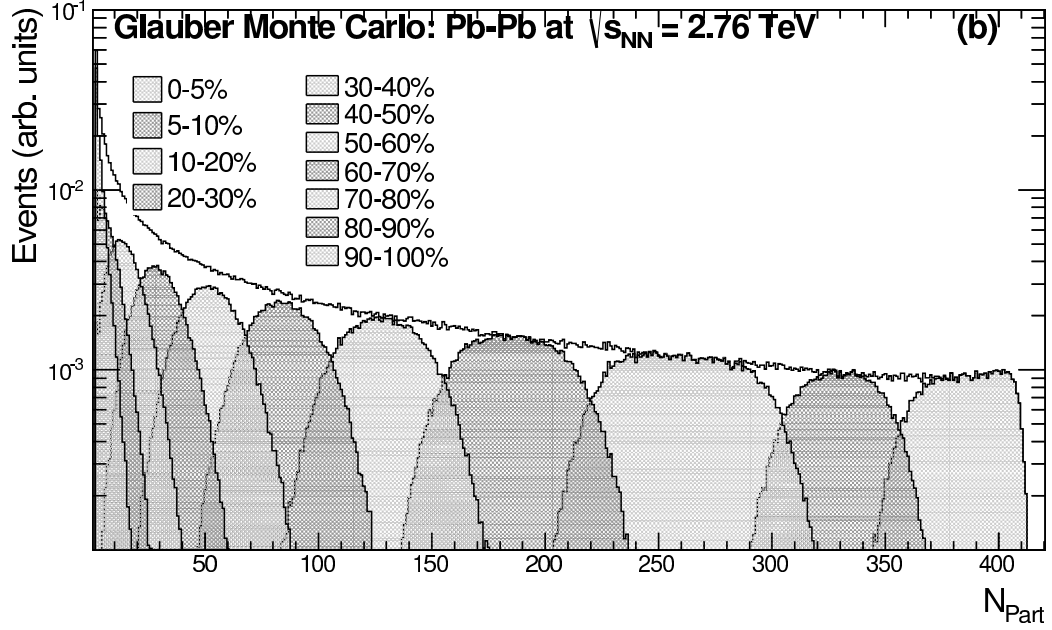


Figure 1.7: A Glauber Monte Carlo simulation of the total number of nucleons involved in $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collisions [19]. A nucleon from one ion that interacts with a nucleon from another is called a participant. In a peripheral collision, each ion only has a handful of participants. Central collisions may have ~ 200 participants per ion. The estimated participant ranges of the centrality classes are shown.

initial shape of the collision impacts the final momentum distribution of the event.

Using a Fourier decomposition of the azimuthally-dependent momentum distribution [21], anisotropic flow can be quantified in terms of flow coefficients v_n

$$E \frac{d^3N}{d^3p} = \frac{1}{2\pi} \frac{d^2N}{p_T dp_T dy} \left(1 + 2 \sum_{n=1}^{\infty} v_n \cos[n(\phi - \Psi_{\text{RP}})] \right), \quad (1.2)$$

where E , p , $p_T = \sqrt{p^2 - p_z^2}$, p_z , and ϕ are the energy, momentum, transverse momentum (radially outward from beam axis), longitudinal momentum (along beam axis), and azimuthal angle of a particle's momentum as measured in the lab frame, respectively; $E \frac{d^3N}{d^3p}$ is the momentum distribution of measured particles; $y = \frac{1}{2} \ln \left(\frac{E+p_z}{E-p_z} \right)$ is a particle's rapidity; and Ψ_{RP} is the azimuthal angle of the reaction plane, which is defined by the collision impact parameter (vector pointing from the center of one nucleus to the center of the other) and beam direction. The Fourier components v_n are then defined as

$$v_n(p_T, y) = \langle \cos[n(\phi - \Psi_{\text{RP}})] \rangle \quad (1.3)$$

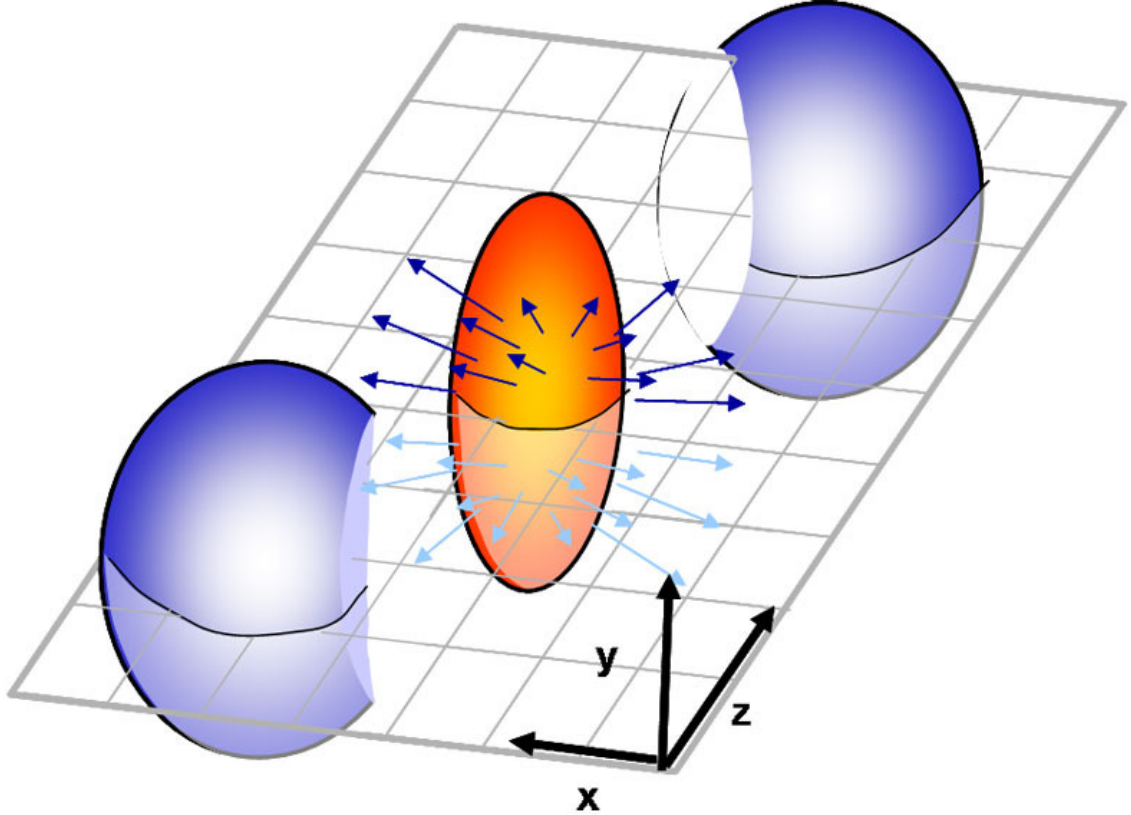


Figure 1.8: Cartoon of a non-central heavy-ion collision [20]. The blue shapes are the recently-collided ions, while the overlap region of the expanding fireball is orange. The almond-shaped fireball is taller than it is wide, which results in a larger pressure gradient along the x -axis than along the y -axis. This pressure gradient causes the fireball to expand more rapidly along the x -axis than along the y -axis.

where the angle brackets represent an average over all particles and sum over each event for a given p_T and y bin. Elliptic flow and triangular flow, v_2 and v_3 , are the most studied coefficients. Elliptic flow owes much of its existence to the initial spatial anisotropy of the almond-shaped fireball, but triangular flow comes from event-by-event energy density fluctuations in the interaction region.

The study of anisotropic flow helps us understand the initial conditions and the evolution of heavy-ion collisions. Theorists have had a great deal of success modeling this evolution, and the evolution of the QGP phase in particular, using relativistic hydrodynamic calculations [11]. This has led some to declare that the quark-gluon plasma is the “most perfect fluid” [1], due to its adherence to ideal fluid dynamics. Nonetheless, viscous effects have been found to have significant effects on the flow coefficients for particles with large transverse momentum. Flow studies continue to be an important tool in understanding the

dynamics of heavy-ion collisions.

1.2.2 Spectra

If the fireball of a heavy-ion collision is in an evolving thermal equilibrium as hydrodynamic models suggest, then it should be possible to quantify the temperature of the fireball at different points in its lifetime. Some temperature information can be learned from spectral analyses, which measure the number of particles produced in collisions. There are a few ways to dissect spectra. For each species of particle (e.g. π , K, p), one can look at how many of them are produced on average for a given event. For a system in thermal equilibrium, the expected number of particles n_i of a particular species in a given energy state E_i follows the distribution

$$n_i(E_i) = \frac{g_i}{e^{(E_i - \mu_b)/kT} \pm 1}, \quad (1.4)$$

where g_i is the degeneracy of the state, μ_b is the baryochemical potential of the system, T is the temperature of the system, and the $+/-$ is for fermions/bosons. The baryochemical potential quantifies the baryon-antibaryon asymmetry of the system; $\mu_b = 0$ means that there are equal numbers of baryons and antibaryons.

Figure 1.9 shows particle yields measured in central Pb-Pb collisions at ALICE. Those yields have been fit with a thermal model [22]. The fit finds that the baryochemical potential is zero, as one expects at LHC energies. Most of the particle species appear to obey a thermal distribution for $T = 156$ MeV, which suggests that relative particle yields cease to evolve below that temperature. That said, the proton and antiproton yields both lie a couple of sigma below the thermal value, and studies are ongoing to determine what physics effects might drive these particles out of equilibrium.

1.3 Femtoscopy

1.3.1 Origin of femtoscopy (HBT)

Hanbury-Brown and Twiss first used two-particle interferometry of photons to measure the angular size of the star Sirius [24], as illustrated in Figure 1.10. Later, it was observed that similar techniques could be employed in a laboratory setting with pions (π) to characterize the size of particle collisions [25] (Figure 1.11). This process of measuring two-particle correlations for pairs of particles with very similar momentum is often referred to as HBT, after the progenitors of the technique. HBT was originally used to study photons and later used to look at pions, so it often carries a connotation of Bose-Einstein correlations, i.e. the quantum mechanical tendency of identical bosons to exist in the same state. This is in contrast to Fermi-Dirac statistics, wherein identical fermions cannot occupy the same state because of the Pauli exclusion principle. Modern analyses look at all manner of hadrons, and

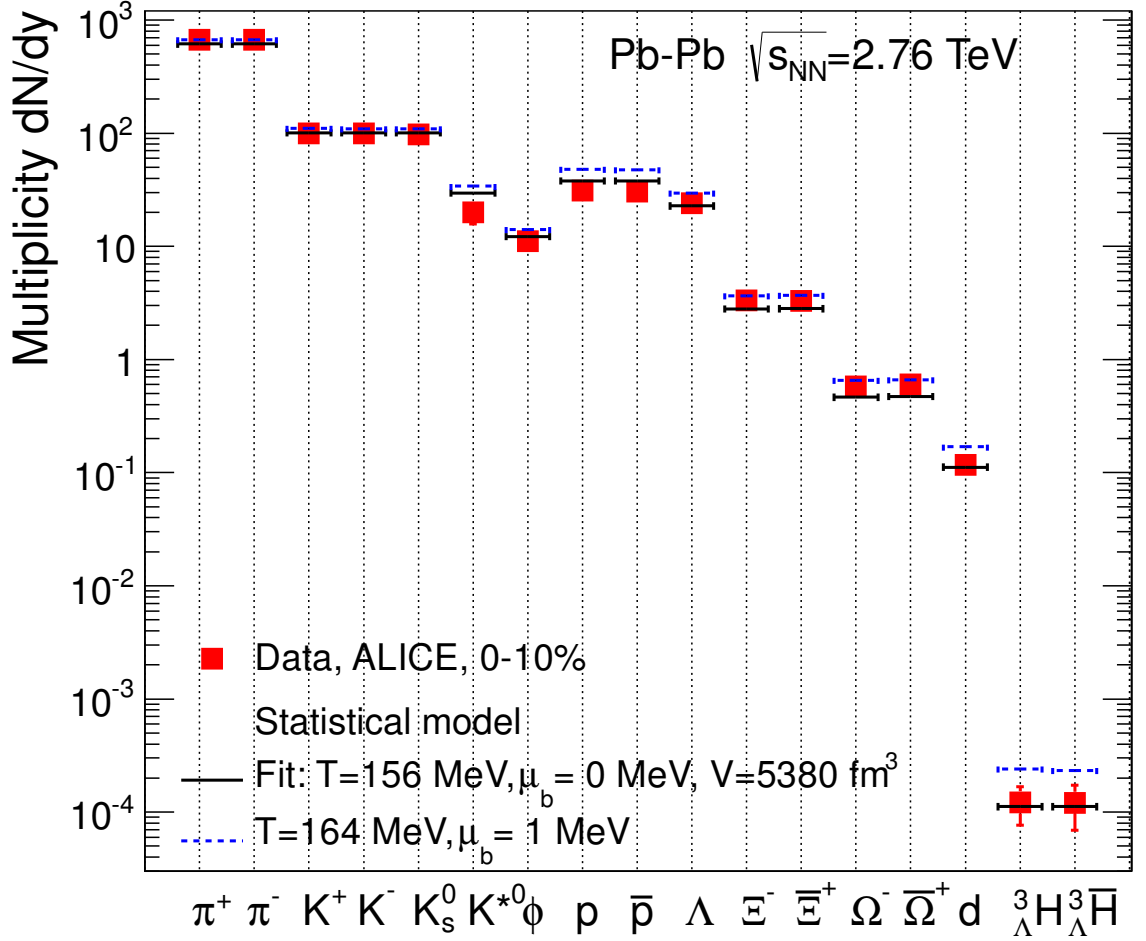


Figure 1.9: Particle yields from central Pb-Pb collisions at ALICE, which are fit with a statistical hadronization model [22]. The fit suggests that the particle abundances, which mostly appear to obey a thermal distribution, cease to evolve below the temperature $T = 156$ MeV. The fit also finds that the baryon chemical potential is zero, which means that there is no asymmetry between the matter and antimatter yields of the system.

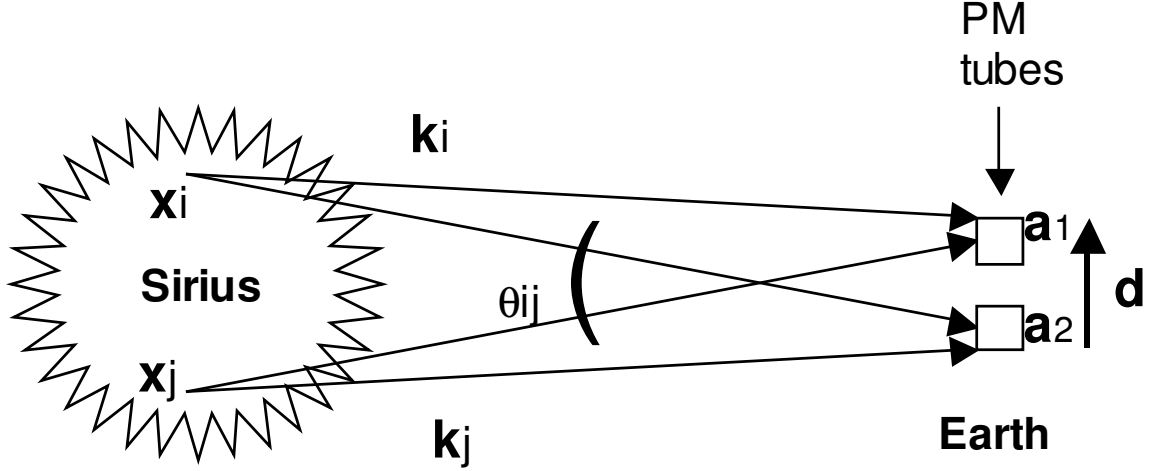


Figure 1.10: Schematic of Hanbury-Brown-Twiss interferometry as used to measure the angular size of the star Sirius [23]. Photons with wave-vectors k_i and k_j are emitted from the star at positions x_i and x_j with angular separation θ_{ij} . The photons are then measured in coincidence by photomultiplier tubes on Earth spaced a distance d apart. By measuring the interference pattern of the light for different photomultiplier tube separations d , one can obtain the average angular size of the star $\theta = \langle \theta_{ij} \rangle$.

therefore the correlations might be subject to bosonic, fermionic, or non-identical particle statistics. In the nuclear physics context, the focus of these analyses is on measuring very small sources, so the more general term *femtoscropy* has come into use.

1.3.2 Femtoscopic correlation functions

The main tool of femtoscopy is the correlation function $C(k)$, which reveals correlations among pairs of particles as a function of the momentum difference $k^\mu = (p_1^\mu - p_2^\mu)/2$ of those particles (this form of k^μ holds for identical particles; the general form is given in Equation 2.2). Correlation functions are typically normalized such that a value of unity for a given bin suggests that there are no correlated effects between the particles in that bin. A value above unity means that some effect is causing an increase of pairs at a given relative momentum, such as the Bose-Einstein enhancement of identical pions. A value below unity (but never below zero) suggests that an effect is suppressing pairs at that relative momentum, such as the Coulomb repulsion of of same-sign pions. The Coulomb attraction/repulsion of opposite-sign/same-sign pions is visible in Figure 1.12, which shows pion correlation functions measured in $\sqrt{s_{NN}} = 2.76$ TeV Pb-Pb collisions by ALICE [26]. Here, the two-particle correlation functions C_2 are shown for 0-5% central collisions as a function of $q \equiv 2\sqrt{-k^\mu k_\mu}$. Each panel shows correlation functions constructed with pairs

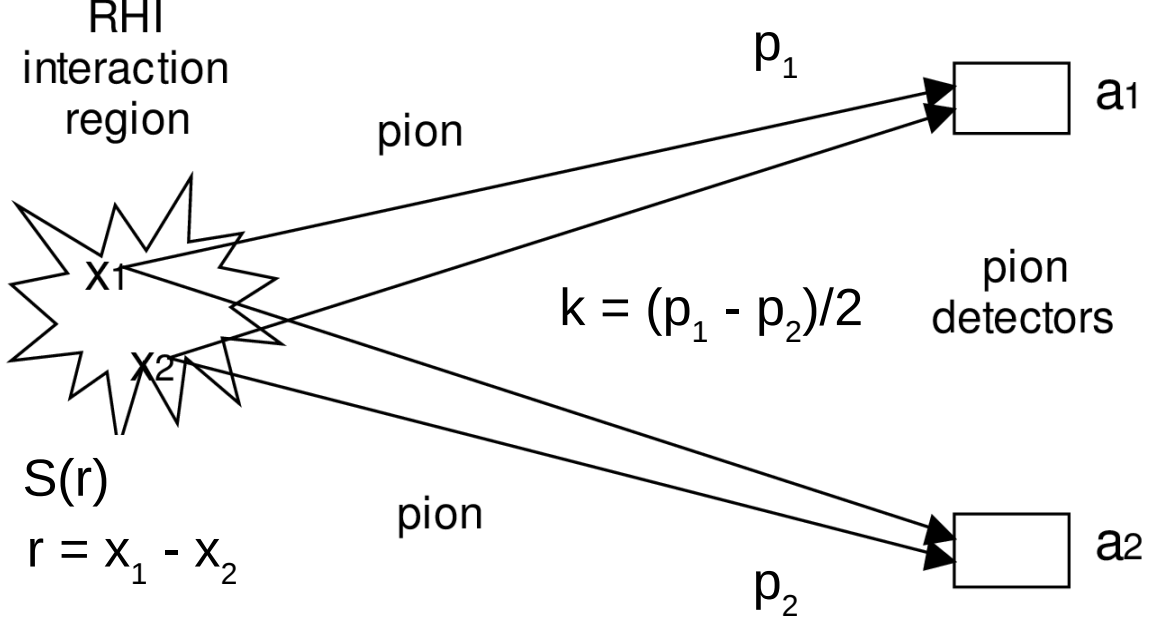


Figure 1.11: Schematic of Hanbury-Brown–Twiss interferometry (femtoscopy) as used to measure the size of the interaction region in relativistic heavy-ion collisions [23]. Pions with momenta p_1 and p_2 are emitted from points x_1 and x_2 , a separation of $r = x_1 - x_2$. The momenta of the pions are measured by a particle detector such as ALICE. By looking at correlations of many pairs of pions, one can measure the size of the two-particle emission function $S(r)$, which describes the separation distribution of the emission points of pairs of particles. In practice, only pairs of particles with small relative momentum $k = |p_1 - p_2|/2 \lesssim 100 \text{ MeV}/c$ are sensitive to the size of the emission region.

from a different average transverse momentum range $k_T = (p_{T,1} + p_{T,2})/2$ (pair transverse momentum is usually denoted as k_T , but throughout this thesis, we will use $P_T \equiv k_T$ to avoid confusion with the momentum difference k). The figure also includes a fit to the data using a Gaussian parametrization of the source size and a non-Gaussian fit with Edgeworth coefficients [27] (more on femtoscopic sources in Sections 1.3.3 and 3.1.2).

The correlation function is constructed as a ratio of a signal distribution over a background distribution. Both distributions are binned in k , sometimes with k as vector and sometimes as a magnitude. The signal distribution sums over pairs of particles where both particles come from the same collision event, aggregated over many events. The background distribution is created to closely match the phase space of the signal distribution, but without any correlated physics effects like Coulomb attraction/repulsion. The uncorrelated, relative-momentum phase space is the combinatoric momentum difference of the single-particle spectra of the particles being studied, $\frac{dN_1}{d^3p}$ and $\frac{dN_2}{d^3p}$. In order to approximate

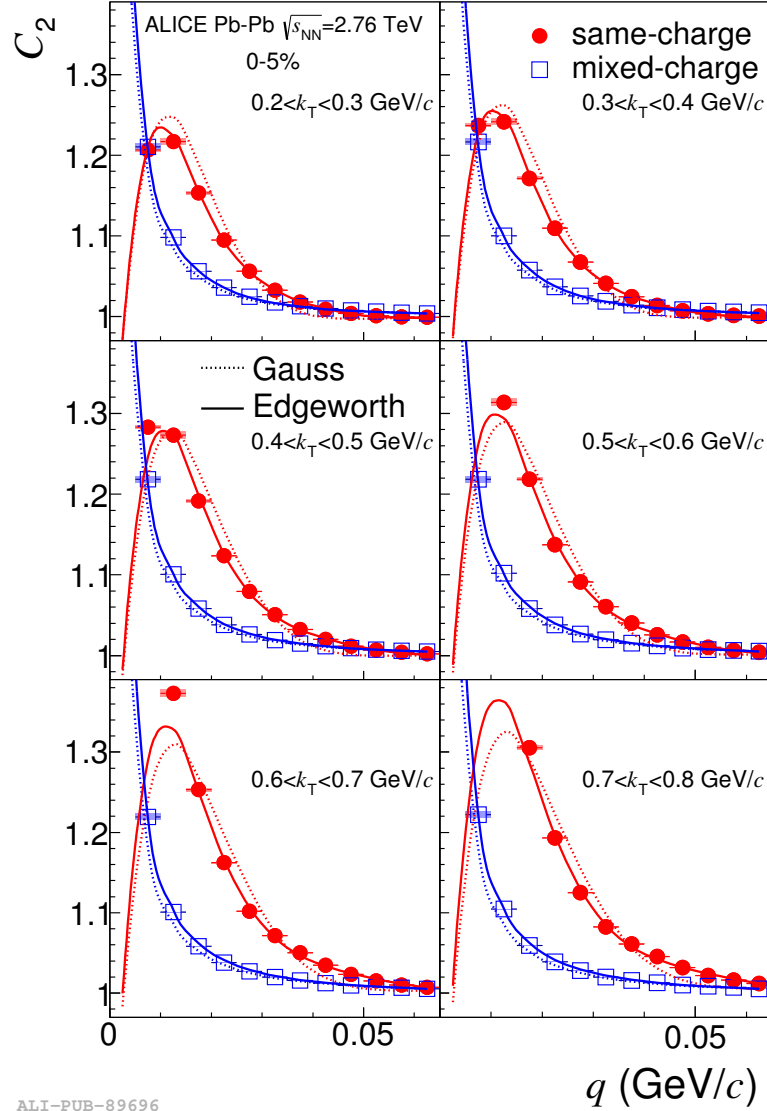


Figure 1.12: Pion correlation functions measured in $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collisions by ALICE [26]. Two-particle correlation functions C_2 are shown for 0-5% central collisions as a function of the relative momentum $q \equiv 2\sqrt{-k^\mu k_\mu}$. Each panel shows correlation functions constructed with pairs from a different average transverse momentum range $k_T = (p_{T,1} + p_{T,2})/2$. Correlation functions of mixed-charge pions (blue boxes) have a strong enhancement (deviation from $C(q) = 1$) at low q arising from Coulomb attraction. Same-charge pions (red circles) have an enhancement from bosonic quantum interference and a dip at very low q from Coulomb repulsion. The figure also includes a fit to the data using a Gaussian parametrization of the source size and a non-Gaussian fit using Edgeworth coefficients [27].

this phase space without correlated physics effects, the background distribution is created from “mixed events”—events where the first particle in the pair is taken from one collision, and the second particle is taken from a different collision. The two events being mixed must be very similar to have their spectra match; they are chosen from the same centrality bin and same longitudinal position bin of the detector. If possible, the two events should even have occurred close to each other in time (hours or days), because detector conditions such as hardware failure can affect the measured spectra. With a carefully chosen background distribution, uncorrelated phase space effects are assumed to divide out, and the correlation function can reveal correlated physics effects. More details on the experimental construction of correlation functions can be found in Section 2.3.

Femtoscopy seeks to exploit the correlation between the relative position of particles at their moment of last scattering and their subsequent momentum at freezeout. The position x of a particular particle at freezeout samples the single-particle emission function $S(x, p)$. This single-particle emission function is directly related to the momentum spectra of the particles

$$E_p \frac{dN}{d^3p} = \int d^4x S(x, p).$$

The relative position of two particles at freezeout $\vec{r} = \vec{x}_1 - \vec{x}_2$ comes from the convolution of the single-particle emission functions of the particles $S_1(x_1, p_1)S_2(x_2, p_2)$ (for identical particles, $S_1 = S_2$). As will be described in further detail in Chapter 3.1, we can also refer to the *two-particle* emission function $S(r, P)$, where $P = p_1 + p_2$.

The final state of the particles, including their momenta, depends on the interactions between the particles and is accounted for by the two-particle wave function of the pair. The correlation function $C(k^*, P)$ is the convolution of the emission function and the wave function:

$$C(k^*, P) = \int d^3r S(r, P) |\Psi(k^*, r)|^2. \quad (1.5)$$

Here $k^* = \sqrt{-k^\mu k_\mu}$ is the magnitude of the relative momentum and $\Psi(k^*, r)$ is the properly symmetrized two-particle wave function of the pair. The wave function acts to Fourier transform the emission function from a relative position space to relative momentum space.

The femtoscopic correlation function can shed light on both the emission function and the interactions that make up the wave-function. In the following two sections, we’ll discuss the source functions and the interactions in greater detail.

1.3.3 Femtoscopic sources

Femtoscopy is built on measuring the interference patterns of particles emitted incoherently from a source, be it a star or a heavy-ion collision. As will be discussed in greater detail in Section 3.1, femtoscopy does not actually measure the size of the full interaction region

of a heavy-ion collision at the time of freezeout. Rather, it measures the size of the region in which particles are emitted with approximately the same momentum, also called the region of homogeneity [28]. If particles were emitted isotropically throughout the interaction region, the measured source size would be the full size of the interaction region. This would be the case for a non-expanding source, one where particles don't interact. If hadrons could not scatter off each other, then they would be emitted from the interaction region isotropically, as seen in the cartoon on the left of Figure 1.13. However, effects such as radial flow create correlations between particle emission direction and position. The pressure of the fireball boosts the momentum of particles outward, as seen on the right hand side of Figure 1.13. Particles can still freeze out with an inward-pointing momentum, but the probability of that happening will be suppressed. In simple terms, a particle near the edge of the fireball is more likely to be emitted outward, where there are fewer particles to scatter off of, than it is for it to be emitted inward where it would have to travel through the bulk of the interaction region without scattering again.

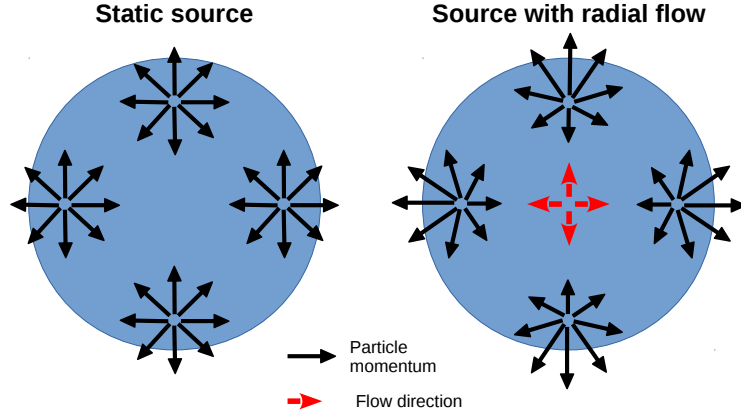


Figure 1.13: A illustration showing possible emission momenta for particles at different places in the fireball. (Left) In a static source where no particles interact, particles are emitted isotropically, and there are no rescatterings to change their momenta. (Right) In an evolving source, the pressure of the fireball creates an outward radial flow of particles. Particles are more likely to be emitted with their momenta pointing outward.

Figure 1.14 illustrates the effect of radial flow, which causes the size of homogeneity region to depend on the average transverse-momentum $P_T = (p_{T,1} + p_{T,2})/2$ of the particles being studied. In the figure, a hadronic rescattering model [23] has been used to determine the freezeout positions of pions in S-Pb collisions at SPS energies. The left plots show pions with $0 < p_T < 100$ MeV/c, while the right plots show pions with $500 < p_T < 600$ MeV/c.

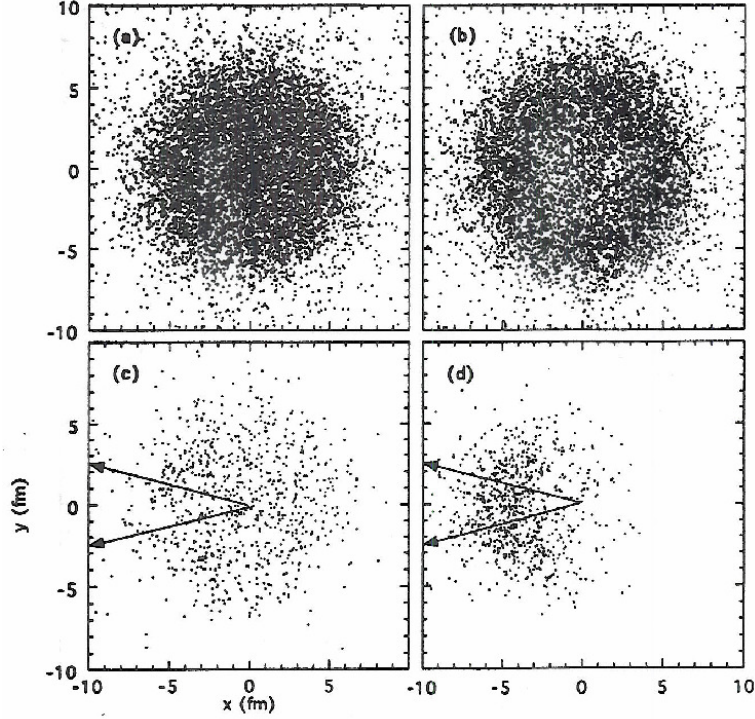


Figure 1.14: Pion freezeout positions in the transverse plane from a hadronic rescattering model [23] of S-Pb collisions at the Super Proton Synchrotron with various momentum cuts applied. The left plots show pions with $0 < p_T < 100$ MeV/c, while the right plots show pions with $500 < p_T < 600$ MeV/c. The top plots show pions with any azimuthal emission direction, while the bottom plots show only pions emitted between $165^\circ < \phi < 195^\circ$, as indicated by the arrows. From the top plots, one sees that low-momentum pions are emitted throughout the fireball, while high-momentum particles tend to be emitted near the outside of the fireball. From the bottom plots, one sees that low-momentum pions travelling generally in the negative x -direction come from all over the fireball, while high-momentum particles are mostly emitted from the left side. From these plots, one can get a qualitative picture of femtoscopic regions of homogeneity, regions where pairs have approximately the same momentum. The region of homogeneity is essentially the size of the fireball for the low-momentum particles, while the region is distinctly smaller for the high-momentum particles.

The top plots show pions with any azimuthal emission direction, while the bottom plots show only pions emitted between $165^\circ < \phi < 195^\circ$, as indicated by the arrows. From the top plots, one sees that low-momentum pions are emitted throughout the fireball, while high-momentum particles tend to be emitted near the outside of the fireball. From the bottom plots, one sees that low-momentum pions travelling generally in the negative x-direction come from all over the fireball, while high-momentum particles are mostly emitted from the left side. From these plots, one can get a qualitative picture of femtoscopic regions of homogeneity, regions where pairs have approximately the same momentum. The region of homogeneity is essentially the size of the fireball for the low-momentum particles, while the region is distinctly smaller for the high-momentum particles.

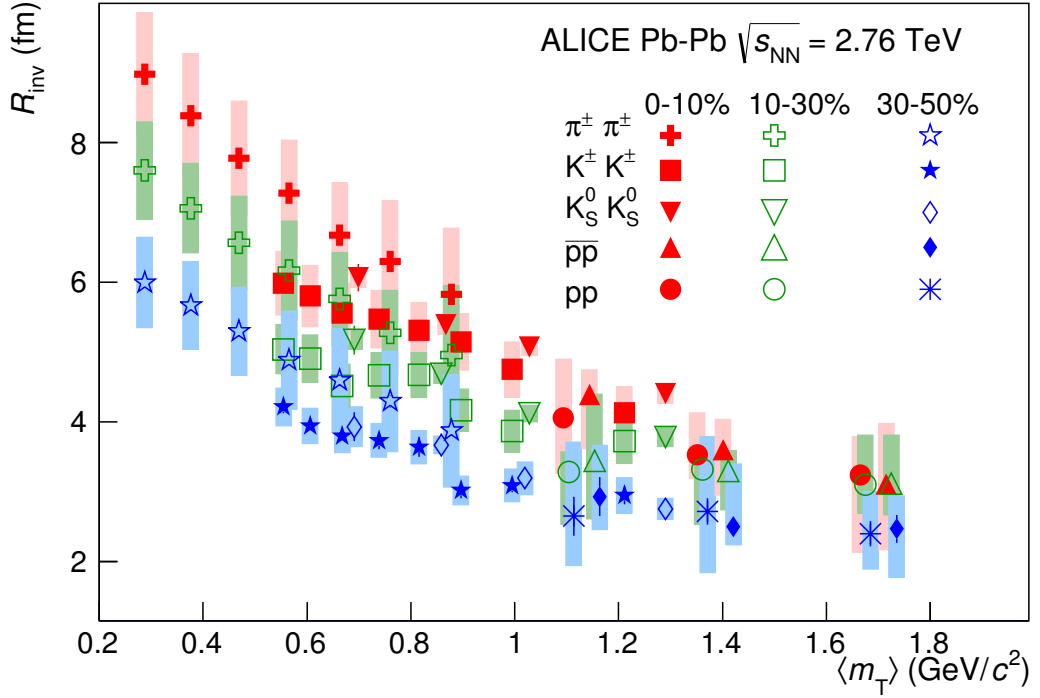


Figure 1.15: Femtoscopic source radii R_{inv} vs transverse mass $m_T = \sqrt{m^2 + P_T^2}$ measured by ALICE in $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collisions [29]. The radii are measured for pions, kaons, and protons using a one-dimensional Gaussian parametrization. The colors show different centrality classes: red, green, and blue for 0-10%, 10-30%, and 30-50% respectively. There is a strong centrality dependence to the radii, with more central events having larger radii. Within each species, there is also a clear trend of the radii falling with increasing transverse mass, though there is at best a rough transverse mass scaling between species.

Femtoscopic analyses traditionally study pions [25, 12] because of their ready availability. However, measurements of heavier particles such as kaons [30] and baryons [31] can serve to complement the pion results. One motivation for studying an assortment of heavier particles is to test the hydrodynamic prediction that radial flow should cause the source radii to scale not with the transverse momentum P_T of the pair, but actually with the transverse mass $m_T = \sqrt{m^2 + P_T^2}$ of the particles [32, 13], where m is mass of the particles being studied. To be more precise, radial flow is expected to cause the size R of the homogeneity regions to fall roughly as $R \propto \sqrt{T/m_T}$, where T is the kinetic freezeout temperature of the fireball. Figure 1.15 shows femtoscopic radii measured by ALICE for pions, kaons, and protons. There is a strong m_T -scaling trend within each species. The trend is looser between species; higher mass particles do have smaller source radii than lower mass particles, but the radii of different species are not consistent at a given transverse mass. There is some dispute about whether radii should be dependent on P_T , m_T , or some nonlinear combination of the two [33]. Measurements of different species of particles at various P_T and m_T values may help to resolve this issue.

As we have discussed in this section, femtoscopy can measure the size of homogeneity regions in fireballs. In doing so, it yields a snapshot of the final size and shape of the interaction region. While it isn't a complete picture of the fireball, it is nonetheless the best picture experimentally available.

1.3.4 Hadronic interactions

Femtoscopic correlation functions are sensitive to the interactions between particles. For identical, charged pion analyses, the quantum interference effects and large Coulomb interactions drown out the interactions of the strong nuclear force. However, for pairs of particles such as protons (pp) or neutral kaons ($K_s^0 K_s^0$), the strong interactions are significant. As we will discuss in Section 3.1.3, these interactions are taken into account via the two-particle wave function. For some pairs of particles, the strength and range of their interaction is well known. For example, decades of experiments have provided a wealth of pp scattering information [34]. While there are some predictions for hadron scattering lengths coming from chiral perturbation theory [35], there are many hadron-hadron pairs for which little to no experimental data is available. And the interactions of hadron-antihadron pairs have been explored even less.

More information on these interactions can improve several disparate areas of research, including rescattering calculations used in modeling the evolution of heavy-ion collisions [36]; binding energy models used in hypernuclear structure theory [37, 38]; and, in the case of the lambda baryon (Λ), the equations of state which describe the properties of neutron stars [39, 40].

One place where the absence of scattering information is relevant is with the Ultrarelativ-

ativistic Quantum Molecular Dynamics model (UrQMD) [36]. UrQMD simulates particle creation and subsequent evolution in heavy-ion collisions. It is also used in conjunction with other heavy-ion collision models [41, 42, 43]; those models simulate the early- and mid-stage qualities of events via macroscopic transport models such as hydrodynamics, and then use UrQMD to handle the late hadronic stage of the collisions. UrQMD calculates the positions and momenta of particles at fixed time steps, and it uses known particle cross-sections to determine how particles scatter off each other. If no experimental cross-section data is available for collisions of a particular pair of hadrons, then UrQMD extrapolates from known data. In the case of baryon-antibaryon pairs, the annihilation cross-sections at a given center of mass energy $\sqrt{s}$ are all treated as having the same annihilation cross-section as proton-antiproton pairs at the same $\sqrt{s}$. Femtoscopy gives us a tool with which to measure these interactions so as to improve the accuracy of future rescattering models.

There have been several recent femtoscopy success stories at the Relativistic Heavy Ion Collider (RHIC) from the STAR Experiment (Solenoidal Tracker at RHIC). In one study of $\bar{p}p$ pairs, STAR measured the s-wave scattering parameters of $\bar{p}p$, and found that they are consistent with those of pp [44]. In another study, STAR saw unexpected differences between the measured femtoscopic radii of $p\Lambda$ and $p\bar{\Lambda}$ pairs, with the latter pairs being about a factor of two smaller [45]. A reanalysis of that data implemented a new method for accounting for signal contaminations in the correlation functions [46]. Rather than correct for impurity with a simple scaling factor, the new method attempted to quantify the effect of the residual correlations (the femtoscopic correlations arising from the contaminations). With the new method in place, the extracted $p\bar{\Lambda}$ radius was consistent with the $p\Lambda$ radius ($R_{p\bar{\Lambda}} = 2.83 \pm 0.12$ fm vs. $R_{p\Lambda} = 3.09 \pm 0.3^{+0.17}_{-0.25} \pm 0.2$ fm). The analysis presented in this thesis, $\Lambda\Lambda$ and $\Lambda\bar{\Lambda}$ femtoscopy, uses this approach for handling residual correlations; the details can be found in Sections 3.2 and 3.3.

1.3.5 Λ femtoscopy

The frontier of femtoscopy lies in studying heavier particle pairs. By virtue of their larger mass, particles such as p , Λ , and the cascade baryons (Ξ^- and Ξ^0) can yield femtoscopic radii at much larger m_T than are accessible via pions. In addition, femtoscopy gives access to the scattering parameters of these particles, information that is sorely needed for a number of heavier particle pairs. Finally, the process of identifying V0s, neutral particles such as Λ and K_s^0 whose decay into charged particles looks roughly “V”-shaped, differs significantly from the process of identifying charged particles. Figure 1.16 shows a cartoon of a Λ decay. Because of the differences in their analysis techniques, V0 particle femtoscopy and charged particle femtoscopy can therefore provide useful consistency checks for each other, as seen in the recent charged and neutral kaon analyses at ALICE [29].

As mentioned above, there is a need for experimental study of exotic particle interac-

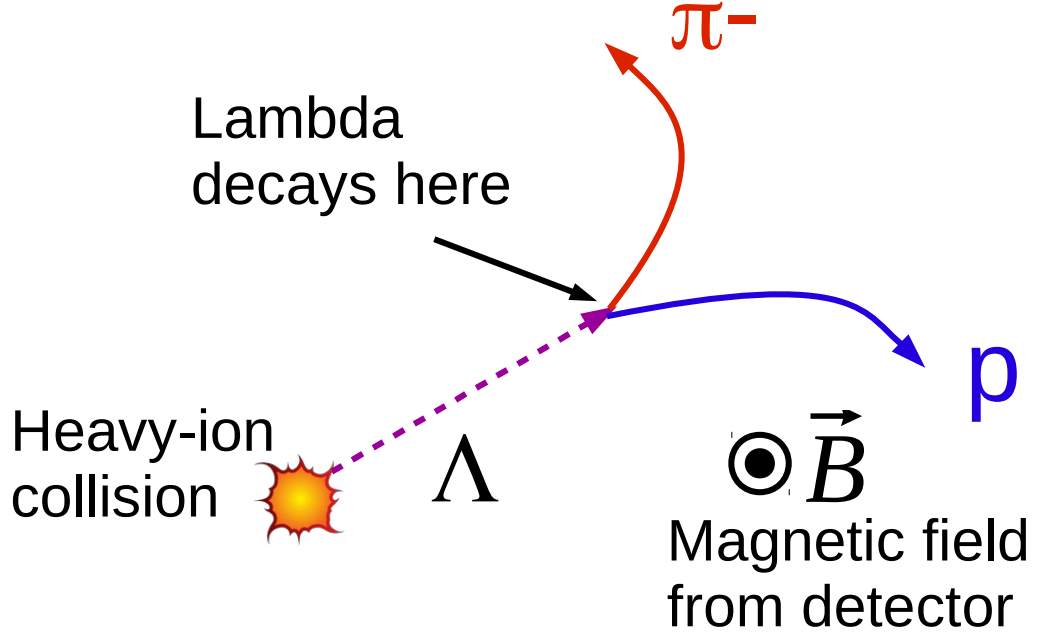


Figure 1.16: Cartoon of the decay of a lambda baryon (Λ). A Λ moving near the speed of light decays into a proton and a pion after traveling a few centimeters away from the fireball in which it was created. The Λ is charge-neutral, so it is not detected directly; instead it is reconstructed by looking for pairs of charged particles that fit a certain decay topology. See Section 2.2 for more details about Λ reconstruction.

tions. For example, there is a sizable disparity in the $p\Lambda$ singlet-state scattering length f_0 between the leading order ($f_0 = 1.9$ fm) and next-to-leading order ($f_0 = 2.9$ fm) chiral effective field theory calculations [47]. There are only a handful of measurements of the $\Lambda\Lambda$ scattering length: one via the decay of a hypernucleus (a nucleus which contains one or more strange particles) [48, 38, 37], and one via femtoscopy at STAR [49], and another via femtoscopy at the Super Proton Synchrotron (SPS) [50]. The SPS experiment did not have enough data to resolve any nuclear interaction. The other results are inconsistent with each other — their scattering lengths differ in sign, tantamount to a difference between attractive and repulsive interactions — so further measurements are necessary.

Research on Λ particles is statistically quite demanding compared to similar research on pions. In $\sqrt{s_{NN}} = 2.76$ TeV central Pb-Pb collisions at the LHC, there are about 4 Λ created per 100 π [51]. Nonetheless, modern particle colliders, including RHIC and the LHC, are in a unique position to perform Λ femtoscopy. At the LHC, Pb-Pb collisions

at $\sqrt{s_{\text{NN}}} = 2.76$ TeV have mid-rapidity Λ yields ~ 2.5 times larger than those measured in the aforementioned fixed-target SPS experiment, which had Pb-Pb collisions with a beam energy $E = 158$ A GeV ($\sqrt{s_{\text{NN}}} \approx 17.3$ GeV) [19, 52]. But the relevant quantity for femtoscopy is number of pairs, not number of particles, so the 2.5 times increase in particles leads to about 6 times more pairs. Moreover, at the LHC, the baryochemical potential $\mu_b \approx 0$ MeV [22], meaning that the ratio of baryons to antibaryons is unity. Therefore, the statistics for $\bar{\Lambda}$ analyses should be comparable to those for Λ analyses.

Several physics effects can be seen in Λ and $\bar{\Lambda}$ correlation functions. The primary effect seen in $\Lambda\bar{\Lambda}$ correlations should be the final state interactions (FSI) of the strong nuclear force. There should be a depletion of baryon-antibaryon pairs at low relative momentum due to annihilation channels. Suppressions of this sort have been seen in $p\bar{p}$ interactions [31, 44], though FSI effects in charged particle studies are somewhat obfuscated by Coulomb effects. $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$ correlation functions are expected to be consistent with each other. In addition to FSI effects, they will also have a suppression at low relative momentum due to Fermi-Dirac quantum interference. The primary goal of this analysis is to extract the FSI scattering parameters of these pairs, as well as their femtoscopic radii.

In this thesis, we present Λ and $\bar{\Lambda}$ femtoscopy from $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collisions measured at the LHC by ALICE. In Chapter 2, we discuss the experimental details involved in measuring Λ and constructing femtoscopic correlation functions. Chapter 3 describes how we analyzed the correlation functions. This includes a discussion of the theoretical formalism used to fit the data (Section 3.1), an account of how we treat signal contaminations (Sections 3.2–3.4), a walkthrough of our fit procedure (Section 3.5), and a breakdown of our extracted scattering parameters and radii (Section 3.7).

Chapter 2

DATA MEASUREMENT

In this chapter, we discuss the experimental details involved in measuring Λ correlations. Section 2.1 covers minutiae about the data sets used in this analysis, the event selection criteria, the relevant subsystems of the ALICE detector, and the code used to perform the analysis. We describe the process of reconstructing Λ , including the cuts employed to reduce contamination from secondary or fake Λ , in Section 2.2. Section 2.3 details how correlation functions are constructed. Finally, Section 2.4 contains a study of the systematic uncertainties on the correlation functions, and plots of the correlation functions with error bars can be found in Section 2.5.

2.1 General details

2.1.1 Data selection

The $\sqrt{s_{NN}} = 2.76$ TeV Pb-Pb collision data analyzed in this analysis were recorded at the end of 2011. This dataset, LHC11h, is comprised of a number of runs. A “run” refers to the data taken during a continuous recording session of the ALICE detector, as well as to the duration of the recording session. Runs can include anything from minutes to days of data collection. Not all runs are used in this analysis. For example, a particular run might only exist for detector calibration purposes, or a subdetector essential for this analysis might have been offline during data collection. Runs were selected if they were listed with a Reconstruction Pass 2 global quality of 1 (“Good run”) in the ALICE Run Condition Table. Approximately 45 million combined central, semi-central and minimum bias events were analyzed. The following runs were used:

170593, 170572, 170388, 170387, 170315, 170313, 170312, 170311, 170309, 170308, 170306, 170270, 170269, 170268, 170230, 170228, 170207, 170204, 170203, 170193, 170163, 170159, 170155, 170091, 170089, 170088, 170085, 170084, 170083, 170081, 170040, 170027, 169965, 169923, 169859, 169858, 169855, 169846, 169838, 169837, 169835, 169591, 169590, 169588, 169587, 169586, 169557, 169555, 169554, 169553, 169550, 169515, 169512, 169506,

169504, 169498, 169475, 169420, 169419, 169418, 169417, 169415, 169411, 169238, 169167, 169160, 169156, 169148, 169145, 169144, 169138, 169099, 169094, 169091, 169045, 169044, 169040, 169035, 168992, 168988, 168826, 168777, 168514, 168512, 168511, 168467, 168464, 168460, 168458, 168362, 168361, 168342, 168341, 168325, 168322, 168311, 168310, 168115, 168108, 168107, 168105, 168076, 168069, 167988, 167987, 167985, 167920, 167915

Approximately a third of these runs were taken with the magnetic field of the ALICE detector in the positive configuration, and two-thirds had the magnetic field in the negative configuration. Having data from both configurations allows for consistency checks; one can see if physics results have any dependence on the detector's magnetic field. In this analysis, χ^2 tests of the correlation functions constructed separately among the two field configurations found the results to be consistent with each other.

Analysis was also performed on the Monte Carlo HIJING dataset, LHC12a17a fix. The approximately $4.5 * 10^5$ events in this dataset were reconstructed using calibration information from the following LHC11h runs

170593, 170572, 170388, 170387, 170315, 170313, 170312, 170311, 170309, 170308, 170306, 170270, 170269, 170268, 170230, 170228, 170207, 170204, 170203, 170193, 170163, 170159, 170155, 170091, 170089, 170088, 170085, 170084, 170083, 170081, 170040, 170027, 169965, 169923, 169859, 169858, 169855, 169846, 169838, 169837, 169835

In all cases, the z-component of the primary vertex (the position of the collision within the detector) was required to be within 10 cm of the center of the ALICE detector for an event to be selected.

2.1.2 Centrality flattening

Each event is characterized by a centrality percent. There are 100 centrality bins ranging from 0% (head-on collisions) to 99% (glancing collisions), and the centrality percent is defined such that each bin would be sampled equally using a minimum bias trigger. The LHC11h dataset employed central and semi-central triggers in addition to the minimum bias trigger, so the centrality bins are not equally filled. There are approximately as many events in the 0-10% range as there are in the 10-50% range.

While the events in the 10-50% show an even sampling, the number of events in each of the 0-10% centrality bins differ by a few percent. This analysis uses a technique called centrality flattening to ensure that each of the 0-10% bins are sampled evenly. For each magnetic field configuration, each centrality bin is assigned a weight that brings that bin even with the bin with the smallest overall content (the 9% bin). We use those weights and a random number generator to determine whether or not to throw away any given event.

2.1.3 Detector use

The data in this analysis utilizes information acquired by several of ALICE’s subdetectors. Event centrality was measured using the VZERO detector with a timing cut from the Zero Degree Calorimeter.

The neutral Λ particle is found by looking for the pion and proton that it decays into (more on this in Section 2.2). The Inner Tracking System (ITS), Time Projection Chamber (TPC), Transition Radiation Detector (TRD), and Time of Flight detector (TOF) together provided global particle tracking for these daughter particles. Information from the TPC is used to obtain PID information for V0 candidate daughters. PID information from the TOF is also used for particles with $p > 0.8$ GeV/ c , as the TPC cannot distinguish between pions, kaons, and protons above that threshold.

2.1.4 Analysis code

The code used in this research to comb through the data, reconstruct Λ , and measure correlation functions was written in ROOT, an offshoot of C++ designed for statistical analysis and data visualization. The analysis classes used for this study are checked into the ALICE Collaboration’s official code bank, AliPhysics. The code is stored in the AliPhysics/PWGCF/FEMTOSCOPY/V0LamAnalysis directory, and the classes all have the “AliAnalysisV0Lam” prefix. The code can be run on the ALICE LEGO Train (Lightweight Environment for Grid Operators) via the macro FEMTOSCOPY/macros/AddTaskLamLam.C. The analysis train can run analyses for one or more people simultaneously, thus facilitating efficient usage of resources on the ALICE supercomputing grid.

For this study, the most recent analysis of an ALICE dataset was performed on 06 May, 2016 using AliRoot vAN-20160506-1. We ran the analysis using the Correlations and Fluctuations Physics Working Group (PWGCF) train with dataset LHC11h_AOD145_fieldconfigs.

The code used for fitting the measured correlation functions (see Section 3.5.2) is not part of the ALICE code repository. However, it can be found online at <http://github.com/jsalzwedel/FemtoFitting>.

2.2 Particle selection

2.2.1 Λ reconstruction

Λ are neutral particles, and as such we do not directly detect them. Instead, Λ are topologically reconstructed by looking for the daughter tracks that they decay into. The general term for a particle reconstructed this way (including $\bar{\Lambda}$ and K_s^0) is V0.

The ALICE event reconstruction software identifies potential V0 candidates by looking for pairs of charged particles that seem like they might be the decay products of a V0; namely

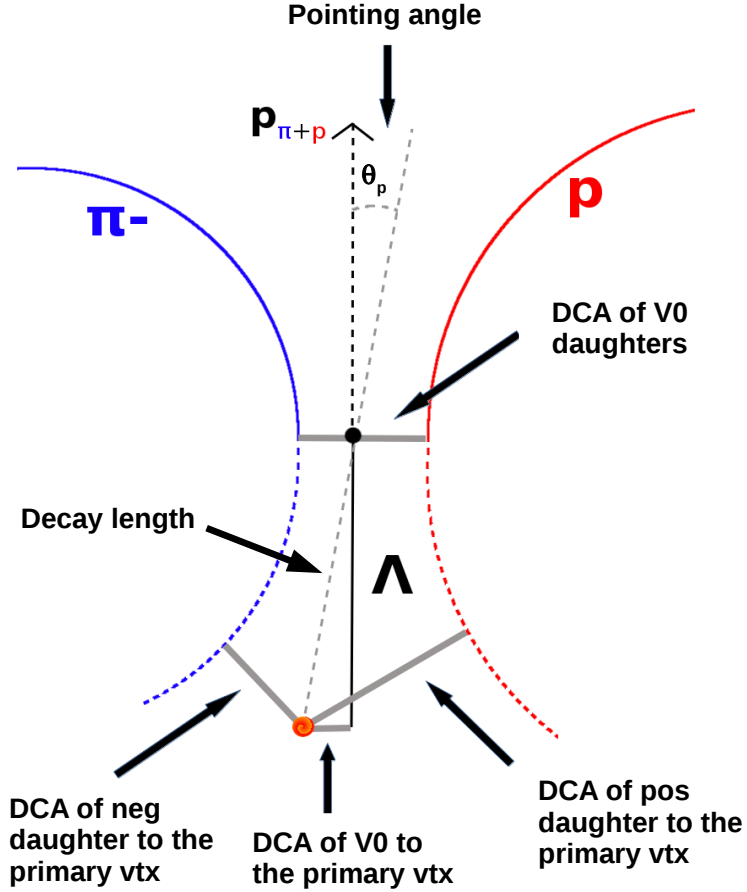


Figure 2.1: Relevant topological quantities used to reconstruct neutral V0 particles from charged tracks [53]. In addition to the reconstructed track trajectory and momentum, one needs the distance of closest approach (DCA) of each track to the primary vertex (the location of the collision), the distance of the tracks to each other, the DCA of the V0 to the primary vertex as determined by the total momentum of the daughter tracks, the decay length of the V0, and the pointing angle of the V0 (the angle between the reconstructed momentum of the V0 and its decay length vector).

that they both trace back to the same point some distance away from the primary vertex. This software, called the V0-finder, uses looser versions of the cuts described below to come up with a list of V0 candidates for each event. These V0 candidates (called AliAODv0 in the software) are constructed offline by the AliV0Vertexer using combined TPC-ITS tracking (or TPC-only if the ITS is inactive). After obtaining a list of V0 candidates, the following cuts were applied for each V0's daughter tracks. See Figure 2.1 for an illustration of these topological quantities.

- The V0 was required to have two daughters, and those daughters were required to have different charges.
- The daughter tracks were required to have hits on at least 80 pad rows of the TPC.
- The daughter protons and antiprotons were required to have $p_T > 0.5 \text{ GeV}/c$, and the daughter pions required $p_T \geq 0.16 \text{ GeV}/c$.
- Both daughters were required to be in the range $|\eta| < 0.8$.
- The TPC was used for PID purposes. A daughter was considered a proton and/or pion if the PID resolution task returned a respective $dE/dx \ N_\sigma < 3$.
- Daughter tracks were required to have a TOF $N_\sigma < 4$ if the TOF signal was available.
- The distance of closest approach (DCA) of the proton daughter track to the primary vertex was required to be $> 0.1 \text{ cm}$ (i.e. the proton daughter track is not allowed to be closer than 0.1 cm to the primary vertex at any point along its trajectory)
- The pion daughter was required to have a DCA to the primary vertex $> 0.3 \text{ cm}$.
- The two daughters were required to have a DCA to each other $< 0.4 \text{ cm}$ (the vertex software computes this DCA using momentum resolution weighting, so the units here should probably be called “effective cm”).

Each V0 is then subjected to the following cuts:

- The V0 is required to have $|\eta| < 0.8$.
- The V0 is required to have a proper decay length $L_0 = \gamma L < 60 \text{ cm}$, where L is the distance from the primary vertex to the estimated decay position as measured in the lab frame, and $\gamma = (1 - v^2/c^2)^{-1/2}$ is the relativistic Lorentz factor.
- The cosine of the pointing angle (the angle between the V0's momentum and its decay length vector) must be greater than 0.9993.
- The DCA of the V0 to the primary vertex must be less than 0.5 cm.

- A Λ was required to have a proton daughter and a π^- daughter, as determined using the above PID cuts. For a $\bar{\Lambda}$, an antiproton and π^+ daughter were required.
- For both Λ and $\bar{\Lambda}$, the reconstructed invariant mass $m_{\text{inv}}^2 = m_{\Lambda}^2 = (E_{\pi} + E_p)^2 - (\vec{p}_{\pi} + \vec{p}_p)^2$ was required to fall within $|m_{\text{inv}} - m_{\text{PDG}}| < 3.8$ MeV, where $m_{\text{PDG}} = 1.11568$ GeV is the known rest mass of the Λ as cataloged by the Particle Data Group [54]. Note that this cut value was determined from the p_{T} -integrated invariant mass peak. Future investigation of p_{T} -dependent invariant mass peak and subsequently enforcing a p_{T} -dependent cut could lead to better background reduction.

The optimal cut values were determined by comparing distributions of reconstructed Λ in Monte Carlo HIJING events [55] with simulated detector effects (the data from the MC events undergo a simulated reconstruction using the same software as for real data). The data set (LHC12a17a_fix) included injected strange signals (extra Λ , $\bar{\Lambda}$, Ξ , and Ω in each event). In this analysis, we have estimated the reconstruction efficiency of primary and secondary Λ in an effort to account for feeddown effects (see Section 3.3.3). The added particles sample a p_{T} range that allows for easy detector reconstruction, which skews those efficiency estimates. Therefore, we have performed our MC analyses with the injected signals removed. An injected primary particle (AliAODMCParticle) has AliAODMCParticle::GetLabel() greater than AliGenHijingEventHeader::NProduced()−1. Likewise, an injected secondary particle (e.g. Λ from a Σ baryon decay) has a mother particle with AliAODMCParticle::GetLabel() greater than AliGenHijingEventHeader::NProduced()−1.

Λ were reconstructed and their corresponding AliAODMCParticle objects were consulted in order to classify the Λ into different categories: Real, primary Λ ; secondary Λ coming from decays of Σ (or Σ excited states); secondary Λ coming from decays of Ξ , Ω , and other sources; fake Λ ; and K_{s}^0 that have been misidentified as Λ . For each cut parameter (e.g. cosine of pointing angle), a plot was made showing the distributions of V0 candidates of the various types. These plots contain only the parameter distributions of candidates that pass all the other cuts. Figures 2.2 and 2.3 show these distributions for the different Λ types. Figures 2.4 and 2.5 show the equivalent results for $\bar{\Lambda}$.

For the construction of Figures 2.2, 2.3, 2.4, and 2.5, each Λ type was preselected using AliAODMCParticle information, and only V0 candidates of that type were reconstructed in that analysis run. This was done for ease of histogramming the distributions. Roughly equal numbers events were analyzed for each run type, with the exception of the analysis of the primary Λ . That analysis run utilized only about half the data set. Therefore, the primary Λ ($\bar{\Lambda}$) distributions in Figures 2.2, 2.3, 2.4, and 2.5 have been scaled by a factor of 2 to compensate.

In these figures, both the shape and magnitude of each distribution is relevant for determining the cuts. For each cut type, a cut value was selected such that a looser cut

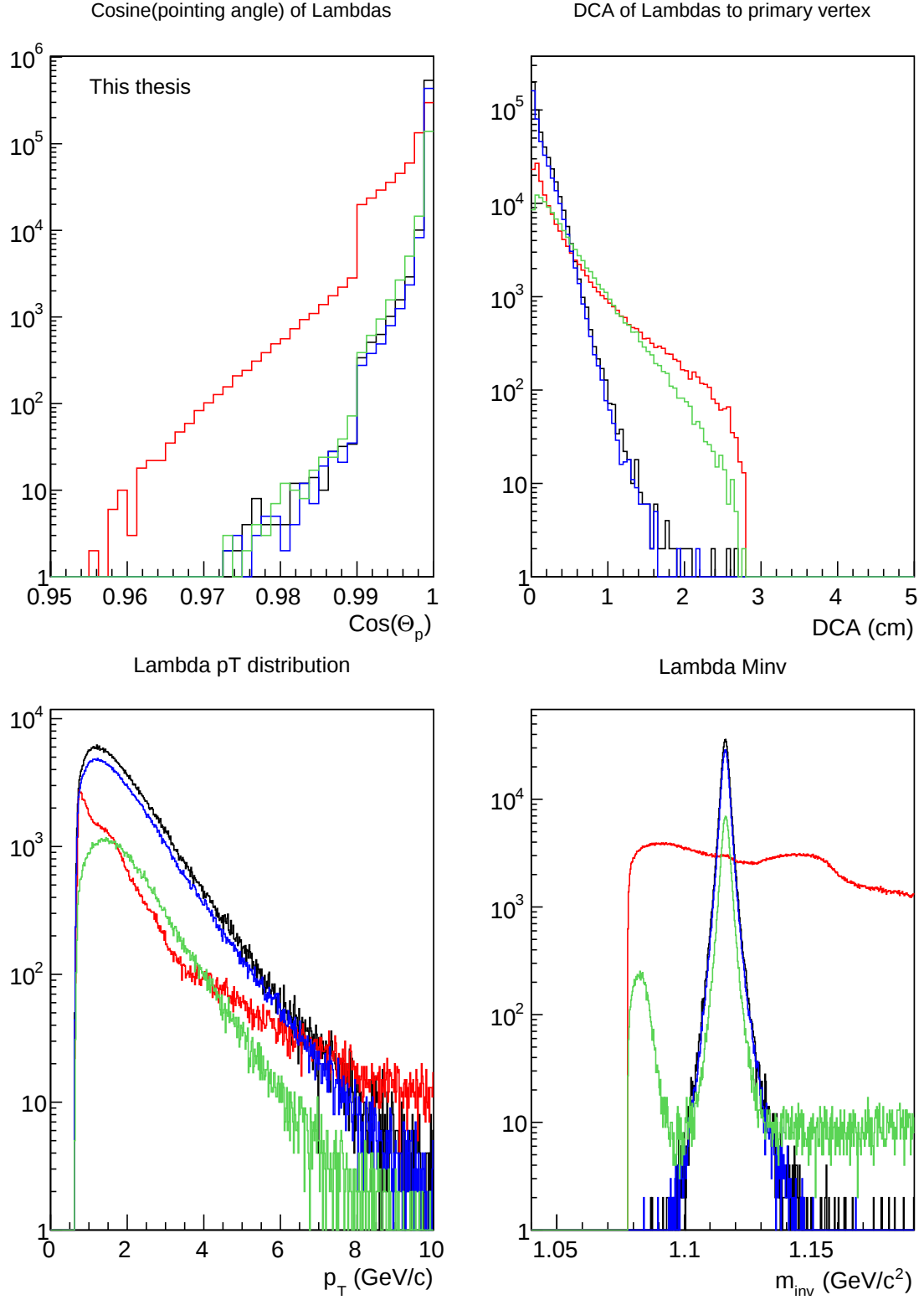


Figure 2.2: Λ cut distributions shown for real Λ (black), Λ from Σ decay (blue), secondary Λ from other sources (green), fake Λ (red). Optimal cut values were set such that a looser cut would differentially add more fake and (non- Σ) secondary Λ than primary Λ

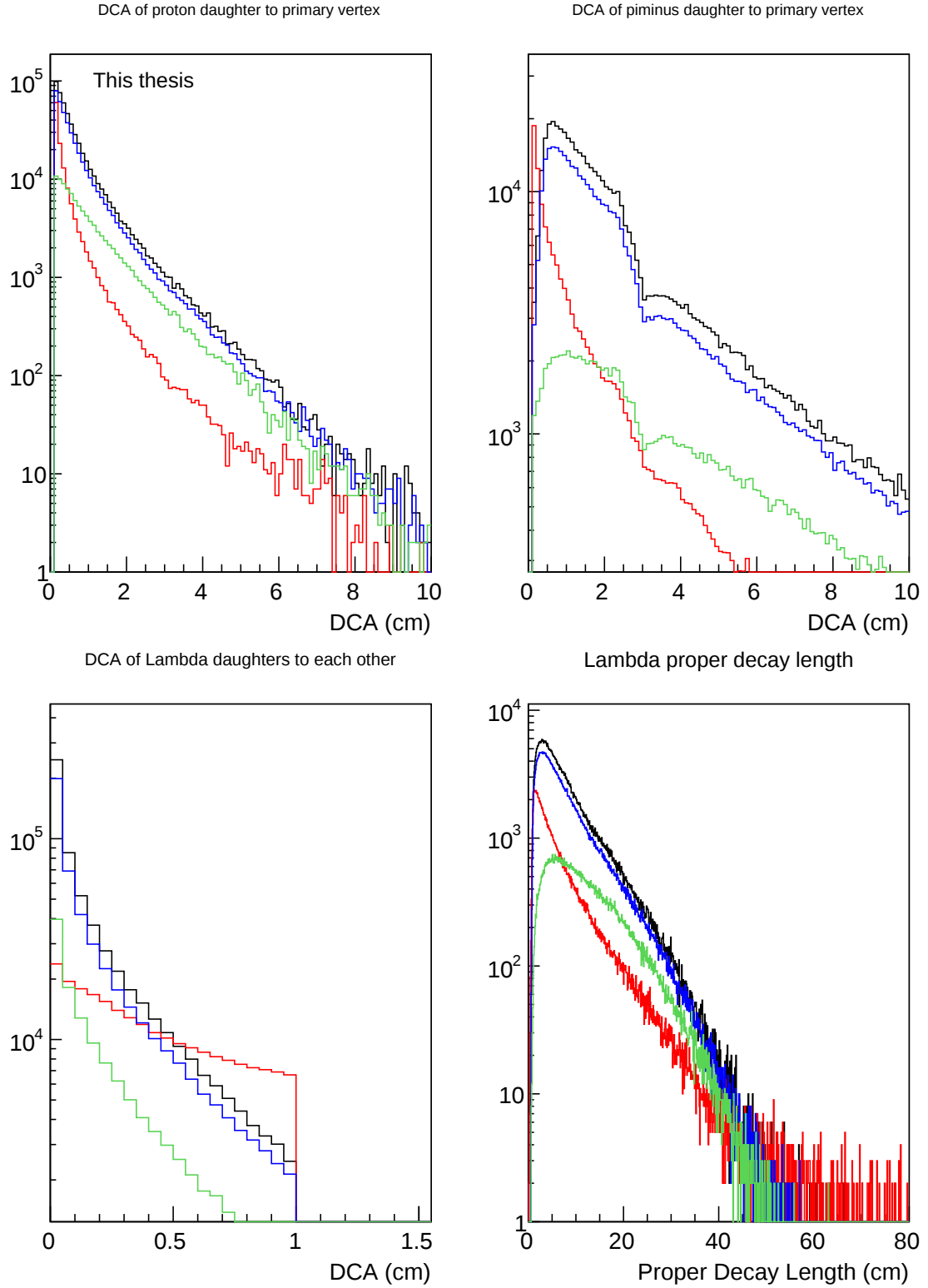


Figure 2.3: Λ cut distributions shown for real Λ (black), Λ from Σ decay (blue), secondary Λ from other sources (green), fake Λ (red). Optimal cut values were set such that a looser cut would differentially add more fake and (non- Σ) secondary Λ than primary Λ

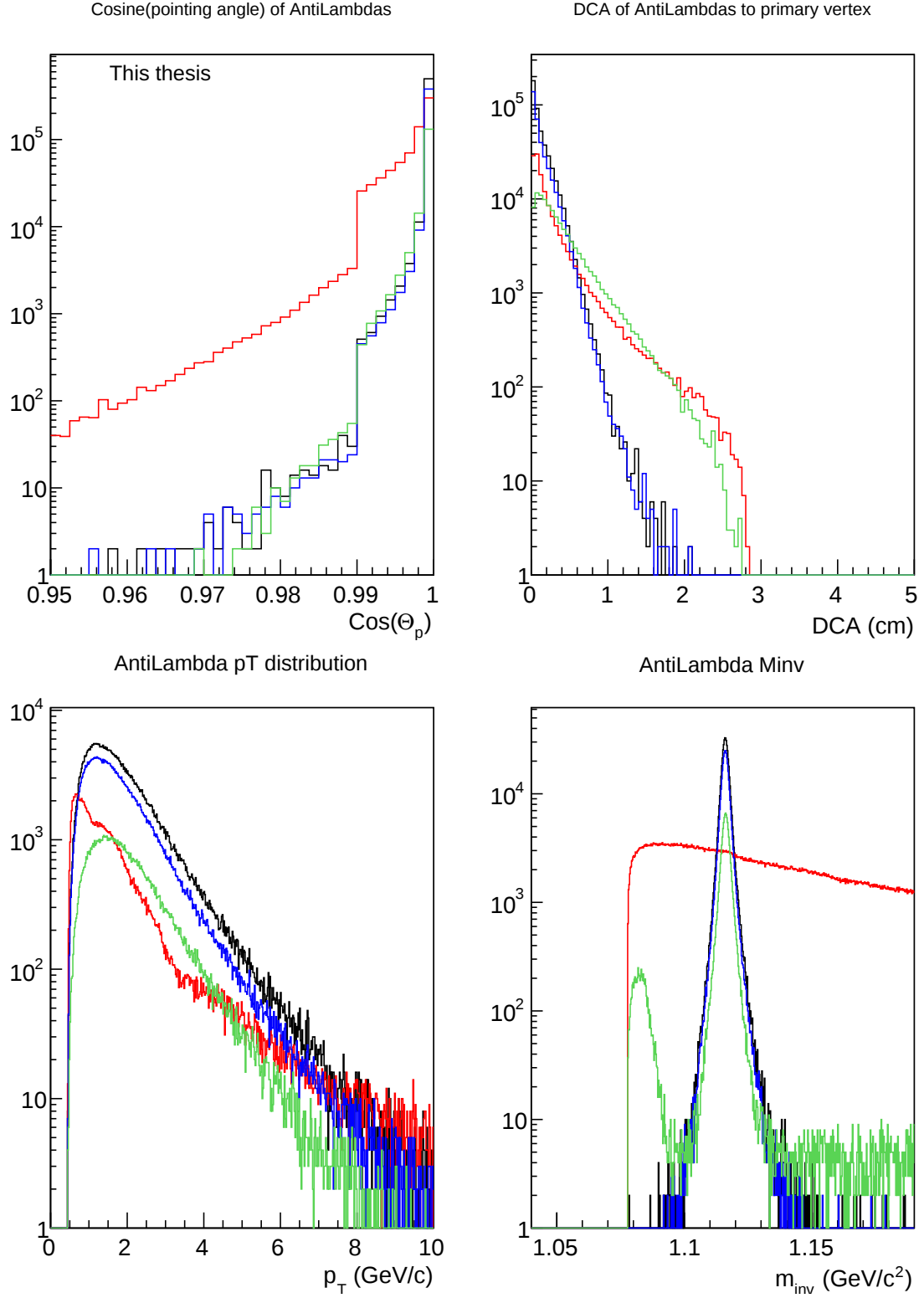


Figure 2.4: $\bar{\Lambda}$ cut distributions shown for real $\bar{\Lambda}$ (black), $\bar{\Lambda}$ from Σ decay (blue), secondary $\bar{\Lambda}$ from other sources (green), fake $\bar{\Lambda}$ (red). Optimal cut values were set such that a looser cut would differentially add more fake and ($\bar{\Lambda}$ $\rightarrow$ $\bar{\Lambda}$ $\rightarrow$ $\bar{\Lambda}$) secondary $\bar{\Lambda}$ than primary $\bar{\Lambda}$

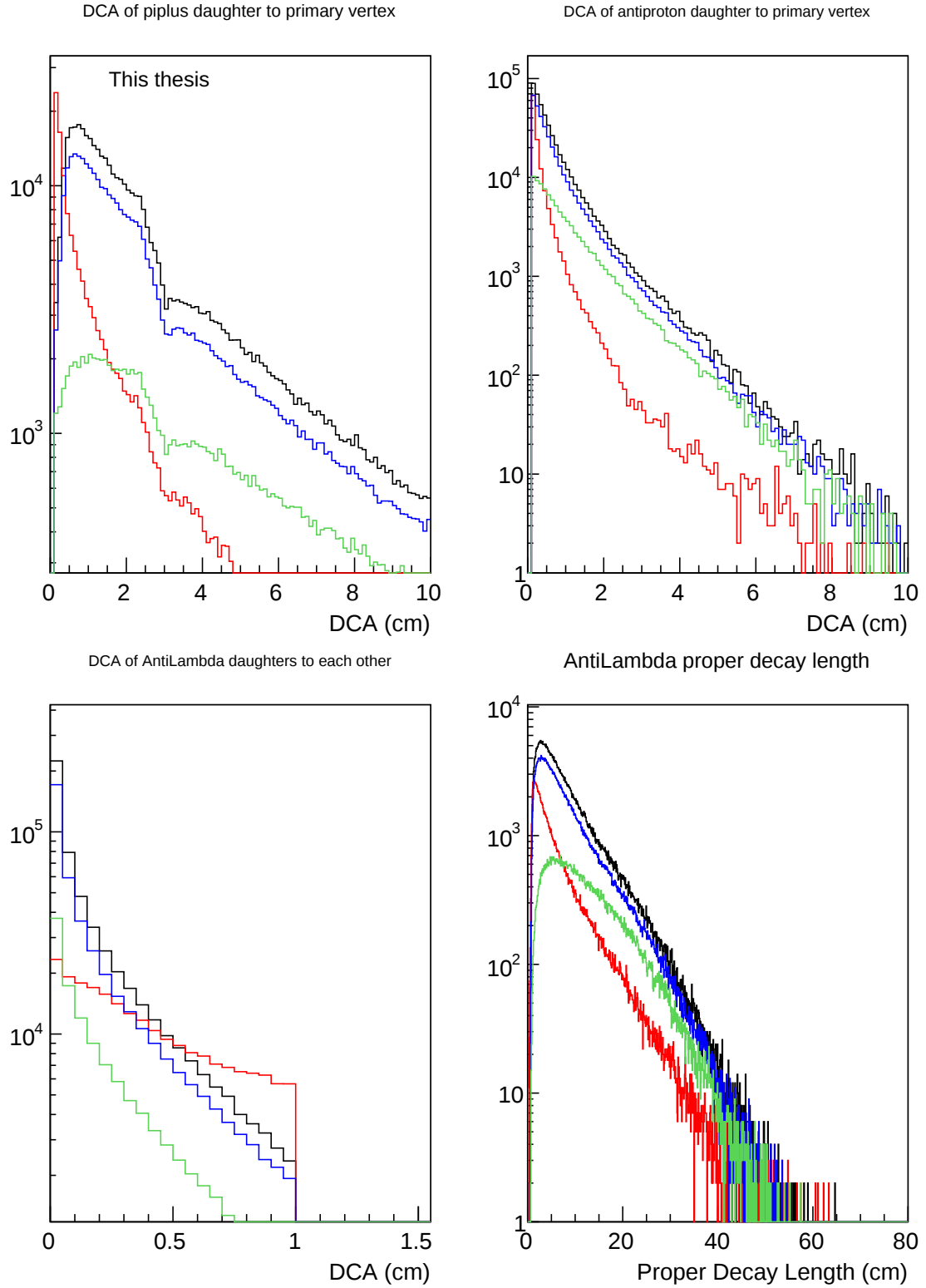


Figure 2.5: $\bar{\Lambda}$ cut distributions shown for real $\bar{\Lambda}$ (black), $\bar{\Lambda}$ from Σ decay (blue), secondary $\bar{\Lambda}$ from other sources (green), fake $\bar{\Lambda}$ (red). Optimal cut values were set such that a looser cut would differentially add more fake and (non- Σ) secondary $\bar{\Lambda}$ than primary $\bar{\Lambda}$

would differentially accept more fake or secondary Λ than primary Λ . Ideally, these cuts would be selected to reduce the inclusion of all types of secondary Λ . However, it can be seen from the distributions that Λ that come from Σ decay display virtually the same cut parameter distribution shapes as primary Λ . Only the magnitude of the Σ curves differ from the primary Λ curves. This is due to the short decay length of the Σ decay: $c\tau \approx 20$ pm for Σ^0 , and $c\tau \approx 6$ fm for $\Sigma(1385)$. As a result, secondary Λ from Σ look identical to primary Λ from the perspective of topological cuts, and they cannot be selectively removed from the analysis.

The systematic errors associated with these cut choices are discussed in Section 2.4.3. Further discussion of the reconstruction efficiency of the different Λ types can be found in Section 3.3.3.

The reconstructed invariant mass distribution for Λ and $\bar{\Lambda}$ can be seen in Figures 2.6 and 2.7. An approximation of the signal purity was estimated using a ratio of real and background (falsely reconstructed) counts. The background was estimated using a fourth order polynomial. The number of real Λ was then estimated by counting the bin content and subtracting the background. The signal quality was found to be $real/(real+background) \approx 0.95$ for both Λ and $\bar{\Lambda}$. The $\bar{\Lambda}/\Lambda$ ratio is estimated to be about 93%.

2.2.2 Shared daughter cut

Occasionally two or more V0s are reconstructed that both claim the same daughter track. However, a given track can only be the daughter of, at most, one V0. Therefore, if several reconstructed V0s all lay claim to the same daughter, at most only one of them can be a true V0. It was decided to employ a cut that would compare characteristics of the different V0s to determine which of them was most likely to be the true parent particle. After the best V0 was found, the competing V0s were removed from the list of V0 particles and not used during same- or mixed-event pair construction.

A study was done using MC event truths to determine which cut would be most successful at removing fraudulent V0s. The following characteristics were examined as candidates for the cut criterion:

- Closeness of invariant mass to the PDG mass value.
- DCA of the daughter particles to each other.
- Cosine of the pointing angle.
- DCA of the V0 to the primary vertex.

The process of culling V0s with shared-daughters was done by looping through V0s on an event-by-event basis after the V0 reconstruction process was completed. For each V0

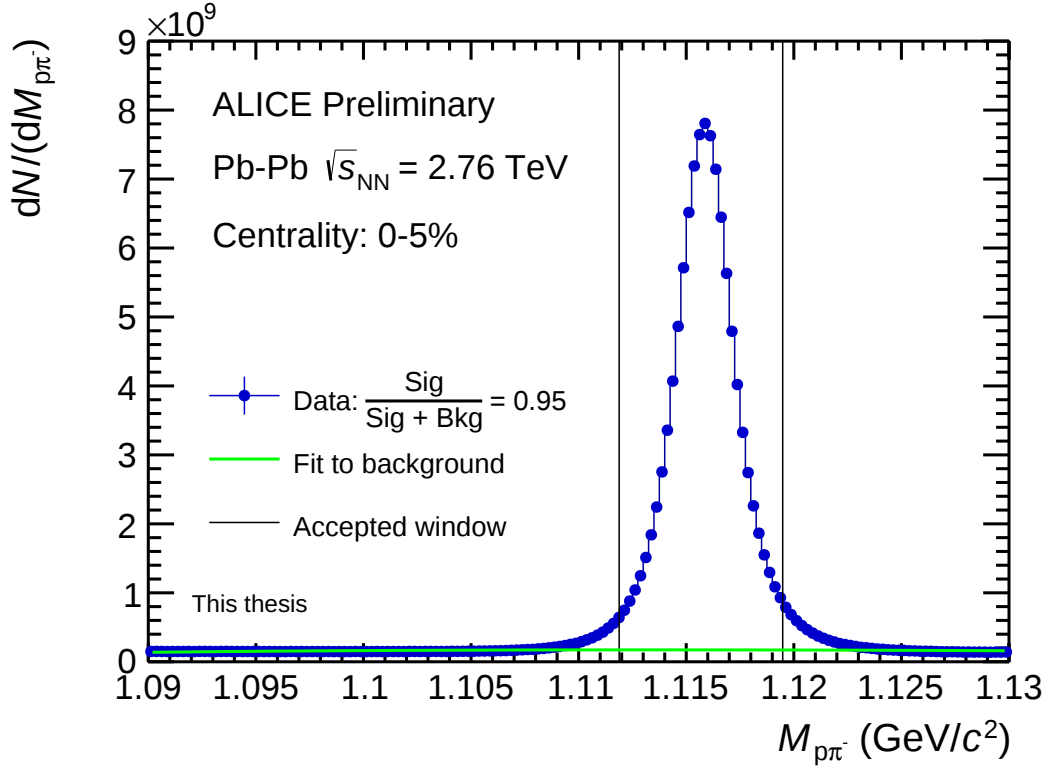


Figure 2.6: Invariant mass distribution for reconstructed Λ using the optimal analysis cuts. The plots show V0s reconstructed from centrality integrated LHC11h data. The green line shows a fourth order polynomial fit to the background, which is used to estimate the number of real and fake Λ . The estimated ratio of real Λ to all reconstructed Λ in the signal region ($|m_{\text{inv}} - m_{\text{PDG}}| < 3.8$ MeV/c²) is approximately 0.95.

in the list of successfully reconstructed V0s (i.e. V0s that passed all cuts), the V0 was compared with each other V0 in the event to determine if they shared a daughter. Two V0s “share” a daughter if both V0s have a charged daughter track with the same track ID. When two V0s were found that shared a daughter, their values of the cut criterion were compared. The daughter with the stricter value (e.g. smaller DCA to primary vertex) was adjudged to be the better V0 candidate, and the other V0 was flagged as “bad” for failing the shared-daughter cut. V0s that failed the shared-daughter cut were not used in any pair construction (same event or mixed event).

However, it is possible for a V0 to share a daughter or daughters with two different V0s. In the aforementioned scheme, it is therefore possible for an excessive number of V0s to fail the shared daughter cut (excessive in that some of those failed V0s only shared a

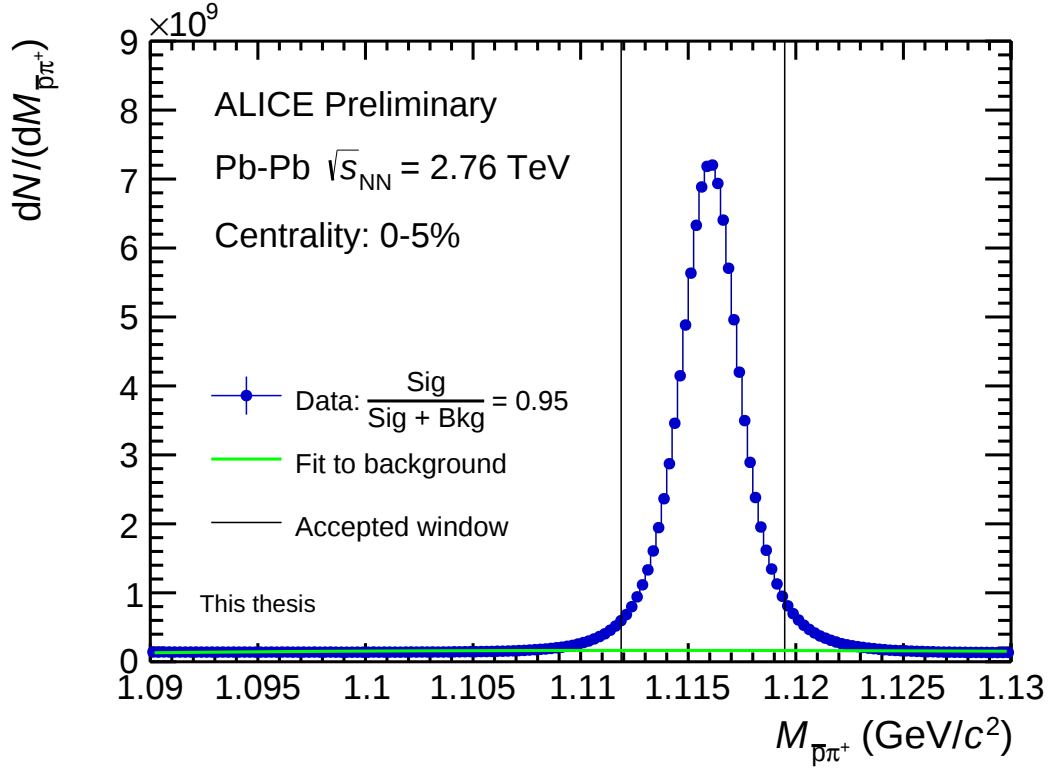


Figure 2.7: Invariant mass distribution for reconstructed $\bar{\Lambda}$ using the optimal analysis cuts. The plots show V0s reconstructed from centrality integrated LHC11h data. The green line shows a fourth order polynomial fit to the background, which is used to estimate the number of real and fake $\bar{\Lambda}$. The estimated ratio of real $\bar{\Lambda}$ to all reconstructed $\bar{\Lambda}$ in the signal region ($|m_{\text{inv}} - m_{\text{PDG}}| < 3.8 \text{ MeV}/c^2$) is approximately 0.95.

daughter with another failed V0). To avoid throwing away too many V0s, a second loop was added to shared-daughter cut method to un-flag V0s that no longer share a daughter with any “good” V0. Example: V0s “A” and “B” both share a daughter, and V0s “A” and “C” share a different daughter. “A” has a better DCA than “B”, so “B” is flagged as bad. Then “C” is found to have a better DCA than “A”, so “A” is flagged as bad. Finally, “B” is re-flagged as good because it no longer shares a daughter with any other good V0.

This method was adapted to run iteratively over the list of V0s candidates until a stable list of non-sharing V0s was found. Then the MC truths of each V0 (those marked “good” and those marked “bad”) were examined. With this information, it is possible to calculate the percentages of true V0s and fake V0s that are removed via this process. This analysis was run several times, each time using a different comparison criterion from the list above.

Based on the MC truth study, keeping the V0 with the closest DCA to the primary vertex was found to successfully keep 87% of true V0s that had shared daughters. In comparison, the cosine of pointing angle criteria was also successful 87% of the time ($\cos(\Theta_p)$ and DCA to primary vertex are strongly correlated), the DCA of daughters to each other criteria was successful 81% of the time, and the absolute value of mass difference criteria was successful 79% of the time. Based on these results, the DCA to primary vertex criteria was selected to be used for the shared-daughter cut before correlation function construction.

One of the benefits of this cut is that it helps remove a splitting-like effect (see Section 2.3.3 for more on splitting), wherein a true V0 shares each of its daughters with two separate fake V0s. By the nature of the reconstruction cuts, all three of the V0s will be close in momentum space. Therefore a pair constructed from the two fake V0s will contaminate the low relative-momentum region of the correlation function. That pair of fake V0s does not share any daughters with each other, so it is not kept out of the correlation function by simple daughter ID checks. But the iterative daughter-sharing cut described above is capable of removing this fraudulent pair, since either of the fake V0s is likely to be removed for sharing a daughter with the true V0.

2.3 Correlation function construction

2.3.1 Relative momentum

This analysis studies two-particle correlations as a function of the one-dimensional relative momentum k^* . To be precise, for a pair of particles with four-momenta p_1 and p_2 , k^* is the momentum of one of the particles in the pair's center of mass frame. It is defined as

$$k^* = \sqrt{-k^\mu k_\mu}, \quad (2.1)$$

where k^μ is half of the four-vector relative momentum

$$k^\mu = \frac{p_1^\mu - p_2^\mu}{2} - \frac{(p_1 - p_2) \cdot P}{2P^2} P^\mu, \quad (2.2)$$

P is the four-vector momentum sum.

The literature is inconsistent when it comes to the notation for relative momentum. In some cases, k^* refers both to the invariant quantity and to the four-vector. In this thesis, we will try to consistently use k^* to designate the invariant quantity, and use k^μ or simply k when speaking of four-momenta. In some papers, q or q_{inv} is used to represent the relative momentum. Though we won't use q in this analysis, for reference $q^\mu = 2k^\mu$, and $q_{\text{inv}} = 2k^*$.

2.3.2 Construction

The correlation function is constructed as

$$C(k^*) = \frac{A(k^*)}{B(k^*)}, \quad (2.3)$$

where $A(k^*)$ is the two-particle distribution of the given event and $B(k^*)$ is the distribution of a reference event. In practice, the signal event histogram $A(k^*)$ is constructed by binning the k^* value of each pair of particles in a single event, and repeating this for each event. The reference or background histogram is constructed by binning k^* for pairs of particles taken from different events. The correlation function therefore is the ratio of the two-particle phase-space with correlated effects coming from interactions (the numerator) to the two-particle phase-space without correlated effects (the denominator). Because the denominator pairs are uncorrelated, their two-particle phase space is essentially a convolution of two single-particle phase spaces. Of course, the numerator pairs are also affected by the single-particle phase spaces. We assume that the physics of the numerator pairs can be factorized to separate out the correlated physics effects and the single-particle phase space effects. Thus, dividing by the background pair distribution leaves us with just the correlated physics effects, i.e. the correlation function. See Section 3.1.1 for more details.

In order for the phase-space effects to divide out, it is of course necessary for the mixed events to resemble the true events. For example, a mixed event where half the particles come from a central collision and half the particles come from a very peripheral collision would have a very different two-particle phase space than a central collision would. Care is therefore taken to ensure that pairs are mixed from events with similar characteristics. Mixed events are required to have approximately the same centrality (bin width of 5%) and primary vertex z-position (2 cm bin width). There are approximately five mixed events for each real event. When constructing pairs of V0s from the same event, the V0s are not allowed to share a daughter with the same track ID. The correlation functions are normalized to unity in the $0.3 < k^* < 0.5$ GeV/c range.

2.3.3 Pair cuts

Femtoscopic studies look at the relative momentum of particles, and often the most interesting physics lies at very low relative momentum (around 0.1 GeV/c and lower). As a result, two-track reconstruction effects such as track splitting and track merging, both of which occur for tracks with similar momenta and trajectories, can have a large effect on the measured results. Track splitting means that the track left by a single charged particle is reconstructed as two separate tracks. A track splitting-like effect can also occur from cutting out tracks with too many shared TPC clusters. Track merging occurs when the tracks of two separate particles are reconstructed as a single track. For V0s, splitting/merging

occurs on the level of the daughter tracks, and it affects the reconstruction of the parent V0s. However, one must be somewhat cautious here, since there can be physics effects that lead to a suppression or enhancement of tracks that are close together. In particular, small track separation correlates somewhat with small relative momentum, so effects that suppress (enhance) low- k^* pairs may also lead to some suppression (enhancement) of low-average separation pairs.

In this analysis, two-track reconstruction effects are combated via a cut on the average separation of daughter tracks from different $\Lambda/\bar{\Lambda}$. The average separation distance of two daughter tracks are computed at nine different radii of the TPC. The track positions are measured via `AliExternalTrackParam::GetXYZatR()` at the following radii: 85, 105, 125, 145, 165, 185, 205, 225, and 245 cm. Pairs of daughter tracks are made from one daughter from the first V0 and one daughter from the second V0. Pairs are made for all daughter combinations, and each type of pair is tracked separately. For each pair of daughter tracks, their separation is computed at each of those radii, and their average separation is taken as the mean value. If a track does not have a value at a particular radius (e.g. a low p_T track with a small radius of curvature), that position is excluded from the calculation of the mean. Same- and mixed-event pairs are binned according to this separation distance. To obtain an estimate of the merging/splitting effects both distributions are then scaled by the number of pairs at high average separation distance (10+ cm), and a correlation function (same event pairs/ mixed event pairs) is created.

To ensure that average separation distributions of mixed-event pairs are comparable to the same-event distributions, it is necessary to perform a shifting of the primary vertex for mixed events. Doing so allows the track separations to be calculated as though both events had the same primary vertex. Without this correction, the mixed-event distributions are biased by differences in the primary vertex location. One could imagine an extreme example of this if one used a 20 cm wide z-vertex bin. In that case, tracks from different events could be shifted relative to each other by as much as 20 cm.

The average separation correlation functions for various V0 daughter combinations are available in Figure 2.8. We see that all like-sign pair combinations, as well as proton-antiproton, show a dip at low average separation. This dip tells us that there is a depletion of tracks that are close together. This depletion may come from track merging or from effects similar to track merging, such as the rejection of tracks that have too many shared TPC clusters. For these pairs of tracks, we implement cuts at average separation value where the correlation function looks to reach unity. These cuts are applied uniformly to same-event and mixed-event pairs. pp and $\bar{p}\bar{p}$ pairs are cut if their separation lies below 12 cm, while $\pi^-\pi^-$, $\pi^+\pi^+$, $p\pi^+$, $\bar{p}\pi^-$ and $p\bar{p}$ are cut below 10 cm.

For the other opposite-sign pairs, the splitting/merging picture is less clear, in great part due to the low statistics at low average separation. First and foremost, this tells us

that there is a very low phase space for finding oppositely charged particles. That should not be a surprise, as the ambient magnetic field accelerates them in opposite directions. As it is not clear whether there exists splitting or merging from these pairs, a conservatively large average separation cut has been implemented. $p\pi^-$ and $\bar{p}\pi^+$ are cut at 15 cm, and $\pi^-\pi^+$ are cut at 25 cm.

One alternative/complement to using an average separation cut would be to enforce a relative decay length cut, where V0s are not paired with each other if their difference between their lab frame decay length values is less than some cut value (e.g. a few cm). One advantage of this cut is that the decay lengths of the various particles should be independent of physics effects - i.e. independent of any quantum interference or final state interaction that might occur between the particles. Meanwhile, the decay length cut may minimize daughter splitting effects, as well as some merging of low- and mid- p_T daughters. High- p_T daughters may still be affected by merging, since higher- p_T particles have straighter trajectories in the TPC. A study of relative decay length has not been performed here, but it may prove useful to combat splitting/merging effects in future V0 analyses.

2.3.4 Momentum resolution correction

There is an inherent resolution to V0 momenta that depends on the resolution of the reconstructed daughter tracks. Because of this limit of precision, there is also an associated resolution to the relative momentum of two V0s. This causes the reconstructed relative momentum of the pair to differ from the true momentum of the pair, and it therefore introduces subtle changes to the measured correlation functions which must be accounted for.

Traditionally, the momentum resolution correction is performed on the experimental data to unsmear it into a form that should resemble the theoretical predictions (i.e. fits). In this analysis we have instead chosen to smear the fits to conform to the experimental data. The experimental correlation functions therefore remain uncorrected. As our corrections are employed at the time of fitting, we'll discuss the full methodology in section 3.4.

2.4 Systematic uncertainties of correlation functions

We will now attempt to identify the various systematic effects that may affect the correlation functions. The goal will be to quantify the systematic uncertainties associated with each k^* bin. It should be noted, however, that subsequent fits to the correlation functions are not directly impacted by these error bars. Instead, the systematics on the resulting fit parameters will be computed by refitting many versions of the correlation functions (e.g. those built with different reconstruction cuts) and comparing the results. This will be discussed in greater detail in Section 3.6.1.

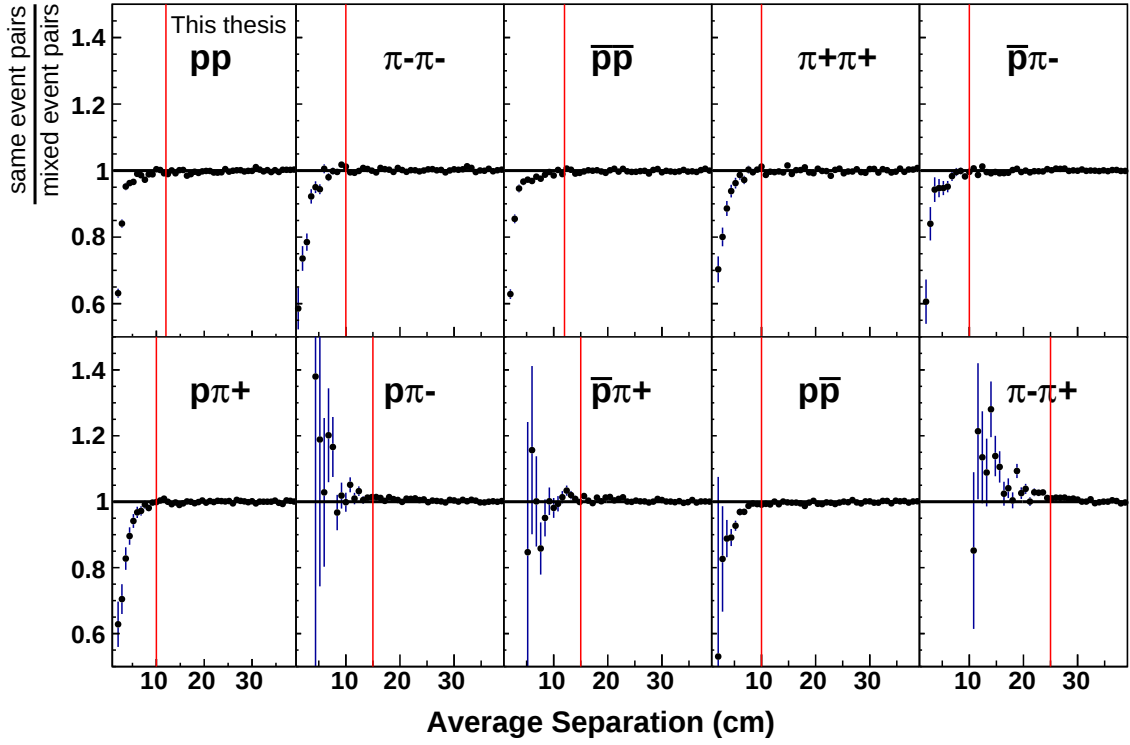


Figure 2.8: Correlation functions of average TPC track separation for pairs of daughters from different V0s. A correlation function is shown for each type of daughter combination. The positions of tracks are computed throughout the TPC, and an average separation value is calculated from those positions. This process is done for same- and mixed-event pairs. The correlation function reveals merging-like effects for all same-sign pairs and also for $p\bar{p}$. To remove splitting/merging effects, V0 pairs are cut if their daughters' average separation is below the indicated red line. Conservatively large cuts have been implemented for oppositely-charged pairs.

2.4.1 Consistency checks for uncorrelated histograms

One of the tools we employ in the search for systematic effects is the `TH1::Chi2Test(TH1* h1)` method, which compares two histograms bin-by-bin to determine a χ^2 . The test is performed assuming the hypothesis that the two histograms are Poissonian samples of the same underlying distribution. A p-value is computed which represents the probability that a measurement could be performed that would yield results this different or more, given the hypothesis of identity. The hypothesis of identity is rejected in instances where the p-value is lower than 0.01. At that threshold, less than one percent of tests will fail of data sets sampled from the same underlying distribution. This corresponds to rejecting the hypothesis of identity for data samples that differ by roughly 3σ . The 0.01 metric is

somewhat arbitrary - it was chosen to be small enough that pure statistical fluctuations should not often lead to false or unnecessary systematic errors. For example, if a p-value of 0.1 were used instead, this would result in roughly 10% of “good” results failing the test and contributing to the uncertainty, which would lead to a gross overestimation of the total systematic error. Similarly, a threshold of 0.001 would probably underestimate the error. Thus, 0.01 was chosen.

In this analysis, `Chi2Test()` will be used to compare two correlation functions which are uncorrelated with each other (uncorrelated in the sense that they are completed independent data samples). One example where this can be employed is the comparison of correlation functions constructed using data taken under different field configurations. This analysis needs all the data available to it, so results taken under different field configurations will need to be merged to improve statistics. But before they are combined, it must be checked that they contain compatible results. The `Chi2Test()` method provides a means of characterizing the degree of similarity of the correlation functions.

Correlation functions that surpass the significance level of 0.01 will be treated as being functionally the same within statistics, and their weighted average will be computed without recourse to systematic uncertainty. In the instances where they do not exceed the significance level, they will be analyzed on a case-by-case basis. In particular, we will look to see if there are any systematic differences between the correlation functions. For example, it may be seen that the lowest ten k^* bins will all be higher in one correlation function than another. If so, the data will still be merged (so long as the results aren’t dramatically different), but a systematic uncertainty will be computed and applied to the final averaged correlation function in the form of systematic error bars.

When using `Chi2Test()` to compare correlation functions, the trouble areas that result in small p-values will likely be the lowest k^* bins (which exhibit most of the interesting physics, and which also suffer the most from two-track effects), and large k^* bins (close to 1GeV/c). The height of the the large k^* bins is noticeably susceptible to statistical (or systematic) deviations in the normalization region. The high k^* bins have relatively small statistical error bars, so small deviations of normalization between correlation functions can result in large χ^2 values. The correlation functions are all normalized to unity in the $0.3 < k^* < 0.5\text{GeV}/c$ range, so there are unlikely to be significant deviations seen in that range.

It is important to emphasize that `Chi2Test()` is only appropriate for comparing two uncorrelated histograms. In the case of correlated histograms (e.g. correlation functions constructed using slightly different two-track cuts or different reconstruction cuts), a different method must be used to test the histograms for consistency.

In this analysis, χ^2 tests of correlation functions with different magnetic field configurations found that the results from the two fields were consistent with each other. Therefore,

the magnetic field configurations do not contribute to the systematic uncertainties on the correlation functions or fit results.

2.4.2 Consistency checks for correlated histograms

Chi2Test() cannot be used to compare correlated data because the test is performed assuming the statistical error bars on the two histograms are independent. That is obviously not the case for two histograms that differ in their content by only a few percent. In this analysis, we follow the advice of Roger Barlow [56] and perform consistency checks of correlated data in the following fashion:

1. Take the difference between the two correlation function histograms

$$\Delta C(k^*) = C_1(k^*) - C_2(k^*),$$

using $C_1 \rightarrow \text{Add}(C_2, -1)$. The resulting histogram shows the usually small differences between the two data sets.

2. Here, ROOT adds the errors in quadrature, which isn't correct for correlated data. To fix this, we manually set the error bars using $\sigma_\Delta(k^*) = \sqrt{|\sigma_1^2(k^*) - \sigma_2^2(k^*)|}$. This error is correct in the specific case (found here) that one histogram is entirely a subset of the other.
3. We now look to see if there are any significant discrepancies or systematic differences between the two original correlation functions. If there is a structural difference, it will generally be largest in the lowest k^* bins, dropping off as k^* increases. This shape is reasonably described by a decaying exponential: $f(k^*) = a \exp(-bk^*)$. The amplitude a tells us the size of the discrepancy. The width b has no significance; it just helps the fit.
4. We fit the histogram with the exponential, and we extract the amplitude a and its error σ_a . For this analysis, we conclude that there is no discrepancy if a is within 2σ of zero. If a is more than 2σ from zero, then we have a source of systematic uncertainty.
5. In the cases where there is a systematic deviation, we take the height of the fit function in the center of each k^* bin as the systematic uncertainty for that bin from this cut. As the fit function can be positive or negative, this leads to asymmetric errors.

Using the steps listed above, we can evaluate if seemingly small changes to the cuts employed in correlation function construction change the resulting data in any significant way.

2.4.3 Systematic errors from reconstruction cuts

In this section we discuss the systematic errors associated with the V0 reconstruction cuts (e.g. cosine of pointing angle, DCA to primary vertex, etc.). Loosely speaking, these systematic checks test the sensitivity of the correlation function to small changes in the Λ purity. The optimal cut values were discussed in Section 2.2.1. For example, the optimal cut value chosen for the DCA of daughter tracks to each other was 4 mm (daughters tracks were required to pass within that distance of each other). However, there is some ambiguity in determining the optimal cut value. For example, any value within $4 \pm 10\%$ would be reasonable. The correlation function should ideally be insensitive to small changes in the cut value. If a small tweak to the cut values produces a statistically significant change in the correlation function, that is a source of systematic uncertainty.

To test for systematic differences, we make correlation functions using Λ ($\bar{\Lambda}$) reconstructed with these different cut values. At any given time, one cut type (e.g. DCA of proton daughter to primary vertex) will be varied away from the default cut value by $\pm 10\%$, while the other cut types (e.g. cosine of pointing angle) are fixed to their default values.

Data has been collected for variations of the following cuts:

- DCA of proton daughter to primary vertex
- DCA of pion daughter to primary vertex
- DCA of daughters to each other
- Proper decay length of Λ
- Eta of Λ
- Cosine of Λ pointing angle
- DCA of the Λ to primary vertex
- p_T of Λ
- Reconstructed mass ($m_{\pi p}$)

In each case, correlation functions have been constructed for each pair type ($\Lambda\Lambda$, $\bar{\Lambda}\bar{\Lambda}$, and $\Lambda\bar{\Lambda}$), and for the three cut values (optimal cut, cut + 10%, cut - 10%). The optimal cut correlation functions are compared with the tighter and looser cut correlation functions. The correlation functions all share much of the same data, so they are highly correlated with each other. We therefore analyze the changes in the correlation function using the method described in Section 2.4.2.

Checking each cut yields a positive, negative, or non-existent systematic uncertainty from that cut for each k^* bin. For a given bin, we take the positive and negative values with the largest magnitude to be the upper and lower asymmetric errors, respectively.

2.4.4 Systematic errors from pair-wise cuts

The systematic uncertainties from the average separation cut are implemented in exactly the same way as the single-particle reconstruction cuts in the previous section. For each cut value, we vary the cut by $\pm 10\%$, construct correlation functions, and find the uncertainty for each k^* bin via the method in Section 2.4.2. In general, the systematic uncertainties are quite small — about an order of magnitude smaller than the uncertainties from the reconstruction cuts.

2.4.5 Systematic errors from shared-daughter cut

At this time, two femtoscopic analyses at ALICE employ a shared-daughter cut (this analysis and the K_S^0 analysis), though each cut is implemented with independent code. The reasons for this cut are intuitive - two V0s cannot actually have the same daughter, regardless of what might arise from the reconstruction process. Moreover, the efficacy of the cut is high: the MC studies of this cut have shown that only 13% real V0s with shared daughters are cut. The removal of a V0 (real or fake) removes all of the same-event and mixed-event pairs that would have included that V0. In most cases, this amounts to an improved purity of the correlation function, since there are fewer pairs of fake and therefore uncorrelated particles. There also a removal of a splitting-like effect - in this case the “split” is an extra V0 close in momentum space to the first. While the two “split” V0s will never be paired together, both are paired with other V0s, resulting in extra pair counts in roughly the same k^* bins. Those extra pairs are removed by virtue of the shared daughter cut.

While the benefits of the shared-daughter cut are apparent, and the disadvantages seem to be minimal (even when you cut a real V0, you are left with a fake V0 that is close in phase space), it is unclear at this time how to determine what systematic uncertainties may arise from this cut and how to quantify that uncertainty.

2.4.6 Combining different sources of systematic error

Systematic errors have been evaluated for reconstruction cuts and average separation cuts. The reconstruction cuts and the separation cuts both yield asymmetric errors for each k^* bin of the correlation function. We then add those asymmetric errors in quadrature (positive with positive, negative with negative) to get the full systematic uncertainty on the correlation functions. The resulting correlation functions with combined systematic

errors are shown in Section 2.5. For all correlation functions and all k^* bins, the systematic uncertainties are much smaller than the statistical error bars.

2.5 Correlation function results

2.5.1 Combining correlation functions

Correlation functions have been constructed using the entirety of the LHC11h data set. As discussed in Section 2.3.2, events were binned in 5% centrality bins. When combining data from different centrality bins, care was taken to combine the correlations by taking weighted averages of the correlation functions.

The naive method of combining correlation functions involves adding numerators from different bins and adding denominators from different bins and then taking the ratio. However, this has the pitfall of potentially combining drastically different phase spaces. As a result, it may introduce artificial signals into the combined correlation function. Instead of using this method, a more correct method is to construct correlation functions for different bins separately. Then one takes a weighted average of the correlation functions, using the number of numerator pairs in each correlation function as the weight:

$$C_{\text{combined}}(k^*) = (\sum_i w_i C_i(k^*)) / (\sum_i w_i) \quad (2.4)$$

where the sum is done over different correlation function bins (e.g. 0-5% and 5-10%) and w_i is the number of numerator pairs in C_i . This averaging is done on a bin-by-bin basis in k^* . Note that this averaging procedure should be used in general to combine data sets with potentially different phase spaces. For example, this procedure is employed to combine data from runs utilizing the two different field orientations (see Section 2.1.1). It is also employed when combining $\Lambda\Lambda$ data with $\bar{\Lambda}\bar{\Lambda}$ data. The statistics are too limited to analyze these correlation functions separately. Instead, we combine them together using the above method for better fitting.

For each centrality range (0-10%, 10-30%, and 30-50%), the correlation functions of the two pair types were found to be consistent with each other (p-values of 0.26, 0.02, and 0.43 respectively). Therefore, no additional systematic error was assessed to come out of their merging.

Figures 2.9, 2.10, and 2.11 show combined $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ correlation functions for the 0 – 10%, 10 – 30%, and 30 – 50% centrality ranges respectively. Figures 2.12, 2.13, and 2.14 show $\Lambda\bar{\Lambda}$ correlation functions for the 0 – 10%, 10 – 30%, and 30 – 50% centrality ranges respectively. Each plot includes the combined systematic errors discussed in Section 2.4. All six correlation functions are included in Figure 2.15.

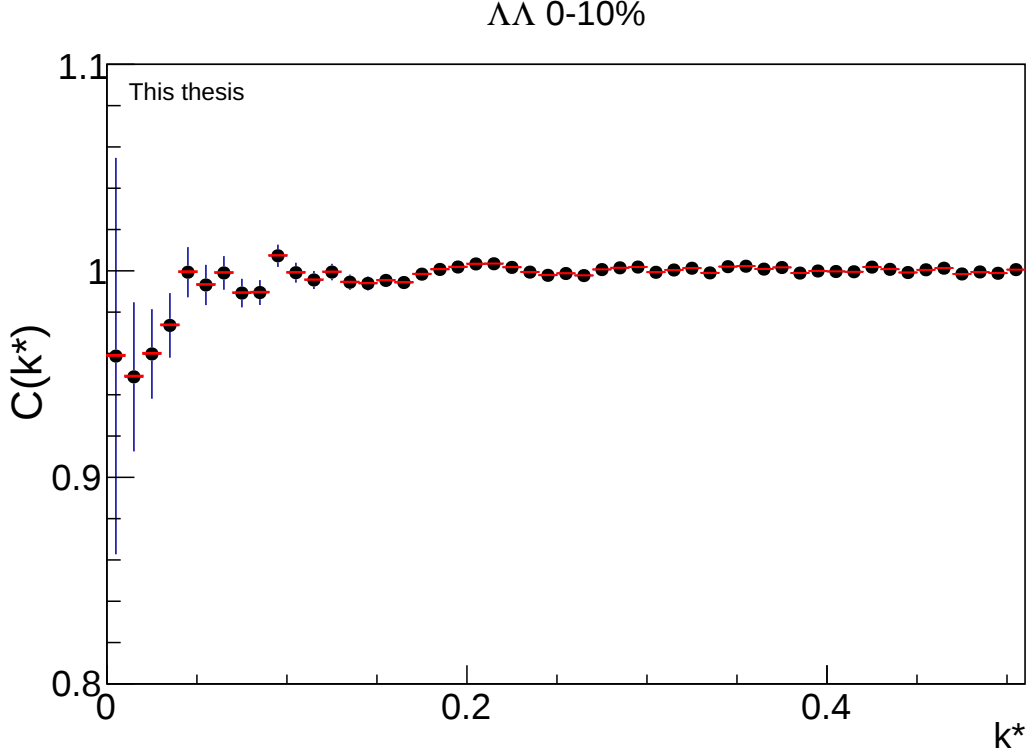


Figure 2.9: $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ correlation function for the 0-10% centrality range with statistical and systematic errors. A dip at low k^* is seen which is expected from quantum interference.

2.5.2 Interpretation of correlation function results

Generally speaking, the interesting part of a relative-momentum correlation function is the region below 100 or 200 MeV/c. Non-unity structures/backgrounds may be visible at large relative momentum - these structures can complicate the fitting procedure, and sometimes even direct attention to flaws in the analysis. They can also exist in the low-relative momentum region, though there they are generally assumed (or at least hoped) to be small. However, the information about the source size and two-particle interactions is predominantly encoded within the low- k^* region. For the time being, this analysis will restrict its attention to behavior in that region. The non-femtoscopic background will be discussed more in Section 3.5.4.

The $\Lambda\bar{\Lambda}$ correlation functions (Figures 2.12, 2.13, and 2.14) exhibit a clear, centrality dependent suppression in the low- k^* region, even with the significant statistical error bars for the 10-30% and 30-50% systems. A signal is also visible in the $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ data, though it is smaller in magnitude and localized to a narrower k^* region.

Several factors may contribute to the small size of this signal. First of all, the physics

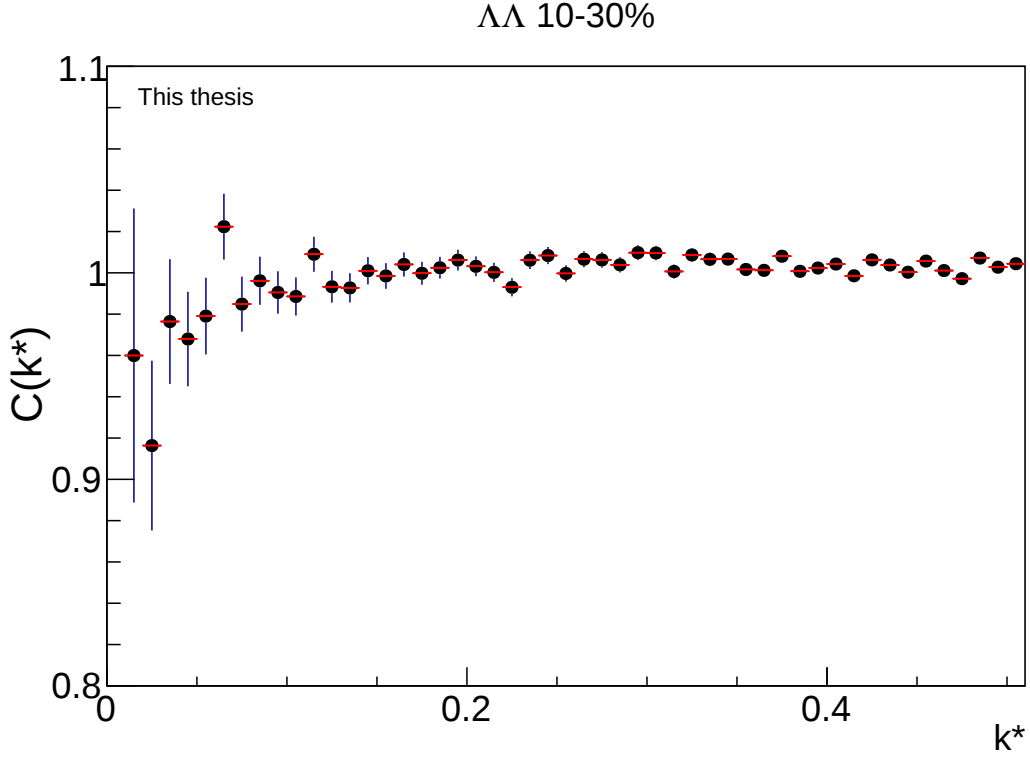


Figure 2.10: $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ correlation function for the 10-30% centrality range with statistical and systematic errors. A dip at low k^* is seen which is expected from quantum interference.

of the identical particle systems (see Sec. 3.1.4) is expected to be confined to just the few lowest bins. There may be competing physics effects – a suppression from quantum interference and an enhancement from the strong final state interactions – that wash each other out. Furthermore, two-track effects like splitting and merging show up in this region, though they are assumed to be well-corrected for by the aforementioned pair-wise cuts. The presence of residual correlations may also muddle or dilute the signal region. This will be discussed further in Section 3.2.

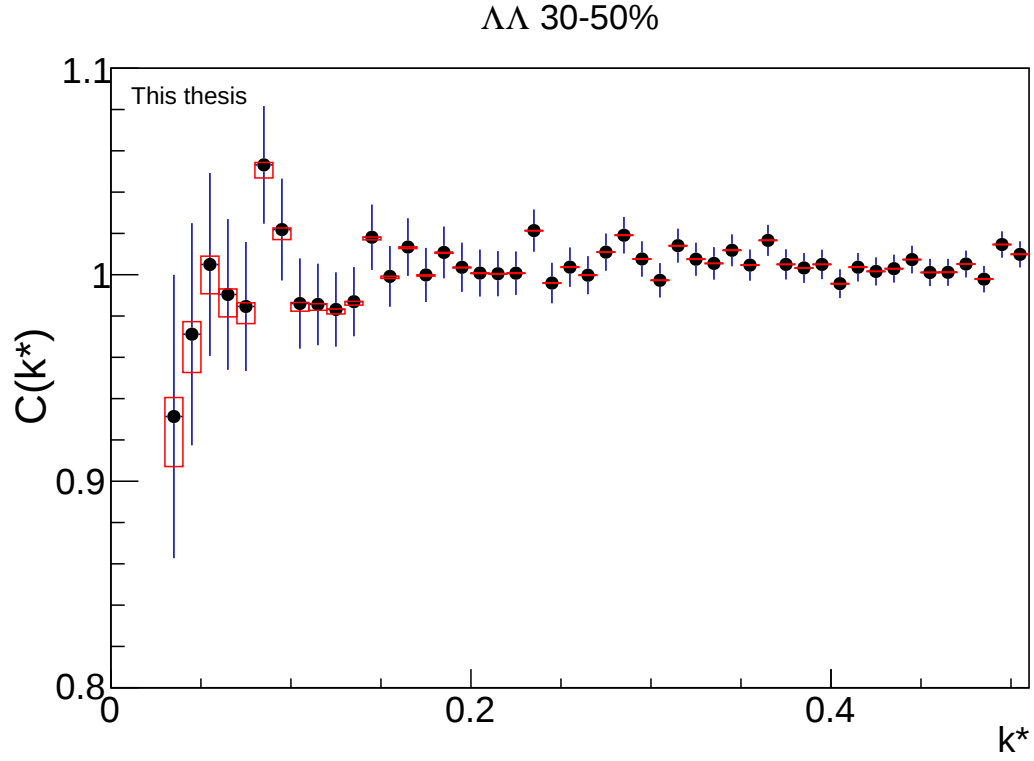


Figure 2.11: $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ correlation function for the 30-50% centrality range with statistical and systematic errors. Due to the poor statistics in the low k^* region, it is not feasible to include this correlation function in the fit analysis.

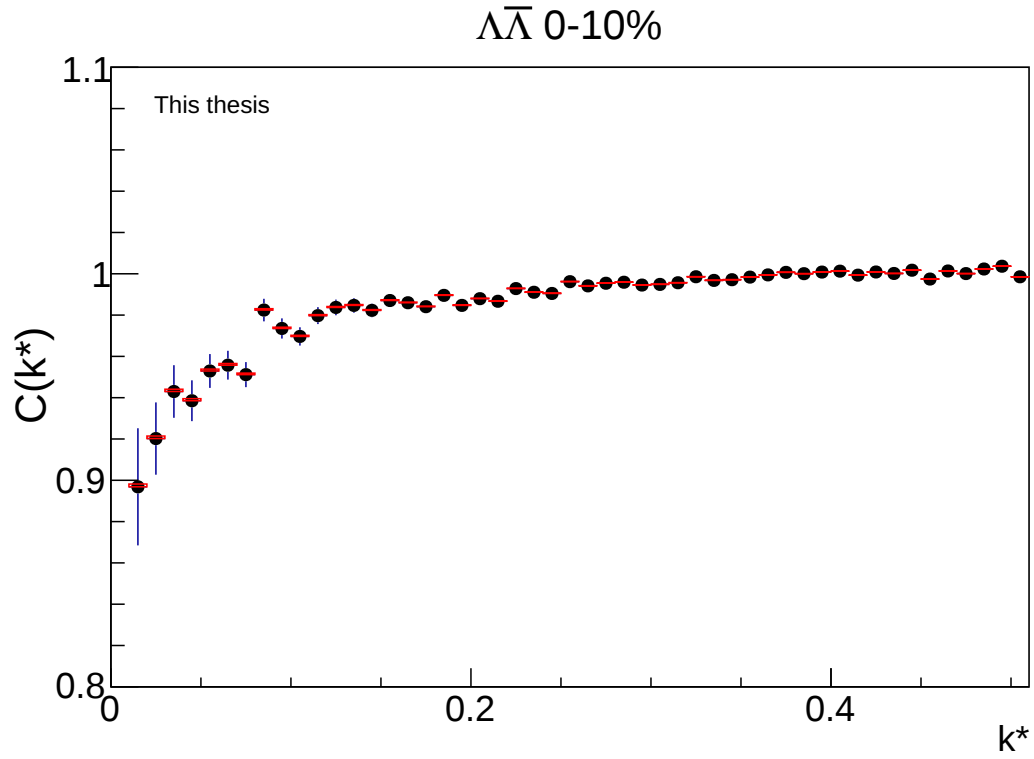


Figure 2.12: $\Lambda\bar{\Lambda}$ correlation function for the 0-10% centrality range with statistical and systematic errors. A wide suppression is seen which is indicative of annihilation.

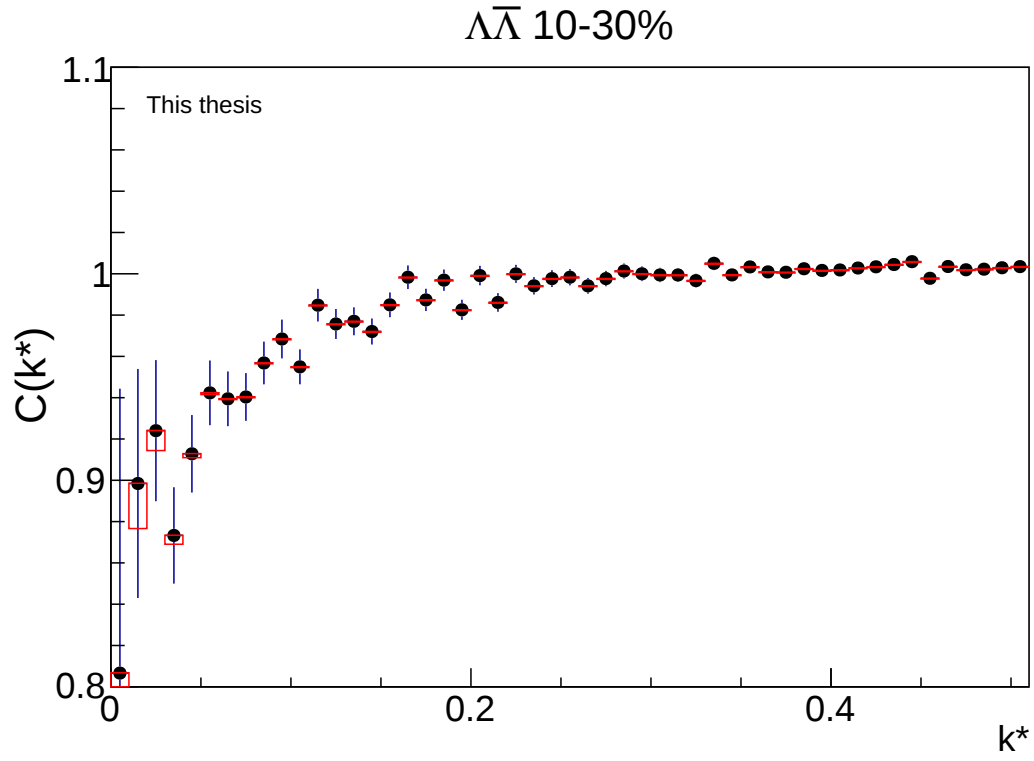


Figure 2.13: $\Lambda\bar{\Lambda}$ correlation function for the 10-30% centrality range with statistical and systematic errors. A wide suppression is seen which is indicative of annihilation.

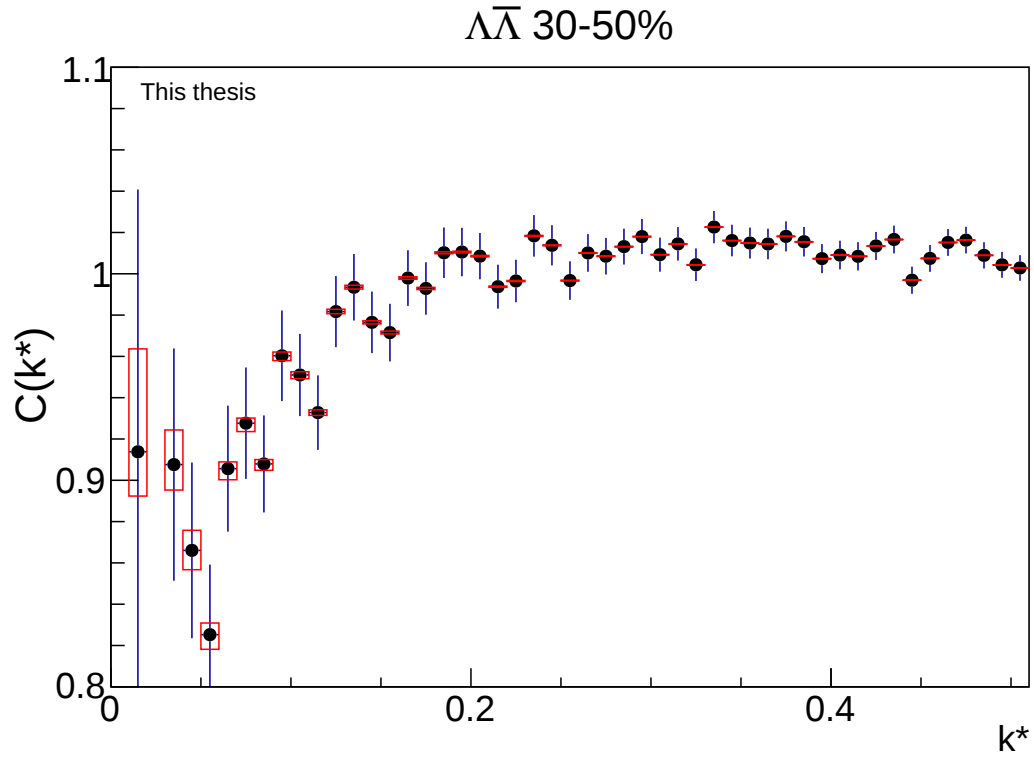


Figure 2.14: $\Lambda\bar{\Lambda}$ correlation function for the 30-50% centrality range with statistical and systematic errors. A wide suppression is seen which is indicative of annihilation.

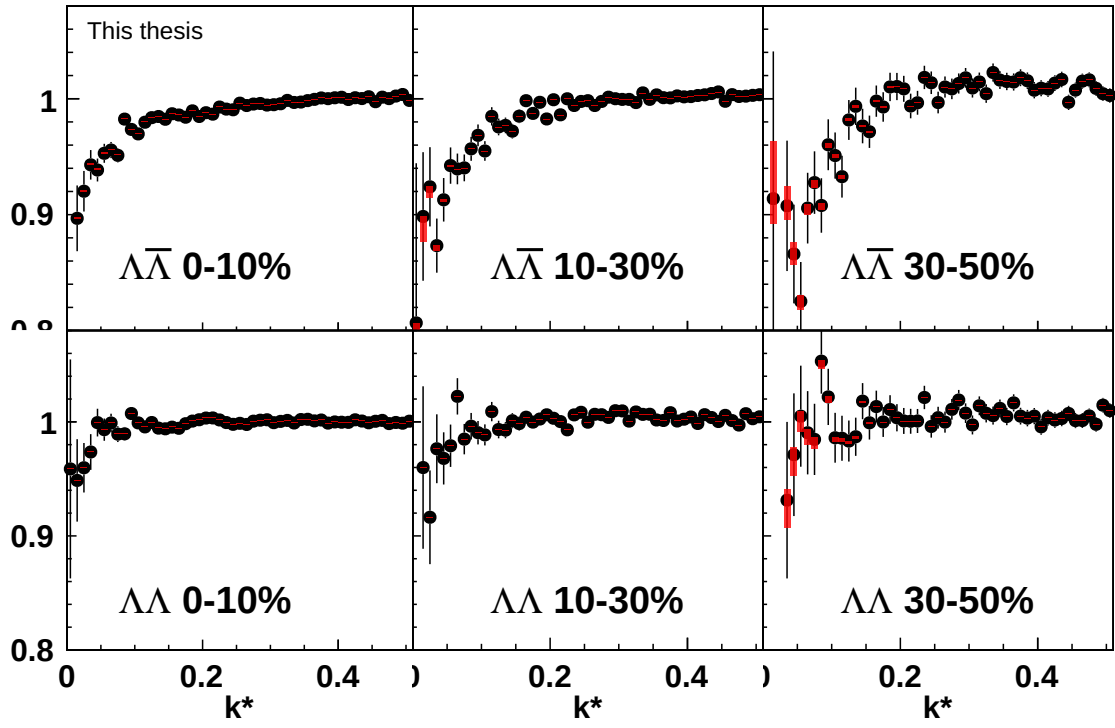


Figure 2.15: $\Lambda\bar{\Lambda}$ (top) and $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ (bottom) correlation functions for all three centrality ranges.

Chapter 3

DATA ANALYSIS

3.1 Correlation function model

In this section, we'll walk through the theoretical framework for analyzing femtoscopic correlations. We'll start with a general description of a two-particle correlation function via the Koonin-Pratt paradigm [57, 58, 13]. From there, we will discuss several ways to write the source function which characterizes the size and shape of the interaction region, and we will give extra attention to the spherically Gaussian source employed in this analysis. After that, we will describe the two-particle wave function that accounts for potentially spin-dependent interactions via low-energy scattering theory. We will use the chosen source function and wave function to derive an analytic form for the one-dimensional correlation function. Finally, we will discuss the observables that are accessible through this model.

3.1.1 Koonin-Pratt equation

As discussed in Section 2.3.2, correlation functions are constructed in such a way as to divide out uncorrelated phase-space effects and therefore to leave behind only the correlated part of the two-particle phase space. We assume that the uncorrelated two-particle phase space is just a product of two single-particle phase spaces. The single-particle emission function $S(x, p)$ (also referred to as the single-particle phase space or Wigner density) is related to the single-particle momentum distribution function of the species:

$$E_p \frac{dN}{d^3p} = \int d^4x S(x, p). \quad (3.1)$$

The two-particle correlation function can be expressed as a ratio of distributions:

$$C(p_1, p_2) = \frac{E_1 E_2 \frac{dN}{d^3p_1 d^3p_2}}{(E_1 \frac{dN}{d^3p_1})(E_2 \frac{dN}{d^3p_2})}, \quad (3.2)$$

where the numerator term is the two-particle momentum distribution, i.e. the distribution of pairs of particles with momenta p_1 and p_2 .

We can similarly write the correlation function as a ratio of emission functions. In that case, the denominator is just a product of two single-particle emission functions. Assuming the particles are created and emitted independently, the numerator is likewise a product of two single-particle emission functions, but we also include the two-particle wave function $\Psi(r, k)$ which introduces correlations between particles.

$$C(p_1, p_2) = \frac{\int d^4x_1 d^4x_2 S_1(x_1, p_1) S_2(x_2, p_2) |\Psi(r, k)|^2}{\int d^4x_1 d^4x_2 S_1(x_1, p_1) S_2(x_2, p_2)}. \quad (3.3)$$

The wave function only cares about the relative separation $r^\mu = x_1^\mu - x_2^\mu$ and the relative momentum k^μ of the particles. The four-vector momentum difference k^μ and its corresponding invariant relative momentum k^* are defined in Section 2.3.1. Note that $S_1 = S_2$ for pairs of identical particles.

As the wave function only cares about relative quantities, let us change variables. Let $p_1^\mu = P^\mu/2 + k^\mu$ and $p_2^\mu = P^\mu/2 - k^\mu$, where $P^\mu = p_1^\mu + p_2^\mu$. With that, we can rewrite the correlation function as

$$C(P, k) = \frac{\int d^4x_1 d^4x_2 S_1(x_1, P/2 + k) S_2(x_2, P/2 - k) |\Psi(r, k)|^2}{\int d^4x_1 d^4x_2 S_1(x_1, P/2 + k) S_2(x_2, P/2 - k)}. \quad (3.4)$$

We assume that for a given total momentum P and small relative momentum k (i.e. for particles with approximately the same momentum), the emission function is independent of the momentum difference

$$\begin{aligned} S_1(x_1, P/2 + k) &\approx S_1(x_1, P) \\ S_2(x_2, P/2 - k) &\approx S_2(x_2, P). \end{aligned} \quad (3.5)$$

To make the integral only care about relative quantities, we can define a two-particle emission function, which describes the separation of particles at the time(s) of their emission

$$S_{12}(r, P) = \frac{\int d^4x_1 d^4x_2 S_1(x_1, P) S_2(x_2, P) \delta(r - x_1 + x_2)}{\int d^4x_1 d^4x_2 S_1(x_1, P) S_2(x_2, P)}. \quad (3.6)$$

With this, we obtain

$$C(P, k) = \int d^3r S_{12}(r, P) |\Psi(r, k)|^2. \quad (3.7)$$

There is one last approximation. It is common to speak of the emission function for a given total momentum P or average transverse momentum $P_T = (p_{1,T} + p_{2,T})/2$. In principle, correlation functions can be created using infinitesimally small momentum bins, in which case a two-particle emission function $S_P(r)$ can be characterized for each of those P bins. In practice, the size of the P or P_T bins is chosen based on the statistical limi-

tations of the experiment; one uses as many bins as one can while still getting sufficient statistics. In any case, it is common to assume that the emission function is independent of momentum within each measured P bin. Because the emission function describes the region where particles have approximately the same momentum, that region is called the region of homogeneity [28].

With these simplifications, we can finally rewrite Eqn. 3.3 as

$$C_P(k) = \int d^3r S_P(r) |\Psi(r, k)|^2. \quad (3.8)$$

Equations 3.3 and 3.8 are both called the Koonin-Pratt equation. In the following sections, we'll discuss the two-particle emission function $S_P(r)$ that we employ to characterize the interaction region, as well as a general wave function that can describe interactions between both identical and non-identical particle pairs in the absence of Coulomb effects.

3.1.2 Source functions

As discussed above and in Section 1.3.3, the two-particle emission function describes the size and shape of the interaction region. However, the size of the source is not the size of the entire fireball at freezeout. Instead, it is the size of the region of homogeneity, the region where particles with total momentum P are emitted with the same momentum, such that $k^* = 0$.

In practice, there are a number of ways to characterize the source function. One of the simplest treatments is to assume that the source is described by a spherically symmetric function such as a Gaussian. One can then write the source with a single, one-dimensional parameter, often a Gaussian radius R . There are several weaknesses to this method. For example the source is likely not spherically symmetric — at the moment of the collision of the ions the interaction region roughly has the shape of a pancake or an almond, and it need not expand isotropically in the longitudinal and transverse directions. Additionally, a Gaussian may not always be the best approximation; sometimes an exponential or an Edgeworth expansion may work better [26, 59]. Despite these flaws, the one-dimensional approach has merit for analyses with limited statistics, since there is less risk of overfitting the data with too many parameters. And as we'll see below, there is an analytic form for the one-dimensional correlation function, which is not necessarily the case for multi-dimensional analyses. Because Λ yields are much smaller than π yields, and the $\Lambda\Lambda$ signal in particular is quite weak (see Figure 2.9), this study employs the 1D approach. We assume a normalized, one-dimensional, spherically gaussian source of width R ,

$$S_P(r) \sim \exp\left(-\frac{r^2}{4R^2}\right). \quad (3.9)$$

Though it is outside the scope of this analysis, it is worth mentioning the multidimensional

mensional fitting methods. There are two common approaches. One is the out-side-long prescription. For a detailed description, see [13]. Here, one decomposes both the source function and the relative momentum k^* into three orthogonal directions. One then calculates the relative momentum of the pair in the longitudinally comoving system (LCMS). The longitudinal direction is defined to point along the beam axis (i.e. the axis along with the ions travel on their collision course), so in the LCMS frame the pair has net longitudinal momentum $P_z = 0$. The out direction points along the total momentum of the pair $\vec{P} = \vec{P}_T$, and the side direction is perpendicular to the other two. Here, one usually assumes a Gaussian approximation for each direction with radii R_{out} , R_{side} , and R_{long} . In some cases, there may also be cross terms such as $R_{\text{out,side}}^2$, as well as effects arising from differences in the particles' emission times.

The other multidimensional approach involves decomposing the 3D correlation function $C(k^*)$ in terms of spherical harmonics $C_{l,m}(|k^*|, \theta, \phi)$ [60]. This has several advantages over the out-side-long method. For one, it is straightforward to plot a particular spherical component of the correlation function $C_{l,m}$, whereas there are ambiguities when making one-dimensional projections of $C(k^*)$. For another, various physical symmetries should cause certain components of $C_{l,m}$ to be zero. If they do not vanish, it usually a symptom that something is procedurally wrong with the analysis. Therefore, inspection of the spherical components can help with analysis debugging. Given the strong signal seen in this analysis's $\Lambda\bar{\Lambda}$ correlation functions (see Figure 2.12), said correlation functions may be a good candidate for spherical harmonic decomposition in a future study.

3.1.3 Wave function

For identical, non-interacting particles, the non-symmetrized two-particle wave function ψ is a product of plane waves

$$\begin{aligned}\psi(p_1, p_2, x_1, x_2) &= e^{-ip_1 x_1} e^{-ip_2 x_2} \\ &= e^{-iP(x_1+x_2)/2} e^{-ikr}\end{aligned}\tag{3.10}$$

where $P = p_1 + p_2$, $k = p_1 - p_2$, and $r = x_1 - x_2$. In the case of interacting particles, we instead have

$$\psi(p_1, p_2, x_1, x_2) = e^{-iP(x_1+x_2)/2} \left[e^{-ikr} + \phi(p_1, p_2, r) \right].\tag{3.11}$$

where $\phi(p_1, p_2, r)$ describes the final state interactions of the particles. Assuming that the range of the interaction potential is smaller than the distance between emission points r , the the magnitude of the momentum difference is much smaller than the mass of the particles ($k^* \ll m$), and that the particles are emitted at the same time, then the interaction term

can be approximated [61] by

$$\phi(p_1, p_2, r) = (1 + (-1)^S) f(k^*) \frac{e^{ik^*r}}{r}, \quad (3.12)$$

where

$$f(k^*) = \left(\frac{1}{f_0} + \frac{1}{2} d_0 k^{*2} - ik^* \right)^{-1} \quad (3.13)$$

is the s-wave scattering amplitude in the effective-range approximation [62], f_0 is the complex scattering length which characterizes the strength of the interaction, d_0 is the effective range of interaction, and S is the total spin of the pair. Thus, for identical particles, the asymptotic solution to the non-symmetrized two-particle wave function ψ takes the form of a superposition of an incoming plane-wave and an outgoing spherical wave

$$\psi(p_1, p_2, x_1, x_2) = e^{-iP(x_1+x_2)/2} \left[e^{-ikr} + (1 + (-1)^S) f(k^*) \frac{e^{ik^*r}}{r} \right]. \quad (3.14)$$

For non-identical particles, the full wave function $\Psi(r, k^*) = \psi(r, k^*)$, and there is no $(1 + (-1)^S)$ prefactor on the scattering amplitude. For identical, unpolarized, spin- $\frac{1}{2}$ fermions, the symmetrized wave function is given by

$$|\Psi(r, k^*)|^2 = \frac{1}{2} \left(\frac{3}{4} |\psi_{p_1 p_2}(r, k^*) - \psi_{p_2 p_1}(r, k^*)|^2 + \frac{1}{4} |\psi_{p_1 p_2}(r, k^*) + \psi_{p_2 p_1}(r, k^*)|^2 \right), \quad (3.15)$$

where $\frac{3}{4}$ of the pairs are in the symmetric spin-triplet state and antisymmetric spatial state, and the other $\frac{1}{4}$ are in the antisymmetric spin-singlet state and symmetric spatial state. In principle, the scattering length f_0 can be a spin-dependent quantity, as is the case for the p Λ interaction [45]. In this analysis, we will assume that singlet ($S = 0$) and triplet ($S = 1$) scattering parameters are spin-independent.

3.1.4 1D analytic model

We then evaluate the Koonin-Pratt equation (Eqn. 3.8) with the use of our two-particle source function (Eqn. 3.9) and symmetrized wave function (Eqn. 3.15). The details of the calculation can be found in [61]. The resulting analytic, one-dimensional, femtoscopic correlation function is called the Lednicky and Lyuboshits equation (sometimes just called the Lednicky equation), which can be broken down as

$$C(k^*) = 1 + C_{\text{QS}}(k^*) + C_{\text{FSI}}(k^*). \quad (3.16)$$

C_{QS} describes plane-wave quantum interference:

$$C_{\text{QS}}(k^*) = \alpha e^{-4k^{*2}R^2} \quad (3.17)$$

where R is the approximate 1D spherical source size, and $\alpha = (-1)^{2j}/(2j+1)$ for identical particles with spins j , and $\alpha = 0$ for non-identical particles. C_{FSI} describes the s-wave strong final state interaction of the particles:

$$C_{\text{FSI}}(k^*) = (1 + \alpha) \left[\frac{1}{2} \left| \frac{f(k^*)}{R} \right|^2 \left(1 - \frac{d_0}{2\sqrt{\pi}R} \right) + \frac{2\Re f(k^*)}{\sqrt{\pi}R} F_1(2k^*R) - \frac{\Im f(k^*)}{R} F_2(2k^*R) \right] \quad (3.18)$$

where $F_1(z) = \int_0^z dx e^{x^2 - z^2}/z$, and $F_2(z) = (1 - e^{-z^2})/z$. The $(1 + \alpha)$ pre-factor accounts for the relevant (non/anti-)symmetrization effects. The scattering amplitude f (Eqn. 3.13) is dependent upon the effective range of interaction d_0 , as well as the complex scattering length f_0 . The real part of the scattering amplitude can contribute either a positive or a negative correlation, but either way the effect is relatively narrow in k^* (on the order of about one hundred MeV/c). The imaginary part of the scattering amplitude accounts for inelastic processes of baryon-antibaryon annihilation. Introducing a non-zero imaginary part to the scattering length produces a wide (hundreds of MeV/c) negative correlation.

Technically, the scattering length and effective range are spin dependent. It should be noted that in the case of identical fermions, “the contributions of s-wave interaction of identical nucleons to the correlation function goes to zero [for summary spin state $S = 1$] (identical nucleons with parallel spins cannot be in the s-wave state)” [61]. As a result, only the spin singlet (i.e. antisymmetric spinor, symmetric spatial wave function) scattering parameters are measured for identical fermions. Such is the case for the $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$ correlation measurements. However, for non-identical particles such as $\Lambda\bar{\Lambda}$, the factor $(1 + \alpha)$ could be replaced with a weighted sum of C_2 calculated using separate spin-dependent scattering parameters. That calculation would use ρ_S , the fraction of pairs in each total spin state S , as the weight. For statistical reasons, this analysis will eschew spin-dependent measurements of the $\Lambda\bar{\Lambda}$ scattering parameters, and instead attempt to fit for a spin-averaged value.

3.1.5 Observables

For analyses such as proton-proton or proton-neutron, the scattering lengths f_0 and effective range of interaction d_0 are well specified by years of scattering data, and they can be fixed to the known values. In the case of lambda-(anti-)lambda femtoscopy, the interaction is heretofore largely unstudied. For this analysis, the scattering lengths and effective range are left as free parameters. Many analyses introduce another free parameter, λ , as rough measure of the pair purity. For example, $\lambda_{\Lambda\Lambda}$ corresponds to the ratio of number of true primary $\Lambda\Lambda$ pairs to the total number of pairs in the correlation function. The remaining fraction is comprised of all the other combinations of pair types, such as pairs where one Λ is real and primary and the other is fake, or pairs between primary and secondary lambdas.

Our implementation of λ parameters is described in Section 3.2.1, and estimates of those parameters are given in Section 3.3. The last fit parameter is R_{inv} , which characterizes the size of the one-dimensional two-particle emission function.

3.2 Residual correlations

In the hadronic rescattering phase of a heavy-ion collision, Λ can scatter off of Σ , Ξ , and Ω baryons, though little is known about the cross section of these interactions. In principle, one could do a femtoscopic analysis of $\Lambda\Sigma^0$, $\Lambda\Xi^-$, or even $\Lambda\Omega$, in order to measure the scattering length and effective interaction range of these pairs. One could reconstruct many pairs of these particles, make correlation functions, and fit them with Eqn. 3.18. In practice, difficulties in topological reconstruction and small particle yields make such analyses challenging at best. However, a ghost of these interactions haunts $\Lambda\Lambda$ femtoscopic analysis in the form of residual correlations.

The aforementioned hyperons all decay into Λ baryons. The momentum difference between these particles and their daughter Λ (caused by the decay) is small relative to the total momentum of the particles, such that daughter and parent carry very similar momenta. As a result, relative-momentum correlation functions between primary particles and secondary particles can be sensitive to the interactions between the primary particles and the parent particles. For example, the correlation between primary Λ and primary Σ^0 should be visible, though somewhat smeared out, in the correlation between primary Λ and secondary Λ . Furthermore, these secondary Λ are often difficult to distinguish from primary Λ . Therefore, measurements of two- Λ correlation functions will contain a mixture of primary correlations and feeddown correlations, and an attempt to parametrize the correlation functions should take this into account.

In this analysis note, we will discuss two different methods of accounting for these residual correlations. In the “Transformed residuals method”, the Lednický equation is used to estimate the shape of the residual correlations *before* they decay (e.g. to estimate the shape of the $\Lambda\Sigma$ correlation function in the $\Lambda\Sigma$ relative momentum space). This pre-feeddown correlation is then smeared into a residual correlation function (in the $\Lambda\Lambda$ momentum space) using a transform matrix that accounts for the kinematic effects of the decay of one or both of the particles. Different sources of residual correlations are determined separately and then summed together with the primary-primary correlation in the fit procedure.

In contrast, the “Gaussian residuals method” attempts to describe the full residual correlation effect phenomenologically. Rather than precisely describe each component of residual correlation, the full shape of the residual correlation is described with a Gaussian parametrization.

3.2.1 Transformed residuals method

One example of the transformed residuals method can be seen in the recent analysis [46] of STAR $p\bar{\Lambda}$ correlation functions measured in 200 GeV Au–Au collisions. In this study, $p\bar{\Lambda}$ and $\bar{p}\Lambda$ correlations were fit with and without residual correlations. The fits with residual correlations included feeddown and feedup contributions from numerous sources. The study found that the fit with residual correlations measured a source size comparable to the previously measured sizes of $p\Lambda$ and $\bar{p}\bar{\Lambda}$. In contrast, the fits that did not include residual correlations measured source radii that deviated significantly from the old $p\Lambda$ and $\bar{p}\bar{\Lambda}$ results. One conclusion from that study was that it is important to make a precise characterization of the λ parameters associated with each type of pair contribution. As mentioned above in Section 3.1.5, the λ parameter can be interpreted as the fraction of total pairs associated with that pair type. We attempt to quantify the residual contamination in this analysis by simultaneously fitting the data for both the primary correlation function and the residual correlations. The λ parameter acts as a weight for combining the primary and secondary correlations, via

$$C_{\text{total}}(k_{\Lambda\Lambda}^*) = \sum_i \lambda_i C_i(k_{\Lambda\Lambda}^*), \quad (3.19)$$

where i runs over the different pair types, and $\sum \lambda_i = 1$, and $C_i(k_{\Lambda\Lambda}^*)$ is a theoretical correlation function (see Eqn. 3.16) expressed in the basis of the relative momentum of primary and/or secondary $\Lambda\Lambda$ pairs.

For example, $\Lambda\Lambda$ correlations with $\Lambda\Sigma^0$ feeddown would be fit using

$$C_{\text{total}}(k_{\Lambda\Lambda}^*) = 1 + \lambda_{\Lambda\Lambda}[C_{\Lambda\Lambda}(k_{\Lambda\Lambda}^*) - 1] + \lambda_{\Lambda\Sigma}[C_{\Lambda\Sigma}(k_{\Lambda\Lambda}^*) - 1], \quad (3.20)$$

where

$$C_{\Lambda\Sigma}(k_{\Lambda\Lambda}^*) \equiv \frac{\sum_{k_{\Lambda\Sigma}^*} T(k_{\Lambda\Lambda}^*, k_{\Lambda\Sigma}^*) C_{\Lambda\Sigma}(k_{\Lambda\Sigma}^*)}{\sum_{k_{\Lambda\Sigma}^*} T(k_{\Lambda\Lambda}^*, k_{\Lambda\Sigma}^*)}, \quad (3.21)$$

$C_{\Lambda\Sigma}(k_{\Lambda\Sigma}^*)$ is the $\Lambda\Sigma$ correlation function calculated from Eqn. 3.16, and T is a transform matrix generated with THERMINATOR [63] (see Section 3.3.2), which shows how the k^* of pairs of particles kinematically transforms when one of the particles decays. Figure 3.1 shows the normalized transform matrix for $\Lambda\Sigma \rightarrow \Lambda\Lambda$. The width of the smearing band is characterized by the decay momentum, 74 MeV/c for $\Sigma^0 \rightarrow \Lambda + \gamma$. See Section 3.4.1 for details about the normalization of T . Further residual correlations (e.g. from $\Sigma\Sigma \rightarrow \Lambda\Lambda$) can be included via additional $C_i(k_{\Lambda\Lambda}^*)$ terms.

Figures 3.2–3.7 show normalized residual correlations for $\Sigma^0\Lambda$, $\Sigma^0\Sigma^0$, $\Sigma^0\Xi^0$, $\Sigma^0\Xi^-$, $\Xi^0\Lambda$, $\Xi^-\Lambda$, and $\Xi^0\Xi^-$. The width of double hyperon decays $YY \rightarrow \Lambda\Lambda$ is generally wider than

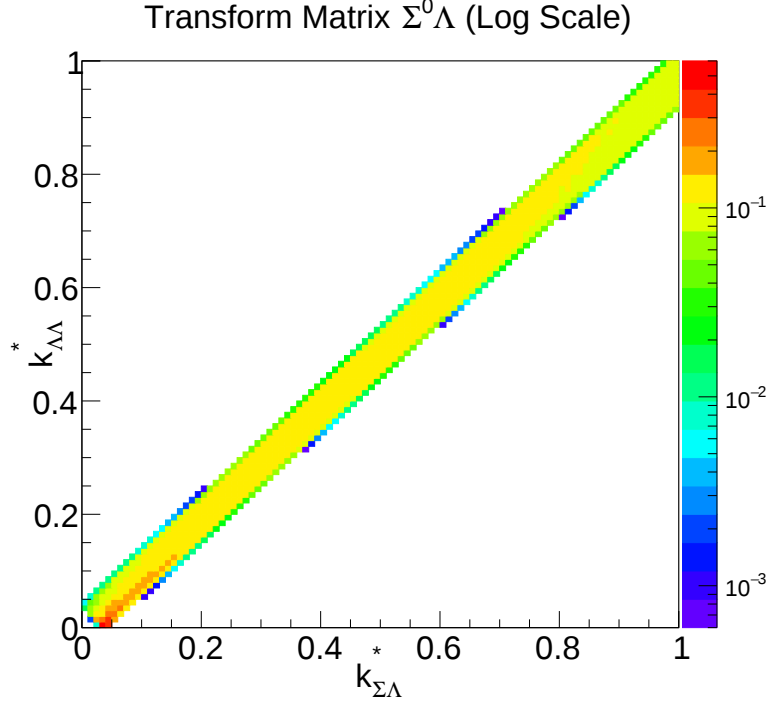


Figure 3.1: THERMINATOR [63] transform matrix showing how the relative momentum of $\Sigma^0\Lambda$ pairs transforms into relative momentum of $\Lambda\Lambda$ pairs after the Σ^0 decays. Each row is normalized to unity.

single decays $Y\Lambda \rightarrow \Lambda\Lambda$ because each decay contributes a momentum shift. The antihyperon decays $\bar{Y} \rightarrow \bar{\Lambda}$ have the same kinematics as the hyperon decays, so Figures 3.1–3.7 are also used for their respective antiparticle-antiparticle and particle-antiparticle pair smearing.

In a simple case with only two pair types ($\Lambda\Lambda$ and $\Lambda\Sigma$), the λ parameters could either be fixed to some estimated experimental value, or they could be left as free fit parameters. However, in this analysis it is necessary to account for all the pair types discussed above. See Section 3.3.4 for estimates of λ parameters as well as estimates of which residual correlations must be included and which can be safely neglected. Because each residual correlation is calculated using the Lednicky equation, it is necessary to estimate, fit, or otherwise constrain the femtoscopic radius, R , and scattering parameters, f_0 and d_0 , of the residual correlation pair. Those constraints are described further in Section 3.5.5.

3.2.2 Gaussian residuals method

The Gaussian residuals method has been used to fit the STAR $p\bar{\Lambda}$ data [64], as well as 200 GeV STAR $\Lambda\Lambda$ data [49].

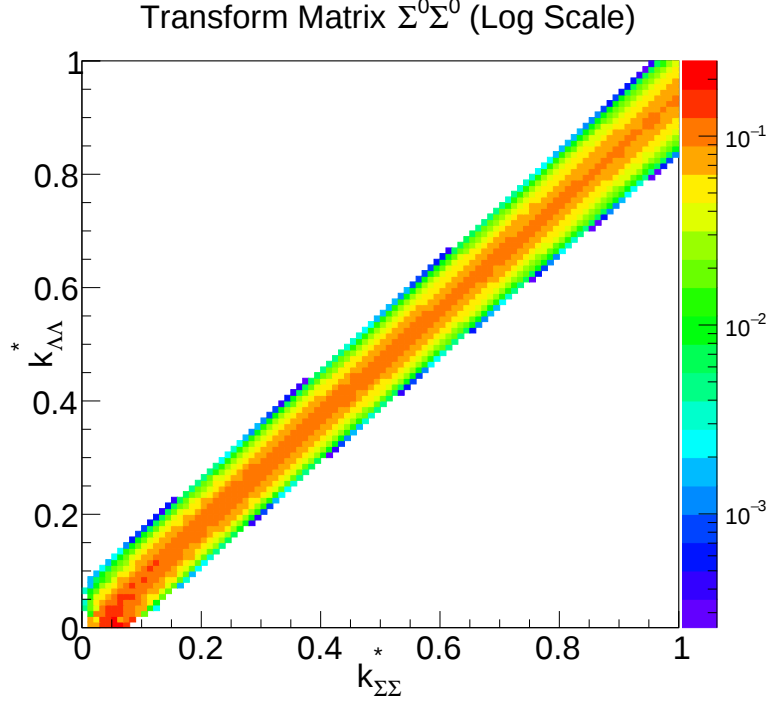


Figure 3.2: THERMINATOR [63] transform matrix showing how the relative momentum of $\Sigma^0\Sigma^0$ pairs transforms into relative momentum of $\Lambda\Lambda$ pairs after both Σ^0 have decayed. Each row is normalized to unity. The smearing effect is wider here than for $\Sigma^0\Lambda$ because there are two decays happening.

In the Gaussian residual method, the measured correlation function is fit with

$$C_{\text{meas}}(k^*) = \lambda_{\text{prim}} * C_{\text{prim}}(k^*) + (1 - \lambda_{\text{prim}})(1 - \beta e^{-4k^{*2}\omega^2}) \quad (3.22)$$

where the phenomenological parameters β and ω help fit the roughly Gaussian-shaped residual correlations that are expected. The interpretation of β and ω is tenuous, so if one has reasonable estimates for the different components (λ parameters) of the residual correlations, as well as for the strength of the residual interactions (the scattering parameters f_0 and d_0), one is probably better off using the more rigorous Transform method of Section 3.2.1.

3.3 λ parameters

We described how λ parameters are used in Section 3.2.1. In this section, we'll estimate the λ parameter for the primary-primary correlation function, $\lambda_{\Lambda\Lambda}$, as well as for all the relevant residual correlations. As stated earlier, λ_{XY} loosely describes the fraction of pairs in the

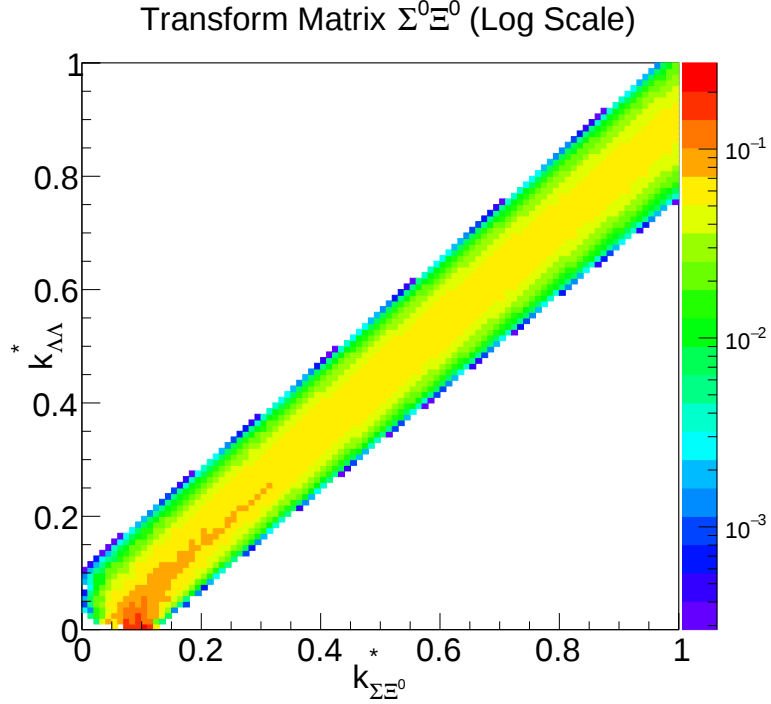


Figure 3.3: THERMINATOR [63] transform matrix showing how the relative momentum of $\Sigma^0\Xi^0$ pairs transforms into relative momentum of $\Lambda\Lambda$ pairs after both parents have decayed. Each row is normalized to unity. The smearing effect is wider here than for $\Sigma^0\Lambda$ because there are two decays happening.

correlation function composed of particles from species X and species Y. In this analysis, there are too many contributions from different types of pairs to leave the various λ as free parameters. Instead, we will estimate them beforehand and use those estimates as fixed inputs in our fit procedure.

First, we will develop a model for calculating λ parameters. In order to do these λ calculations, we will need estimates of typical Pb-Pb event yields for Λ and Λ parent species. We also need to know how detector reconstruction efficiencies affect the reconstruction rate of Λ of different origins.

3.3.1 Model of λ parameters

The experimental correlation function is defined as

$$C_{\text{Exp}}(k^*) \equiv \frac{1}{\gamma} \frac{A(k^*)}{B(k^*)}, \quad (3.23)$$

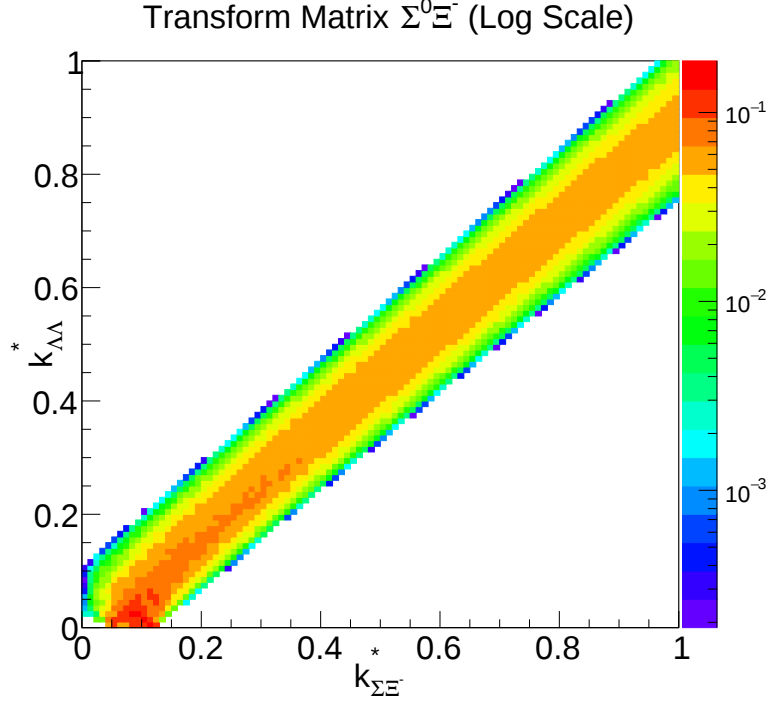


Figure 3.4: THERMINATOR [63] transform matrix showing how the relative momentum of $\Sigma^0 \Xi^-$ pairs transforms into relative momentum of $\Lambda \Lambda$ pairs after both parents have decayed. Each row is normalized to unity. The smearing effect is wider here than for $\Sigma^0 \Lambda$ because there are two decays happening.

where $A(k^*)$ is the same-event relative momentum distribution, $B(k^*)$ is the mixed-event relative momentum distribution, and γ is a normalization factor defined to set the high- k^* background region to unity. In what follows, we will investigate the relationship between $A(k^*)$ and $B(k^*)$.

Let us define N_{ij} as the number of $\Lambda \Lambda$ pairs that originated as type i (e.g. $\Lambda \Sigma^0$ or $\Xi^- \Xi^0$) in event j . Next we define $D_{ij}(k^*)$ as the normalized phase-space distribution $\frac{dN}{dk^*}$ of pair type i in event j . Putting these together, we can write the total relative momentum distribution for a single event as

$$A_j(k^*) = \sum_i^m N_{ij} D_{ij}(k^*),$$

where m is the number of different pair types. Then the total relative momentum distribu-

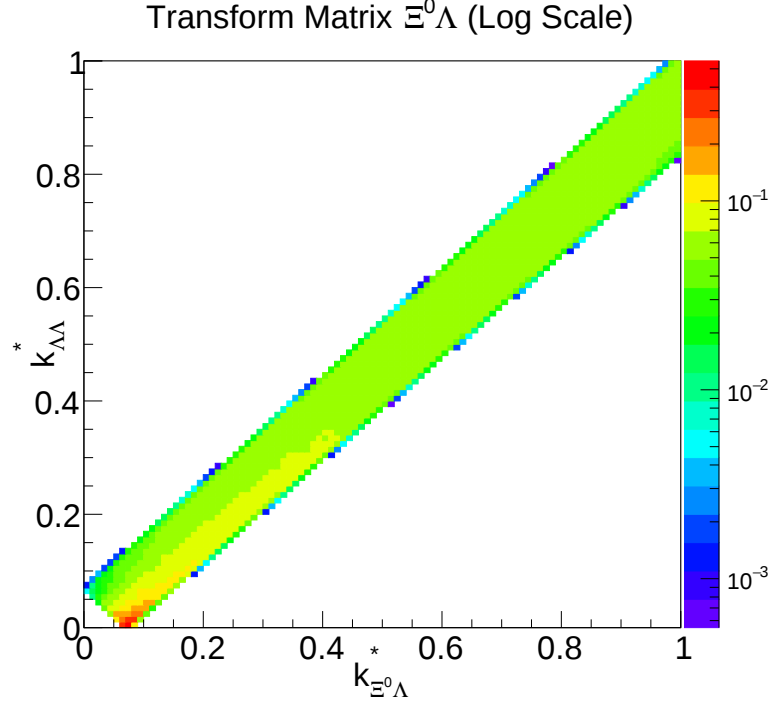


Figure 3.5: THERMINATOR [63] transform matrix showing how the relative momentum of $\Xi^0\Lambda$ pairs transforms into relative momentum of $\Lambda\Lambda$ pairs after the Ξ^0 has decayed. Each row is normalized to unity.

tion, composed of many events, can be written as

$$A(k^*) = \sum_j^n \sum_i^m N_{ij} D_{ij}(k^*),$$

where n is the total number of events. Similarly, the total relative momentum distribution for mixed-event pairs can be written as

$$B(k^*) = \sum_j^{n'} \sum_i^m N'_{ij} E_{ij}(k^*),$$

where N'_{ij} is the number of pairs of type i in mixed-event j , n' is the total number of mixed-events, and $E_{ij}(k^*)$ is the normalized, underlying relative momentum distribution.

Synthesizing Eqns. 3.19 and 3.23 yields

$$C_{\text{Exp}} = \frac{1}{\gamma} \frac{A(k^*)}{B(k^*)} = \frac{1}{\gamma} \frac{\sum_j^n \sum_i^m N_{ij} D_{ij}(k^*)}{\sum_j^{n'} \sum_i^m N'_{ij} E_{ij}(k^*)} = \sum_i^m \lambda_i C_i(k^*). \quad (3.24)$$

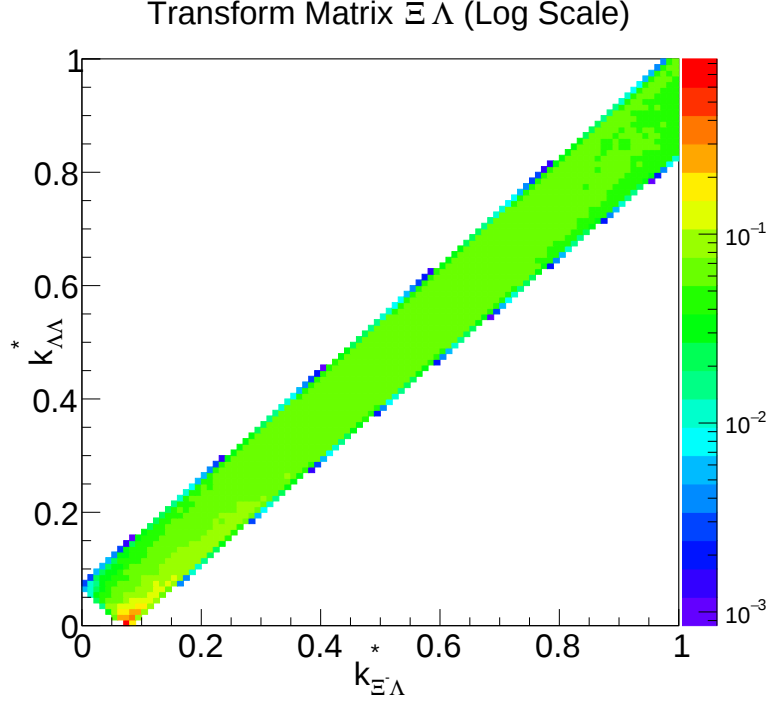


Figure 3.6: THERMINATOR [63] transform matrix showing how the relative momentum of $\Xi^- \Lambda$ pairs transforms into relative momentum of $\Lambda \Lambda$ pairs after the Ξ^- has decayed. Each row is normalized to unity.

We can look at the component for each type of pair separately:

$$\lambda_i C_i(k^*) = \frac{1}{\gamma} \frac{\sum_j^n N_{ij} D_{ij}(k^*)}{\sum_{j'}^{n'} \sum_i^m N'_{ij} E_{ij}(k^*)}. \quad (3.25)$$

We now make several assumptions. First, it is commonly assumed that the same-event relative momentum distribution is factorizable into a term describing the femtoscopic physics effects and a term describing the simple phase-space of relative momentum (i.e. the convolution of two single-particle momentum phase spaces). In that case, we can write

$$D_i(k^*) = C_i(k^*) E_i(k^*), \quad (3.26)$$

where $C_i(k^*)$ is the correlation function one would measure for pair type i if one could detect said pairs with perfect purity. Second, we assume that the same-event phase space distribution of a pair type $D_{ij}(k^*)$ is constant across all events that have the same centrality and z-vertex binning. Then we can pull D_{ij} out of the summation, dropping the event index

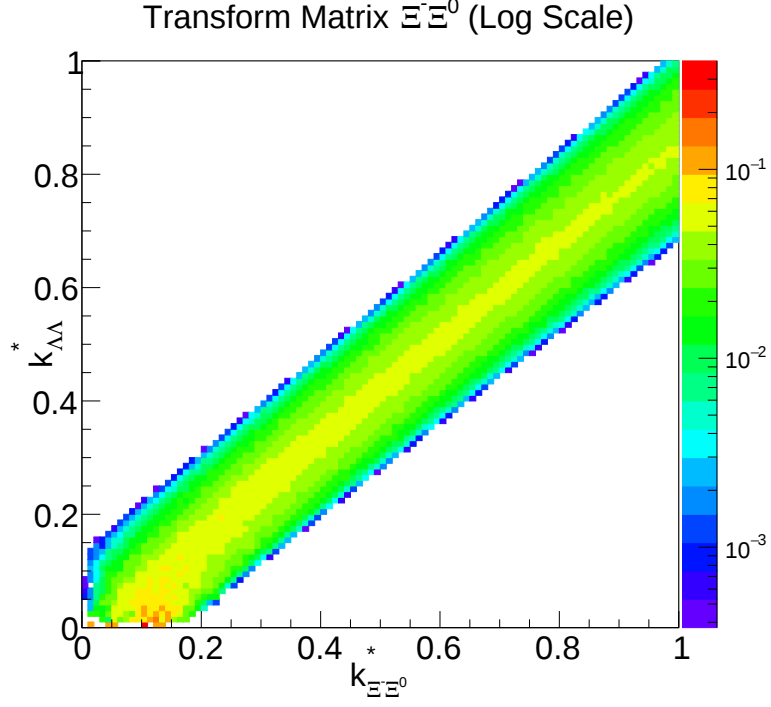


Figure 3.7: THERMINATOR [63] transform matrix showing how the relative momentum of $\Xi^-\Xi^0$ pairs transforms into relative momentum of $\Lambda\Lambda$ pairs after the both parents have decayed. Each row is normalized to unity. The smearing effect is wider here than for $\Sigma^0\Lambda$ because there are two decays happening.

j :

$$\lambda_i C_i(k^*) = \frac{1}{\gamma} \frac{D_i(k^*) \sum_j^n N_{ij}}{\sum_j^{n'} \sum_i^m N'_{ij} E_{ij}(k^*)}. \quad (3.27)$$

Third, we assume that the mixed-event pair phase space $E_{ij}(k^*)$, which only describes the combinatorics of pair relative momentum, is constant across all pair types and all events. Then we can pull out E_{ij} as well:

$$\lambda_i C_i(k^*) = \frac{1}{\gamma} \frac{D_i(k^*) \sum_j^n N_{ij}}{E_i(k^*) \sum_j^{n'} \sum_i^m N'_{ij}}. \quad (3.28)$$

Next, we divide out $C_i(k^*) = D_i(k^*)/E_i(k^*)$ on both sides of the equation to get an expression for λ_i

$$\lambda_i = \frac{1}{\gamma} \frac{\sum_j^n N_{ij}}{\sum_j^{n'} \sum_i^m N'_{ij}}. \quad (3.29)$$

Subject to a choice of normalization parameter γ , this is an expression for λ framed entirely in terms of numbers of pairs of particles in events. Such an expression can be evaluated

using representative Monte Carlo events.

We can clean this expression up further. It is common practice to use integrated pair counts for correlation function normalization. Often, the numerators and denominators are normalized separately:

$$C(k^*) = \frac{A(k^*) / \sum_j^n \sum_i^m N_{ij}}{B(k^*) / \sum_j^{n'} \sum_i^m N'_{ij}} \quad (3.30)$$

This leads to the normalization parameter

$$\gamma = \frac{\sum_j^n \sum_i^m N_{ij}}{\sum_j^{n'} \sum_i^m N'_{ij}}. \quad (3.31)$$

Substituting this into Eqn. 3.29 yields

$$\lambda_i = \frac{\sum_j^{n'} \sum_i^m N'_{ij}}{\sum_j^n \sum_i^m N_{ij}} \frac{\sum_j^n N_{ij}}{\sum_j^{n'} \sum_i^m N'_{ij}} = \frac{\sum_j^n N_{ij}}{\sum_j^n \sum_i^m N_{ij}} \quad (3.32)$$

We see that the number of mixed-event pairs drops out of the equation entirely. We can introduce a $1/n$ factor to both the numerator and denominator to get

$$\lambda_i = \frac{\frac{1}{n} \sum_j^n N_{ij}}{\frac{1}{n} \sum_j^n \sum_i^m N_{ij}}. \quad (3.33)$$

The expression $\frac{1}{n} \sum_j^n$ is, of course, an event averaging. We finally conclude that

$$\lambda_i = \frac{\langle N_i \rangle_{\text{event}}}{\langle \sum_i^m N_i \rangle_{\text{event}}}. \quad (3.34)$$

In plain terms, the λ parameter for pair type i can be computed as the ratio of the average number of pairs of that type per event to the average number of total pairs per event. The following sections will discuss how we use Monte Carlo data to estimate these numbers.

3.3.2 Thermal yield estimates

In order to compute the average numbers of pairs per event, as discussed in the previous section, we need two things: a Monte Carlo generator with reliable particle yields and many events. HIJING is notoriously unreliable for producing accurate hyperon spectra. HIJING Pb-Pb events are generated from a superposition of PYTHIA p—p collisions, so they lack the strangeness saturation of a fully thermalized medium. Instead of HIJING, we'll use THERMINATOR 2 (THERMal heavy IoN generATOR version 2) [63] to generate events.

THERMINATOR 2 is a Monte Carlo event generator designed to use statistical hadronization to produce hadronic freeze-out spectra of relativistic heavy-ion collisions. The user selects a hypersurface model and freeze-out model (several of which are included in the package), as well as rapidity range and the type of primordial particle multiplicity distribu-

tion (e.g. Poissonian). One of the crucial features of THERMINATOR is the inclusion of a large number of resonances from the particle data tables. These play an important role in the analysis, as a sizable number of these resonances decay into hyperons, thus boosting the hyperon yields. The resonances that undergo strong-decay into hyperons do so on time-scales on the order of fm/c or tens of fm/c. These decays happen on the same time-scale as the hadronic rescattering phase, so it is possible that the particle has decayed into its hyperon daughter before kinetic freeze-out. In that case, it is the final-state interaction of the hyperon daughter that is relevant to femtoscopic analyses, not the interactions of the resonant parent. In this analysis, we will assume that all short-lived resonances decay before kinetic freezeout. Under that assumption, we can treat the daughters of these resonances as primary particles.

In this analysis, we used the Lhyquid3v freeze-out model for 2.76 TeV Pb-Pb events. This freeze-out model is not packaged with THERMINATOR 2 — it was obtained via correspondence with the group that authored the THERMINATOR 2 package. There are several hypersurfaces available for this model. We chose the one with the most-central impact parameter, $b = 2.3$ fm, and a hadronic freeze-out temperature of 140 MeV. While the default is to use a rapidity window of ± 12 units, we set our rapidity window to ± 2 units to conserve disk space.

Using THERMINATOR, we generated 2000 events. In those events, we found all the primary and secondary Λ , ignoring those that fall outside the pseudorapidity and p_T cuts used in the data collection. For each Λ we found in a THERMINATOR event, we generated a random number between 0 and 1. We ignored the Λ if the generated number was higher than the Λ 's estimated reconstruction efficiency of $\sim 15\%$ (see Section 3.3.3 for more details on efficiency estimates). We tallied the reconstructed Λ and noted their origins. We also added fake Λ counts to each event to account for reconstruction purity. As discussed in Section 2.2.1, we have an approximately 95% purity. To add fakes, we looked at the average number of real particles tallied per event and determined the average number of fakes corresponding to a 5% impurity. Then for each event, we generated a number of fakes according to a Poissonian distribution peaked at the estimated value. The results without efficiency corrections are visible in Figure 3.8. The results with efficiency corrections can be seen in Figure 3.9.

Finally we used the tallied Λ to compute the number of pairs of each type per event (including real-real, real-fake, and fake-fake), along with the total number of pairs in each event, and we recorded that information. Events with no pairs (i.e. events that had one or fewer reconstructed Λ) were not included in event averaging, as they are skipped in the correlation function measurement of the real data.

The following section describes how we estimated the reconstruction efficiencies that were used in this process. The results of the λ parameter calculations can be found in 3.3.4.

Average Yield of Λ by Origin without Efficiency Correction

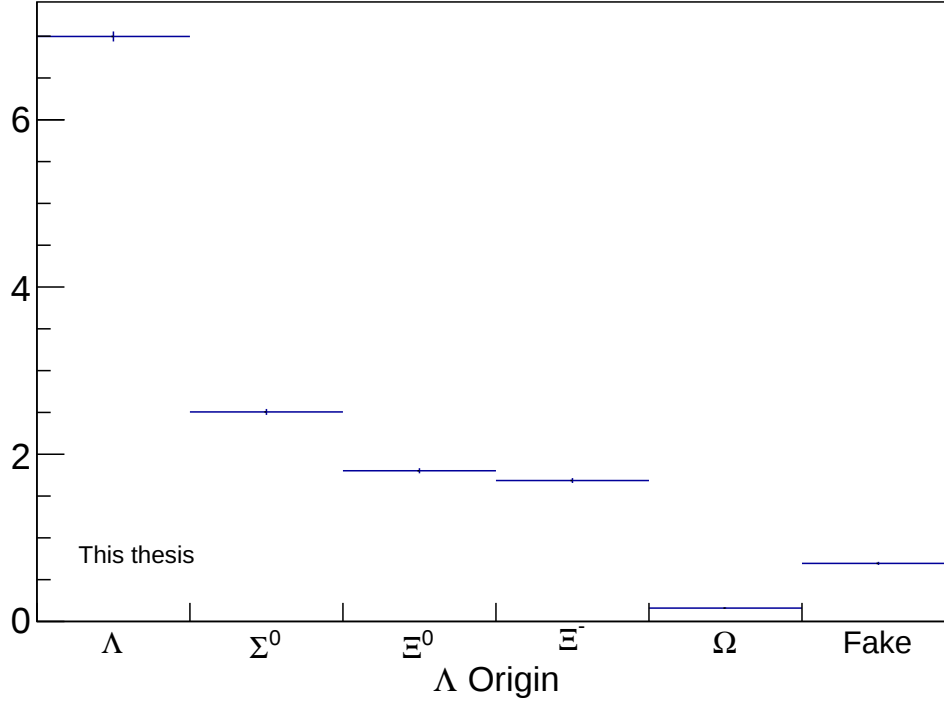


Figure 3.8: Expected yield of Λ in each event arising from different sources, as predicted by the thermal hadronization model in THERMINATOR 2 [63]. The x-axis shows the different origins: primary Λ , secondary Λ coming from various hyperons, and fakes (misidentified particles). Λ coming from the strong decay of short-lived resonances are included in the primary Λ count. The number of fakes in each event are generated via a Poissonian distribution, assuming an $\approx 5\%$ impurity. Reconstruction inefficiencies, which naturally reduce the number of reconstructed particles, are not included in this plot.

Average Yield of Λ by Origin with Efficiency Correction

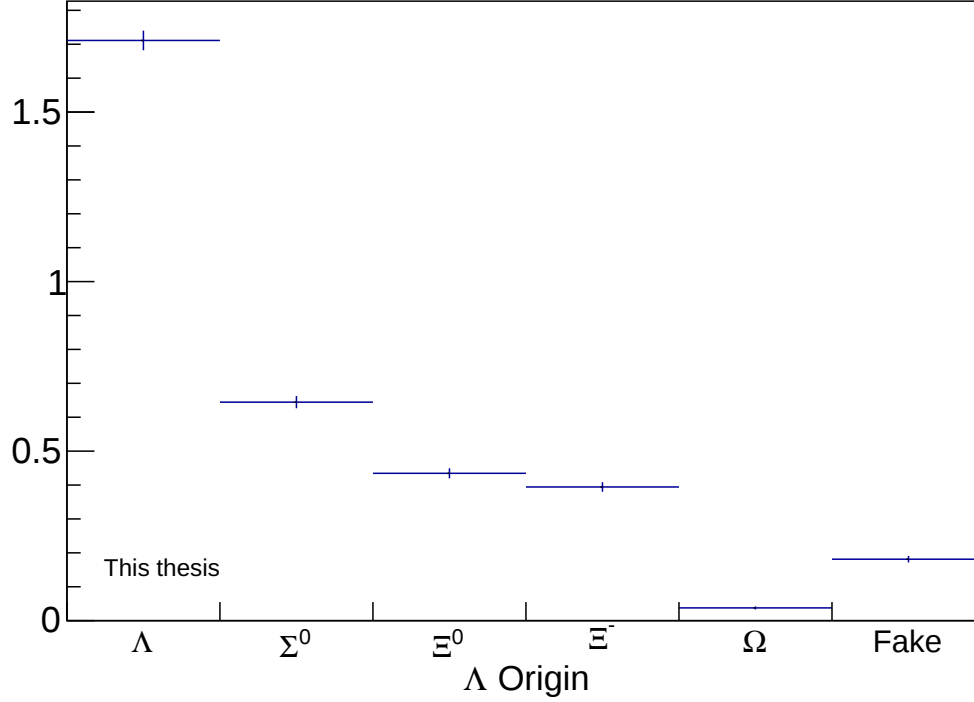


Figure 3.9: Expected yield of Λ in each event arising from different sources, as predicted by the thermal hadronization model in THERMINATOR 2 [63]. The x-axis shows the different origins: primary Λ , secondary Λ coming from various hyperons, and fakes (misidentified particles). Λ coming from the strong decay of short-lived resonances are included in the primary Λ count. The number of fakes in each event are generated via a Poissonian distribution, assuming an $\approx 5\%$ impurity. Reconstruction efficiency corrections (estimated in Section 3.3.3) are included in these yields.

3.3.3 Reconstruction efficiency effects on relative particle yields

In the previous section, we described how we used the THERMINATOR 2 event generator to estimate yields of primary and secondary Λ . Those yields were produced according to a statistical hadronization model. However, in real data we do not reconstruct perfect thermal spectra of particles. Instead, the yields are tempered by reconstruction efficiencies. The ALICE detector does not have perfect track resolution, so some tracks are misidentified or lost. We also employ topological reconstruction cuts to remove fake Λ candidates. Because of the finite detector resolution, some of the real Λ are removed as well. In addition, the Λ topological reconstruction cuts have been selected to preferentially remove secondary Λ coming from Ξ and Ω decays (see Section 2.2.1). This is not 100% effective at removing secondary Λ , and it also causes some primary Λ to be thrown out. Finally, the branching ratio of the Λ weak decay results in roughly a third of the Λ decaying into untrackable neutral particles. These effects naturally result in smaller measured yields and thus smaller quantities of pairs. The reconstruction efficiency is an important component of our estimation of λ parameters. To account for the effects of reconstruction efficiency on the λ parameters, we artificially introduced reconstruction efficiencies into the THERMINATOR data. In this section, we discuss how we used Monte Carlo HIJING events with simulated detector reconstruction to estimate these reconstruction efficiencies.

To investigate the influence of reconstruction efficiencies on particle yields, we looked at the reconstruction rates of the different particle types using the HIJING run specified in Section 2.1.1. The MC-truth yields of the V0s were examined at three stages of reconstruction:

1. In the raw MC event before any detector simulation was added to the MC particles. Here the daughters of the V0s are required to be in mid-rapidity and surpass a minimum p_T value, so as to only look at V0s that could reasonably be found.
2. In the V0 finder before any analysis-side reconstruction cuts were employed.
3. After topological reconstruction cuts were made.

At each stage, the MC truths of all remaining particles of the following types were counted: primary Λ ; secondary Λ from Σ^0 , Σ^* , Ξ^0 , Ξ^- , Ω , and other sources; the respective anti-particles; K_S^0 misidentified as Λ or $\bar{\Lambda}$; other misidentified V0s; and fake V0s reconstructed in the V0 finder. The resulting yields are visible in Figure 3.10.

The results of Figure 3.10 can be better interpreted by taking the ratios of the yields at different stages. Figure 3.11 shows one such ratio. Here, one can see the fraction of Λ of various parentage that were generated in the event that were also recognized as V) candidates by the V0 finder. Note that roughly a third of them are already lost at this stage because of the $\Lambda \rightarrow n\pi^0$ decay channel. One curious feature of this plot is that Λ from

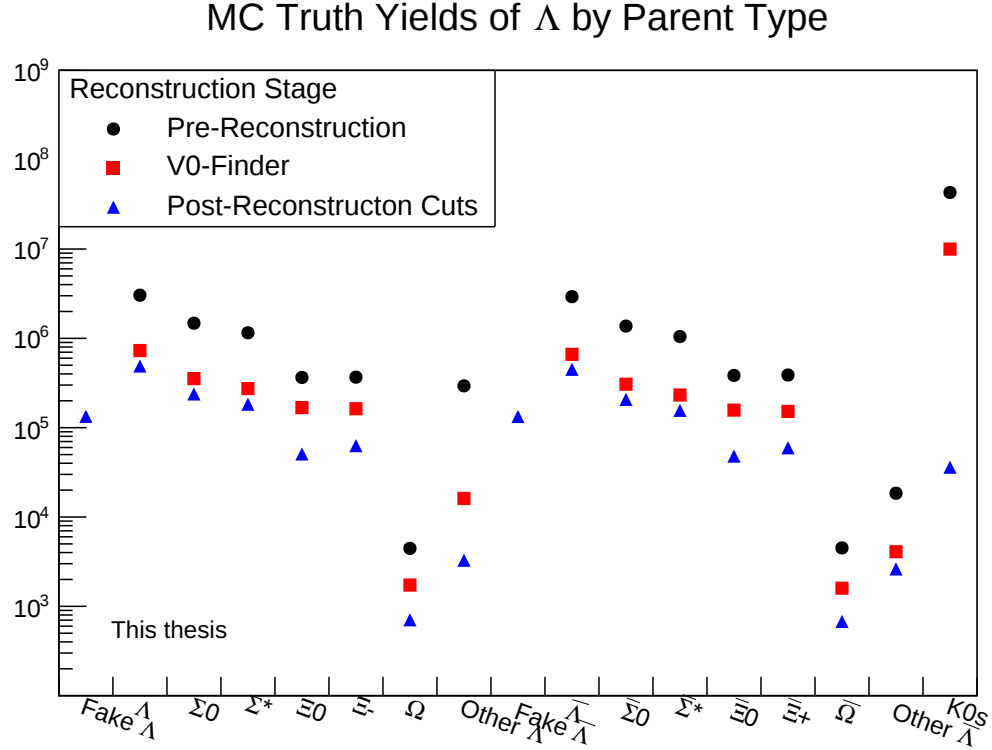


Figure 3.10: MC truth yields of Λ at different stages of reconstruction. The x-axis shows the parentage of the reconstructed particle. The Λ and $\bar{\Lambda}$ bins are primary particle yields. Other bins show the number of Λ ($\bar{\Lambda}$) that come from decays of other particles. The K^0_s bin shows the number of kaons that were misidentified as Λ and $\bar{\Lambda}$. The black points show the yields that exist in the underlying event, the red points show the true parentage of V0-candidates that have been found by the offline V0 finder, and the blue points show yields after all topological reconstruction cuts have been imposed. Only midrapidity particles that exceed a minimum p_T cut have been examined, and injected signals have been excluded from this analysis.

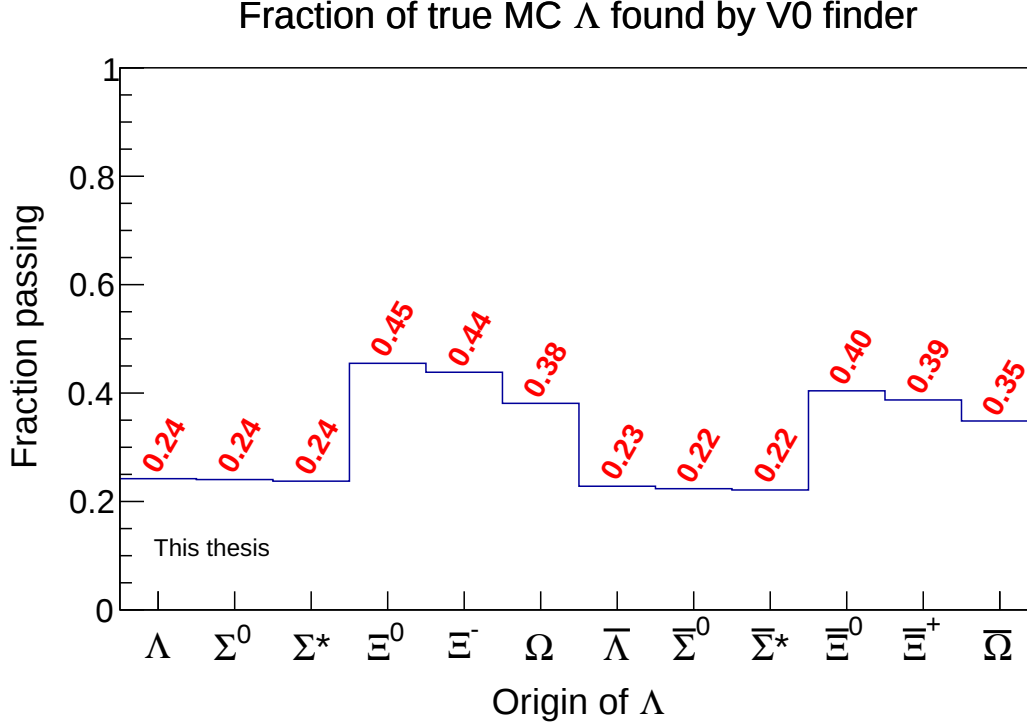


Figure 3.11: Average fraction of true Λ (primary and secondary) in underlying MC events that are labeled as V0 candidates by the V0 finder. The x-axis shows the parentage of the Λ , with the first bin corresponding to primary Λ . About a third of the Λ are lost across the board due to the $\Lambda \rightarrow \pi^0 n$ decay. We note that Λ from weak decays have a much better chance of being found in the V0 finder than primary Λ or those from electromagnetic (Σ^0) and strong (Σ^*) decays. This may be due to the topological cuts in the V0 finder that are designed to exclude primary tracks. Λ from Ξ decays will, on average, have a larger apparent decay length. Their daughter tracks will have a larger average DCA to the primary vertex, so they will be less likely to be ignored as primary tracks.

weak decays have a much better chance of being found than primary Λ . We suspect that this may be due to topological cuts in the V0 finder that have been designed to exclude primary tracks. Λ from weak decays have a larger average decay length than primary Λ , and their daughter tracks have a larger average DCA to the primary vertex. As a result, those daughter tracks are less likely to be ignored as primary tracks.

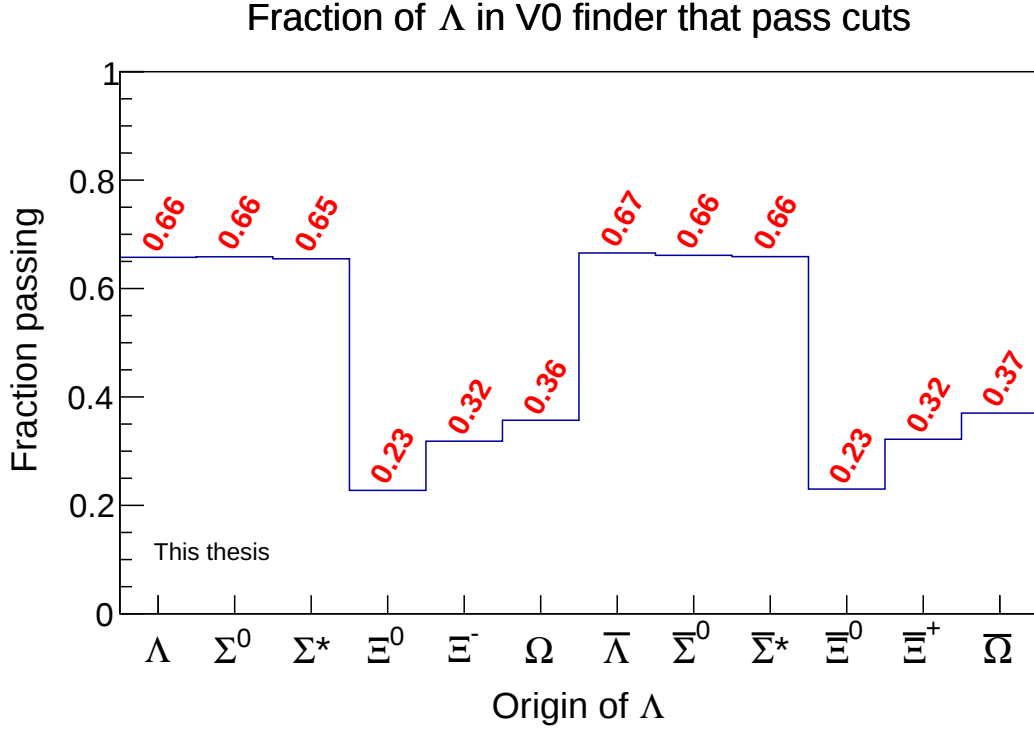


Figure 3.12: Average fraction of true Λ in the V0 finder of MC events that pass all reconstruction cuts. The topological cuts of this analysis (see Section 2.2.1) have been selected to preferentially remove secondary Λ from weak decays. Within the limits of the ALICE tracking resolution, Λ from Σ^0 and resonance decays look like and reconstruct like primary Λ .

Figure 3.12 shows the fraction of V0 candidates in the V0 finder that survived all the reconstruction cuts. From this, one can see our the topological cuts at work — primary Λ , Λ from Σ^0 , and Λ from resonance decays have a much better survival rate than Λ from weak decays.

The final ratio to consider is Figure 3.13, which shows the reconstruction efficiency from the beginning (particles in the underlying event) to the end (all particles passing all

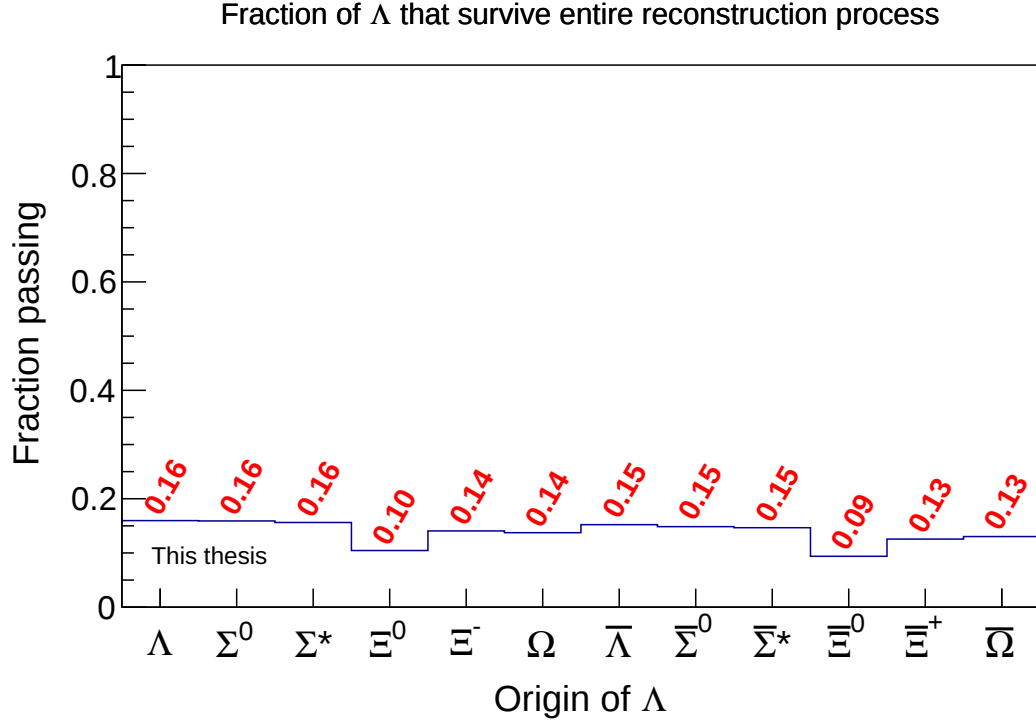


Figure 3.13: Average fraction of true Λ in underlying MC events that pass all reconstruction cuts. The first bin shows the efficiency of reconstructing primary Λ , while the other bins show the reconstruction efficiency of Λ that are daughters of the labeled particles. The results in this plot are the product of the results in Figures 3.11 and 3.12. The reconstruction efficiency of a Λ of any origin is approximately 15%.

cuts). This information, which is just the product of the previous two figures, conveys the overall reconstruction efficiency of Λ of different origins. This efficiency encompasses not only the results of the analysis-side cuts, but also the natural tracking efficiencies of ALICE detector. One can see that Λ of different origins all have approximately the same chance of getting reconstructed — about 15%. When applying the efficiency correction to the yields in Section 3.3.2, we used a flat 15% efficiency for Λ of all origins.

3.3.4 Estimates of λ parameters

Here we present estimates of the λ parameters that describe the relative contributions of feeddown in our analysis. For these results, we have used the model developed in Section 3.3.1, the thermal particle yields estimated in Section 3.3.2, and the reconstruction efficiencies estimated in Section 3.3.3. For ease of reference, our λ model (Eqn. 3.34) is

$$\lambda_i = \frac{\langle N_i \rangle_{\text{event}}}{\langle \sum_i^m N_i \rangle_{\text{event}}}.$$

Figures 3.14–3.16 show the estimated λ parameters for various pair combinations. The elements of the particle-particle matrix are comparable to the elements of the antiparticle-antiparticle matrix, as well as to the diagonal elements of the particle-antiparticle matrix. The sum of the symmetric, off-diagonal elements of the particle-antiparticle matrix (e.g. $\Lambda\bar{\Sigma}^0$ and $\bar{\Lambda}\Sigma^0$) are approximately the same as the respective element ($\Lambda\Sigma$) in the particle-particle matrix. In all three figures, we see that the $\Lambda\Lambda/\bar{\Lambda}\bar{\Lambda}/\Lambda\bar{\Lambda}$ contribution dominates at approximately 25%. The Λ -Y contributions follow between 20% and 10%, except for the $\Lambda\Omega$ pair. The various Ω pairs have a 1% or smaller contribution.

For this analysis, we choose to ignore contributions with an effect smaller than 2%. This excludes Ω pairs entirely, as well as the identical Ξ pairs, $\Xi^0\Xi^0$ and $\Xi^-\Xi^-$. Those pairs, along with pairs containing a fake Λ , are assumed to give no contribution to the total correlation function (i.e. their individual correlation functions are treated as unity for all k^*). This leaves us with the following pairs contributing to the theoretical correlation functions: $\Lambda\Lambda$, $\Lambda\Sigma^0$, $\Lambda\Xi^0$, $\Lambda\Xi^-$, $\Sigma^0\Sigma^0$, $\Sigma^0\Xi^0$, $\Sigma^0\Xi^-$, and $\Xi^0\Xi^-$.

3.4 Momentum resolution corrections

As discussed in Section 2.3.4, we corrected for the momentum resolution of tracks at the fit level. In this section, we examine the effect of momentum resolution on relative momentum distributions, we walk through the math of correcting the theoretical prediction, and we describe how the correction was implemented in practice.

Using a HIJING MC data set, we created a two-dimensional histogram of k_{gen}^* vs k_{recon}^* for pairs of Λ . We removed fake Λ from the analysis. However, because momentum resolu-

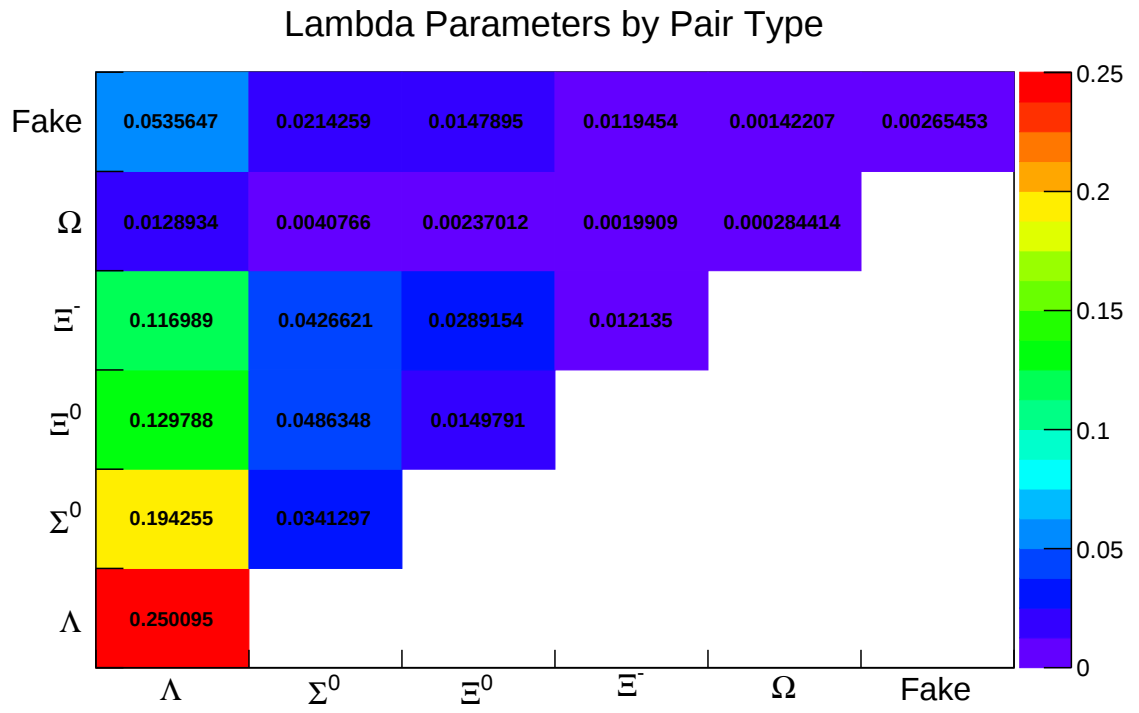


Figure 3.14: Estimated λ parameters for various particle-particle contributions.

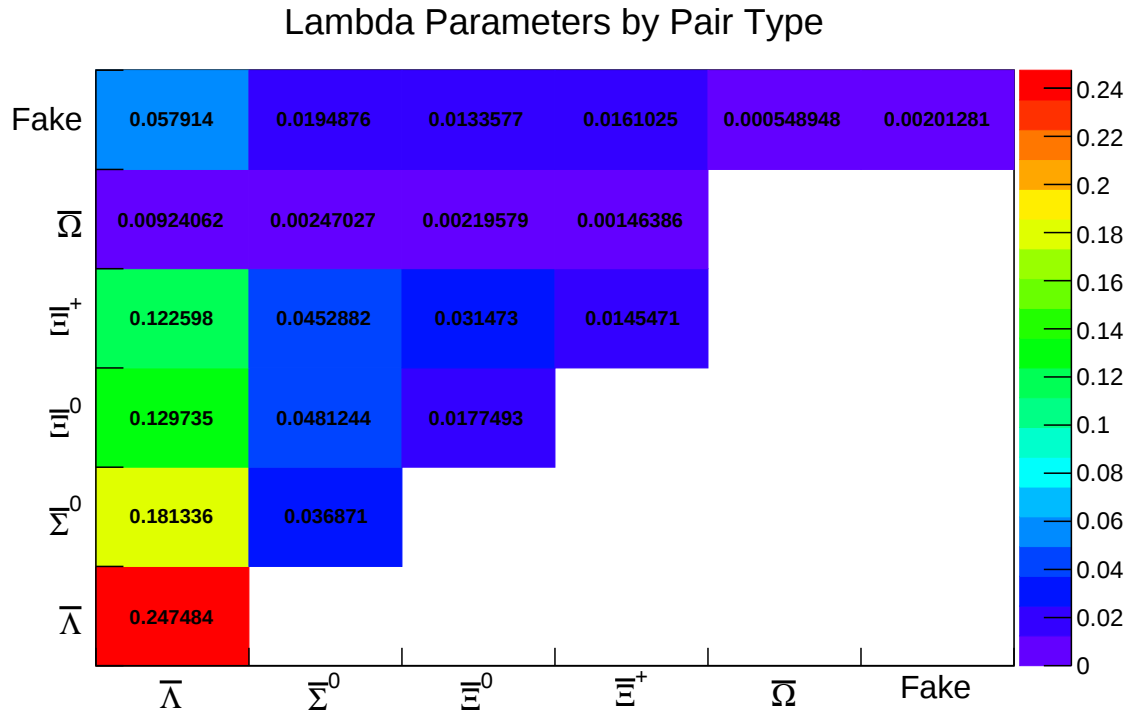


Figure 3.15: Estimated λ parameters for various antiparticle-antiparticle contributions.

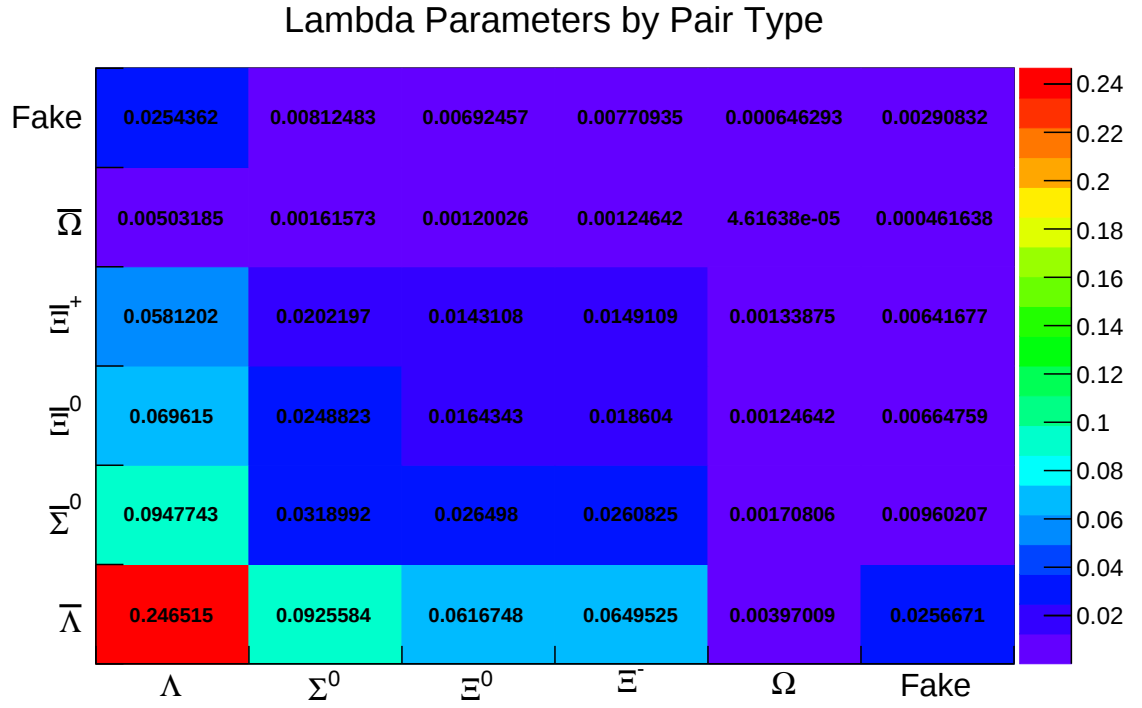


Figure 3.16: Estimated λ parameters for various particle-antiparticle contributions. The sum of the symmetric off-diagonal elements (e.g. $\Lambda\bar{\Sigma}^0$ and $\bar{\Lambda}\Sigma^0$) are approximately the same as the respective element ($\Lambda\Sigma$) in the particle-particle matrix, Figure 3.14.

tion smearing affects the daughters of both primary and secondary Λ , and because we were in need of a large number of pairs, we did not distinguish between primary and secondary Λ . We used mixed event pairs rather than same event pairs to ensure we had sufficient statistics.

Figures 3.17-3.19 show momentum resolution histograms for different pair types. These histograms describe the possible phase space for how relative momentum can be smeared by detector resolution. Each row of each figure is normalized independently of the other rows (see Section 3.4.1). Given a pair with a particular true k^* value, the likelihood of that pair having a particular reconstructed k^* value is given by the weights in the associated bin of the histogram. In general, the momentum shift is marginal. For most pairs, the reconstructed relative momentum differs from the generated relative momentum by no more than a bin or two.

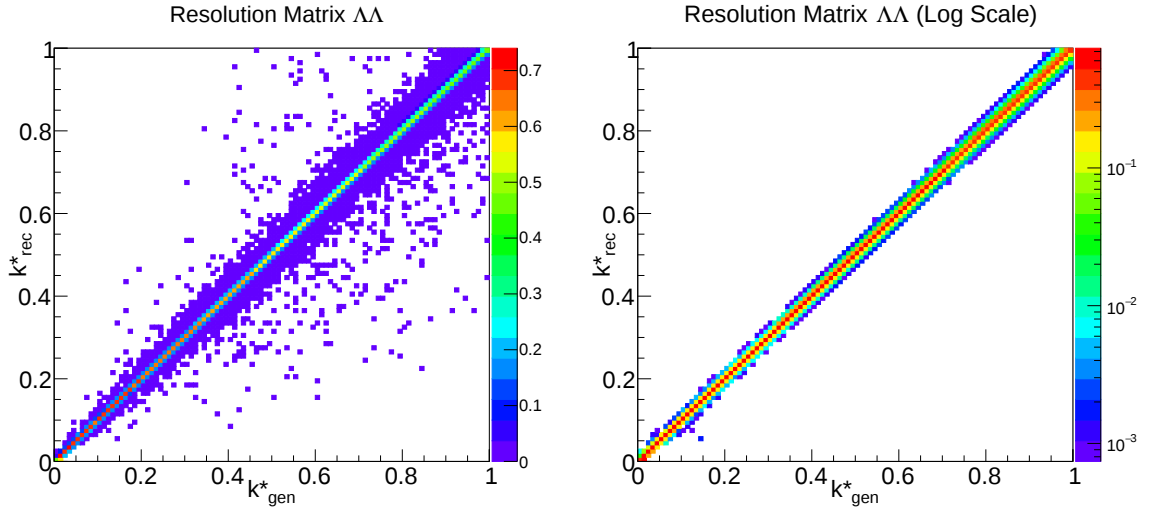


Figure 3.17: Momentum resolution matrices for $\Lambda\Lambda$. In the log scale plot (right), one can see that the resolution smearing is very narrow — most pairs have their k^* shifted by no more than one or two bins.

It is important to note that these histograms describe how a k_{gen}^* distribution can be smeared into a k_{recon}^* distribution, but not the reverse process (unsmearing). Unsmearing would be described by the matrix inverse of these histograms. Unfortunately, attempts to invert the matrices yield significant aliasing, and the results cannot be used.

That said, these momentum resolution matrices can be used to smear the theoretical prediction (i.e. the fit) from k_{gen}^* -space into k_{recon}^* -space using exactly the same procedure

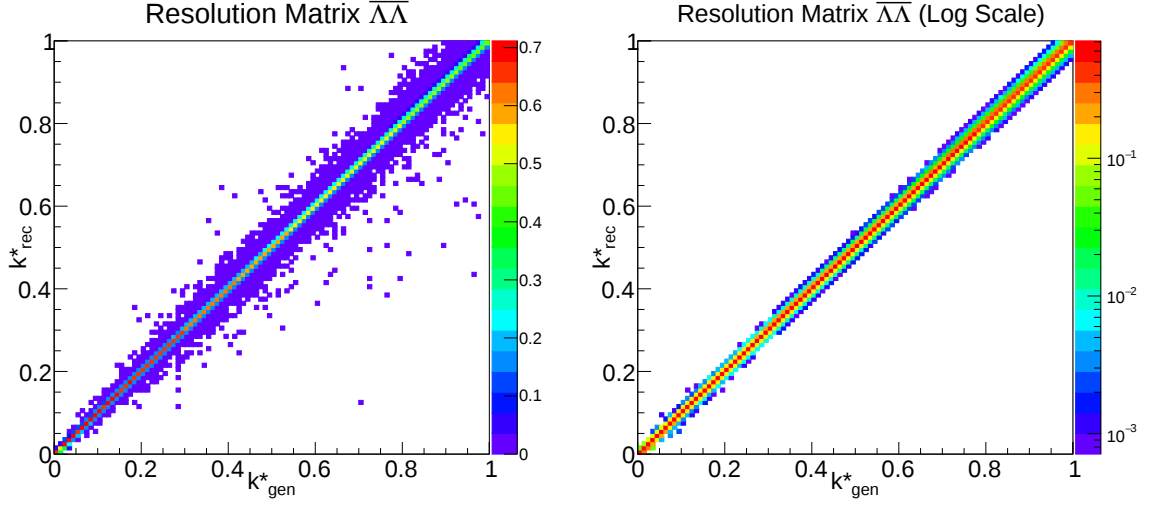


Figure 3.18: Momentum resolution matrices for $\bar{\Lambda}\bar{\Lambda}$. In the log scale plot (right), one can see that the resolution smearing is very narrow — most pairs have their k^* shifted by no more than one or two bins.

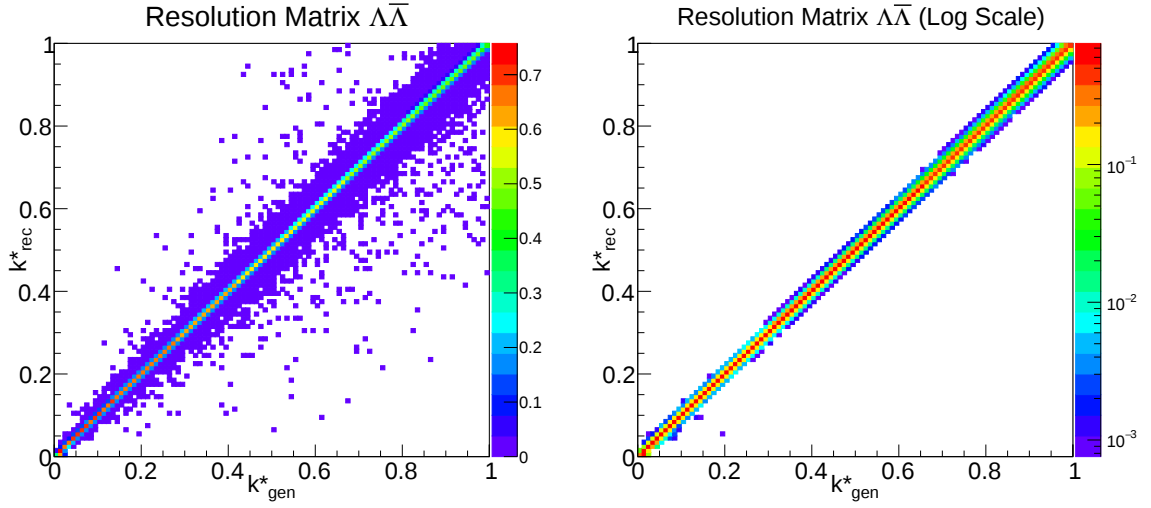


Figure 3.19: Momentum resolution matrices for $\Lambda\bar{\Lambda}$. In the log scale plot (right), one can see that the resolution smearing is very narrow — most pairs have their k^* shifted by no more than one or two bins.

described above for residual correlation smearing:

$$C(k_{\text{rec}}^*) \equiv \frac{\sum_{k_{\text{gen}}^*} M(k_{\text{rec}}^*, k_{\text{gen}}^*) C(k_{\text{gen}}^*)}{\sum_{k_{\text{gen}}^*} M(k_{\text{rec}}^*, k_{\text{gen}}^*)}. \quad (3.35)$$

Here, $M(k_{\text{rec}}^*, k_{\text{gen}}^*)$ is the momentum resolution smearing matrix. One can perform this transformation on the fit function to ensure that both the data and the fit have the same smearing.

There are two main benefits to treating the momentum resolution effects in this way. First, it is no longer necessary to run an iterative analysis (see Section 2.3.4). This method requires a one-time analysis of Monte Carlo data to produce the momentum resolution smearing matrix. Second, the resolution correction can be incorporated directly into the fit procedure as an intermediate step between calculating the fit function and taking the χ^2 difference with the data.

3.4.1 Smear matrices

Momentum resolution affects daughter tracks, and as such the primary or secondary nature of the parent Λ is irrelevant. Therefore, when smearing the fit function, momentum resolution smearing should be applied after residual correlation smearing. In fact, as both types of smearing are effectively forms of matrix algebra, the two smearing effects can be rolled together. In that case, $C(k_{\text{gen}}^*)$ in Eqn. 3.35 is the theoretical correlation function for some pair type expressed in terms of $k_{\Lambda\Lambda}^*$, as given by Eqn. 3.21.

Combined together, they give (in the case of $\Lambda\Sigma$ feeddown)

$$C_{\Lambda\Sigma}(k_{\text{rec}}^{*\Lambda\Lambda}) \equiv \frac{\sum_{k_{\text{gen}}^{*\Lambda\Lambda}} M(k_{\text{rec}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Lambda}) \frac{\sum_{k_{\text{gen}}^{*\Lambda\Sigma}} T(k_{\text{gen}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Sigma}) C_{\Lambda\Sigma}(k_{\text{gen}}^{*\Lambda\Sigma})}{\sum_{k_{\text{gen}}^{*\Lambda\Sigma}} T(k_{\text{gen}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Sigma})}}{\sum_{k_{\text{gen}}^{*\Lambda\Lambda}} M(k_{\text{rec}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Lambda})}. \quad (3.36)$$

This becomes much cleaner if we introduce the partially-normalized (i.e. along one axis) matrices

$$\tilde{M}(k_{\text{rec}}^*, k_{\text{gen}}^*) \equiv \frac{M(k_{\text{rec}}^*, k_{\text{gen}}^*)}{\sum_{k_{\text{gen}}^*} M(k_{\text{rec}}^*, k_{\text{gen}}^*)} \quad (3.37)$$

and

$$\tilde{T}(k_{\Lambda\Lambda}^*, k_{\Lambda\Sigma}^*) \equiv \frac{T(k_{\Lambda\Lambda}^*, k_{\Lambda\Sigma}^*)}{\sum_{k_{\Lambda\Sigma}^*} T(k_{\Lambda\Lambda}^*, k_{\Lambda\Sigma}^*)}. \quad (3.38)$$

Then we get the more compact

$$C_{\Lambda\Sigma}(k_{\text{rec}}^{*\Lambda\Lambda}) = \sum_{k_{\text{gen}}^{*\Lambda\Lambda}} \tilde{M}(k_{\text{rec}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Lambda}) \sum_{k_{\text{gen}}^{*\Lambda\Sigma}} \tilde{T}(k_{\text{gen}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Sigma}) C_{\Lambda\Sigma}(k_{\text{gen}}^{*\Lambda\Sigma}) \quad (3.39)$$

We can contract the $k_{\text{gen}}^{*\Lambda\Lambda}$ indices to get

$$C_{\Lambda\Sigma}(k_{\text{rec}}^{*\Lambda\Lambda}) = \sum_{k_{\text{gen}}^{*\Lambda\Sigma}} S(k_{\text{rec}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Sigma}) C_{\Lambda\Sigma}(k_{\text{gen}}^{*\Lambda\Sigma}), \quad (3.40)$$

where

$$S(k_{\text{rec}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Sigma}) \equiv \sum_{k_{\text{gen}}^{*\Lambda\Lambda}} \tilde{M}(k_{\text{rec}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Lambda}) \tilde{T}(k_{\text{gen}}^{*\Lambda\Lambda}, k_{\text{gen}}^{*\Lambda\Sigma}). \quad (3.41)$$

There are multiple residual correlation transform matrices (one for each pair that decays into $\Lambda\Lambda$), and a corresponding number of full smearing matrices S defined by Eqn. 3.41. For consistency, we also treat the momentum resolution smearing of primary $\Lambda\Lambda$ as a smearing matrix, though here the residual correlation matrix $\tilde{T}$ would be the identity matrix. Figures 3.20–3.35 show the full smear matrices.

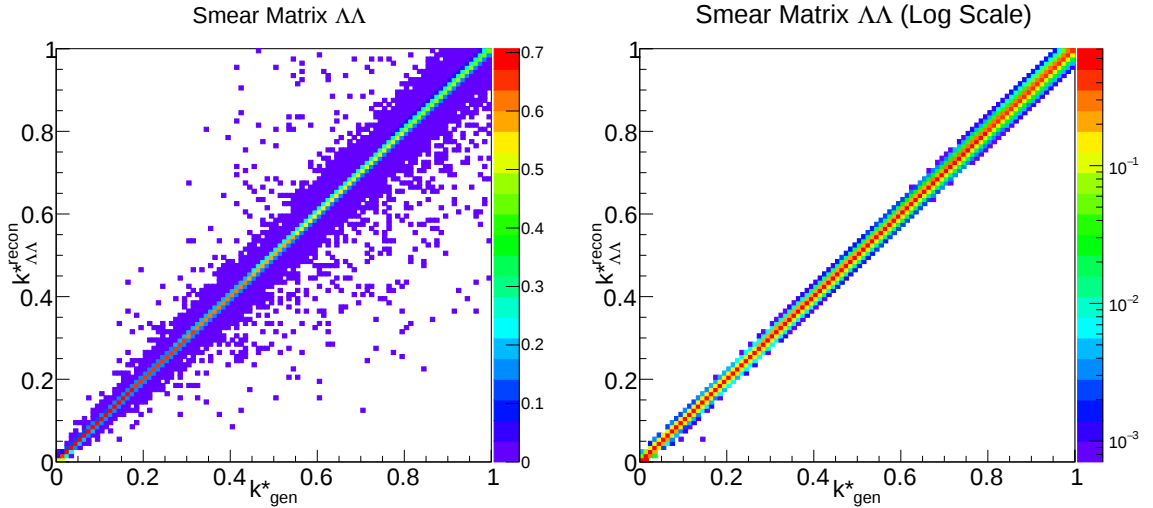


Figure 3.20: Smear matrix used for both $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$. As the residual correlation matrix for primary $\Lambda\Lambda$ pairs is the identity matrix, these plots are identical to the plots in Figure 3.17.

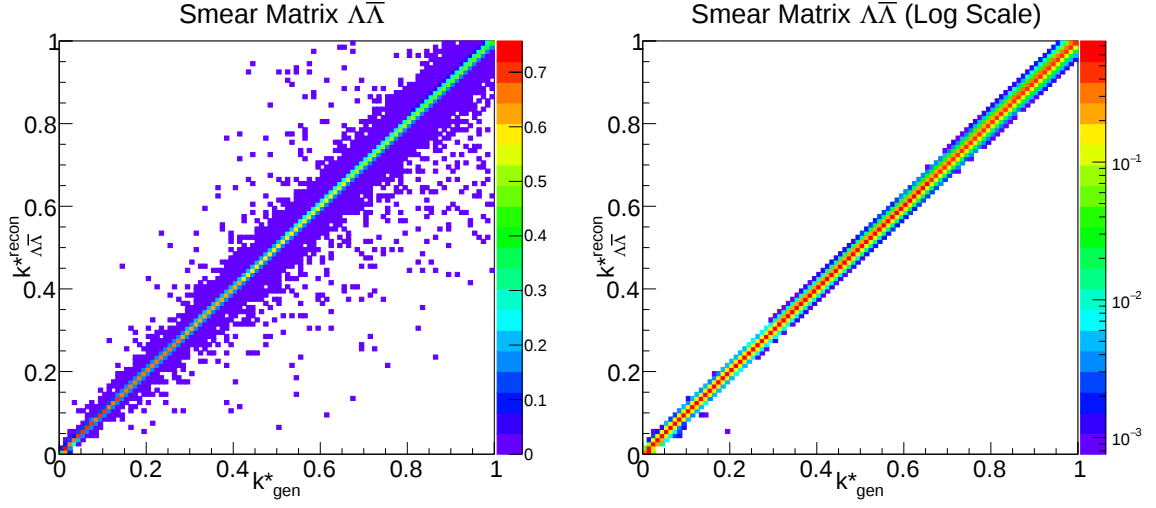


Figure 3.21: Smear matrix for $\Lambda\bar{\Lambda}$. As the residual correlation matrix for primary $\Lambda\bar{\Lambda}$ pairs is the identity matrix, these plots are identical to the plots in Figure 3.19.

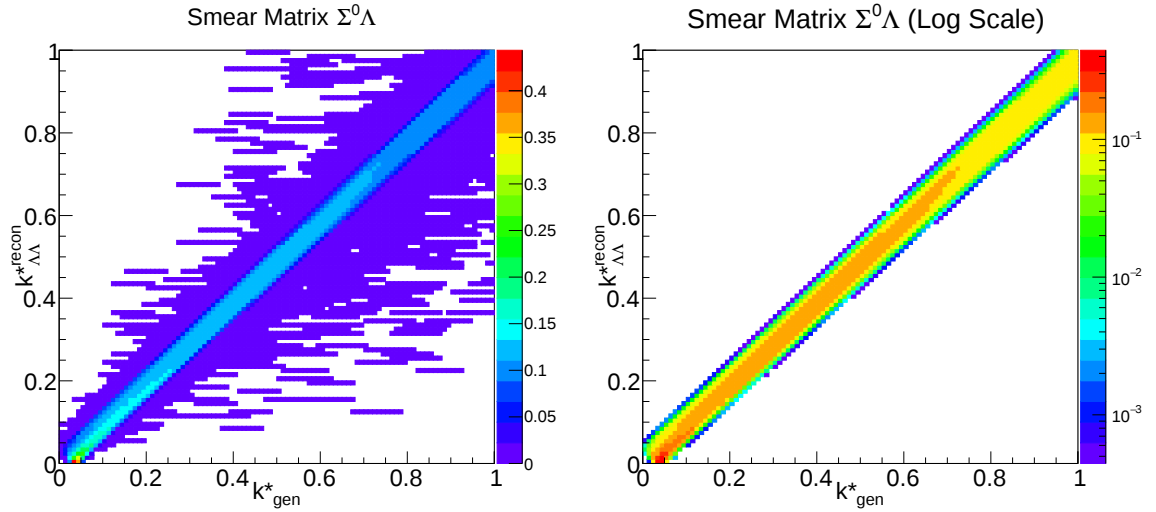


Figure 3.22: Smear matrix for $\Sigma^0\Lambda$ and $\bar{\Sigma}^0\bar{\Lambda}$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

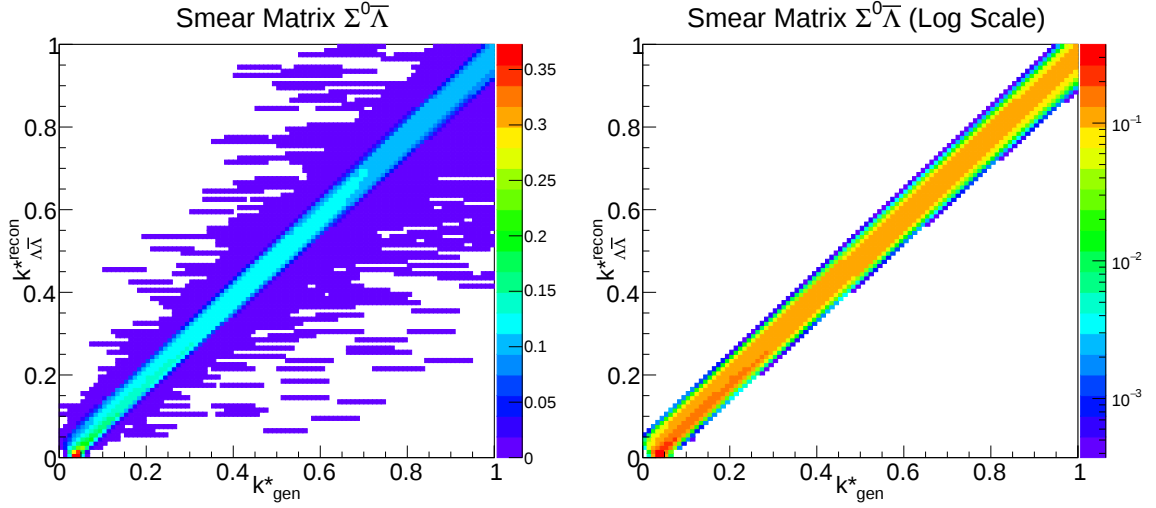


Figure 3.23: Smear matrix for $\Sigma^0\bar{\Lambda}$ and $\bar{\Sigma}^0\Lambda$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

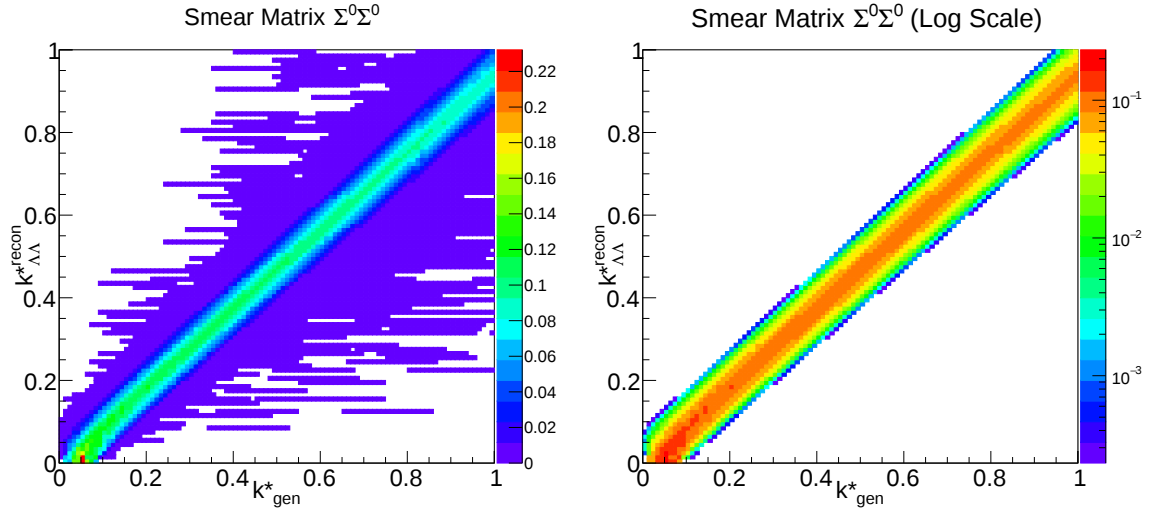


Figure 3.24: Smear matrix for $\Sigma^0\Sigma^0$ and $\bar{\Sigma}^0\bar{\Sigma}^0$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

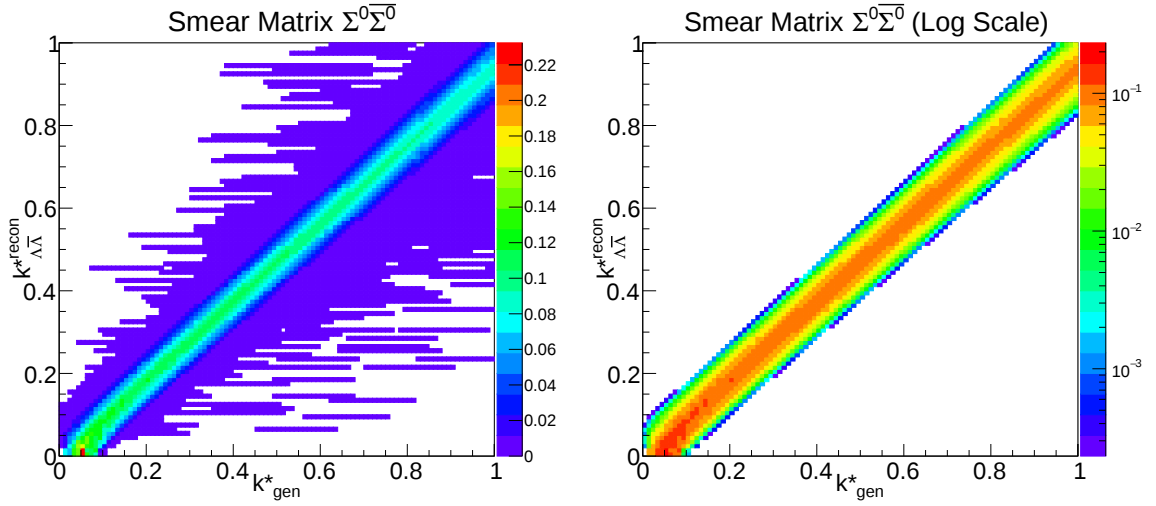


Figure 3.25: Smear matrix for $\Sigma^0\bar{\Sigma}^0$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

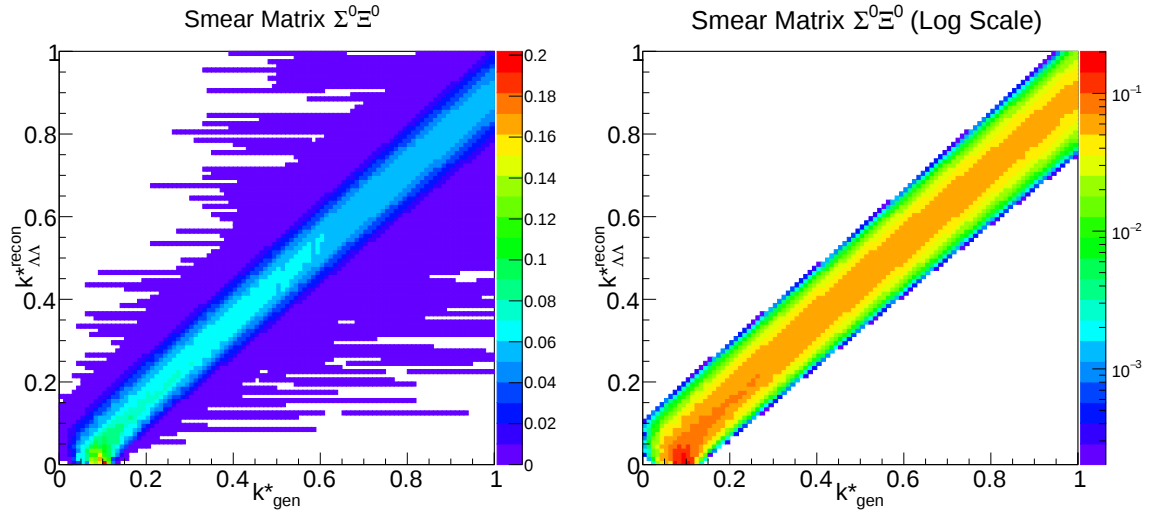


Figure 3.26: Smear matrix for $\Sigma^0\Xi^0$ and $\bar{\Sigma}^0\bar{\Xi}^0$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

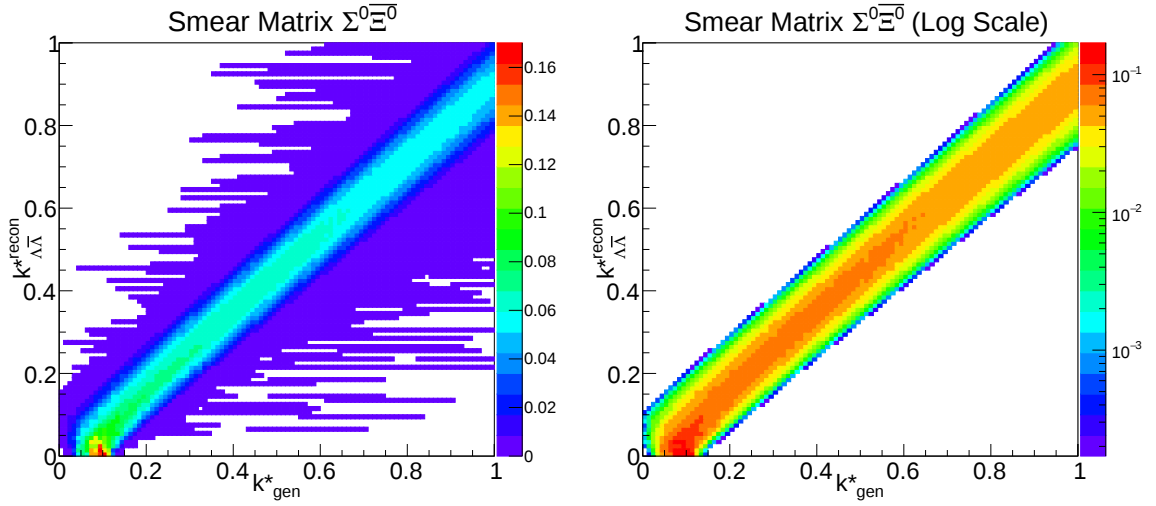


Figure 3.27: Smear matrix for $\Sigma^0\Xi^0$ and $\Xi^0\Xi^0$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

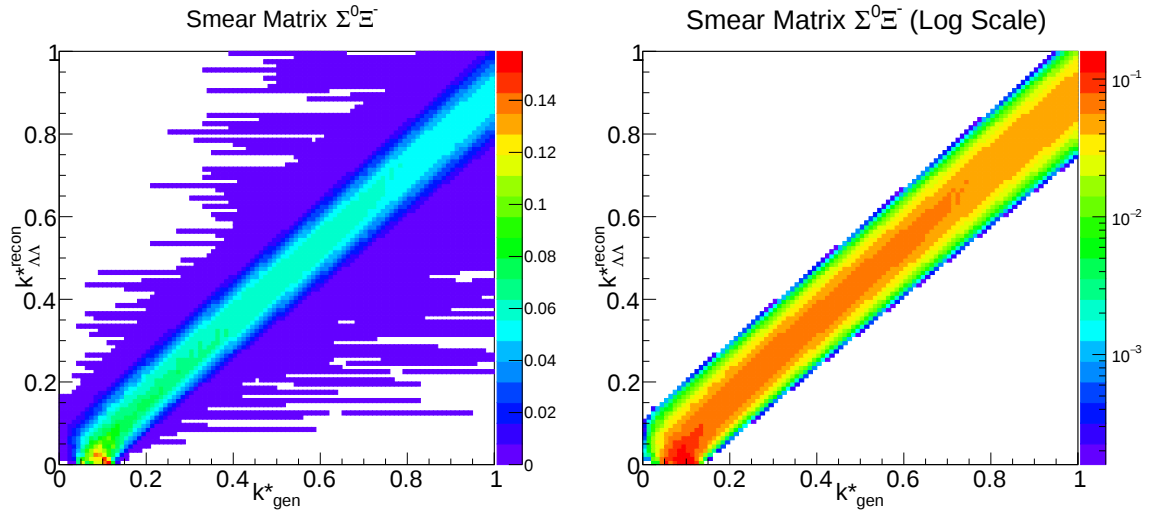


Figure 3.28: Smear matrix for $\Sigma^0\Xi^-$ and $\Xi^0\Xi^+$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

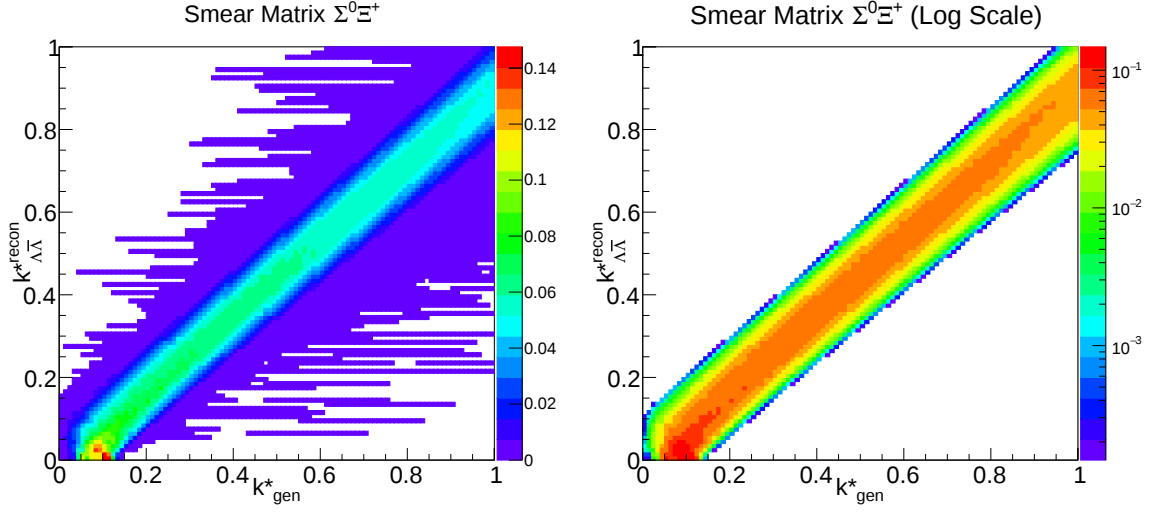


Figure 3.29: Smear matrix for $\Sigma^0\Xi^+$ and $\bar{\Sigma}^0\Xi^-$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

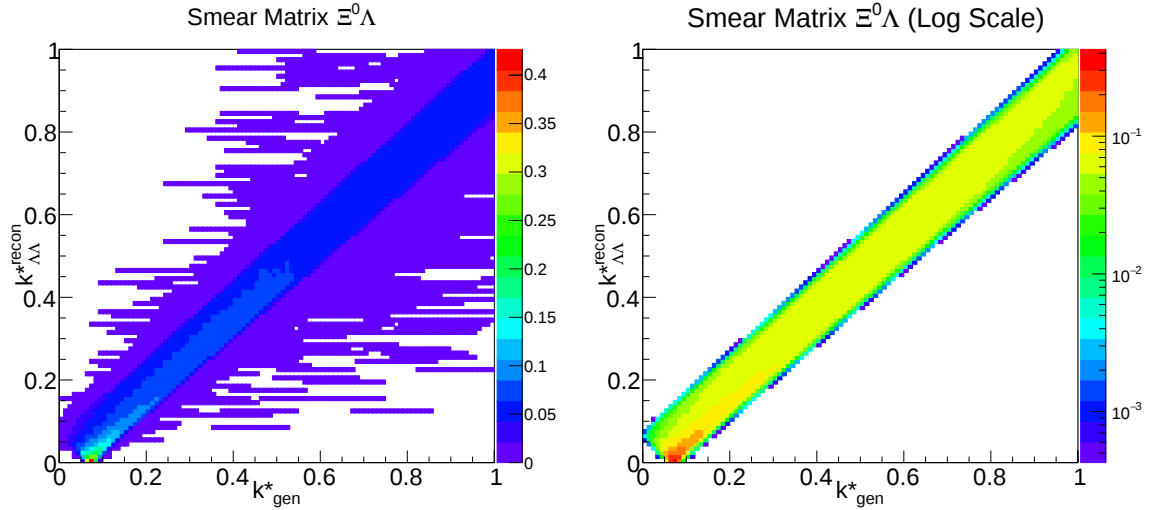


Figure 3.30: Smear matrix for $\Xi^0\Lambda$ and $\bar{\Xi}^0\bar{\Lambda}$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

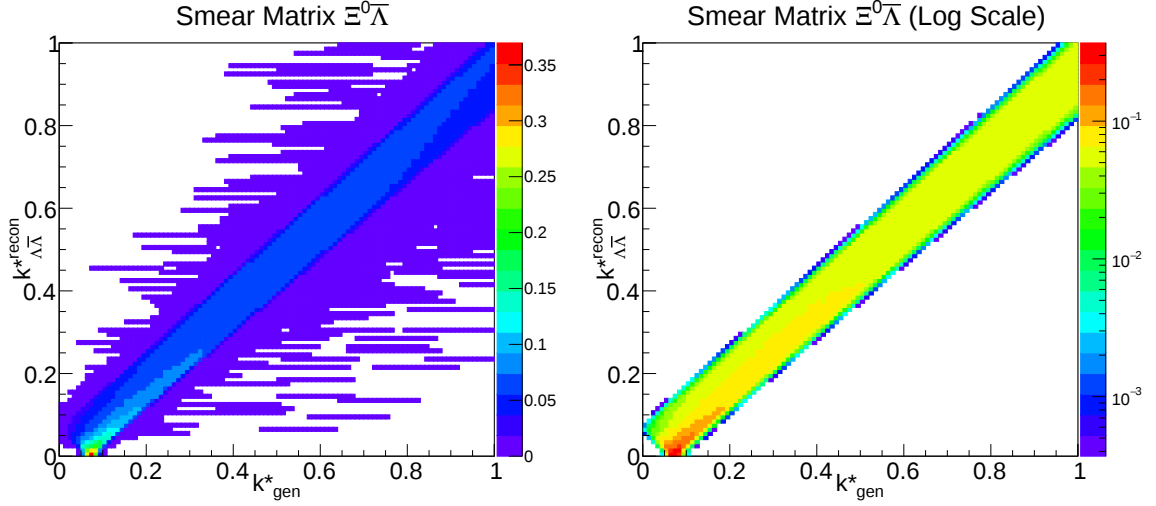


Figure 3.31: Smear matrix for $\Xi^0\bar{\Lambda}$ and $\Xi^0\Lambda$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

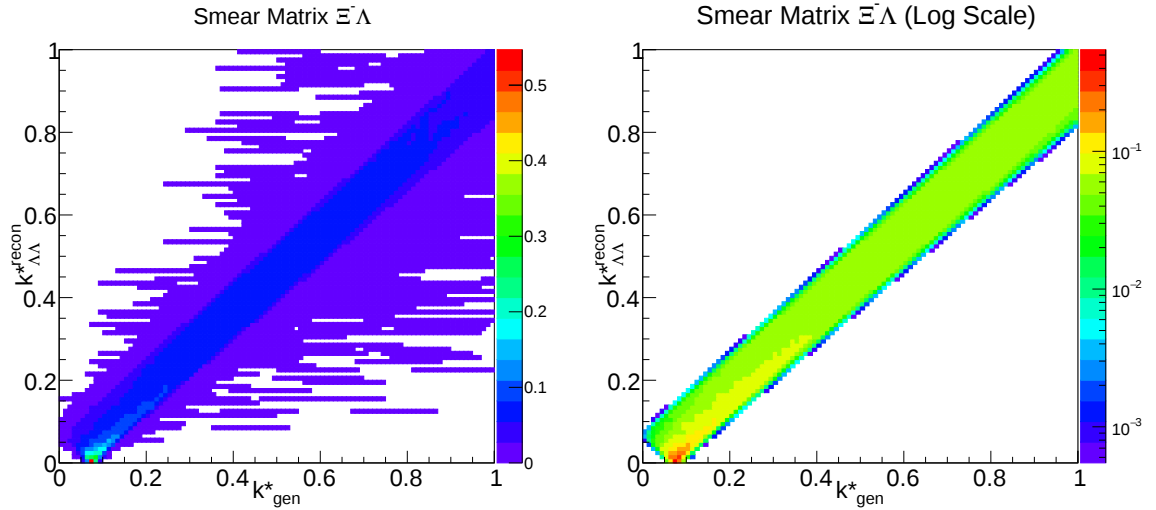


Figure 3.32: Smear matrix for $\Xi^-\Lambda$ and $\Xi^+\bar{\Lambda}$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

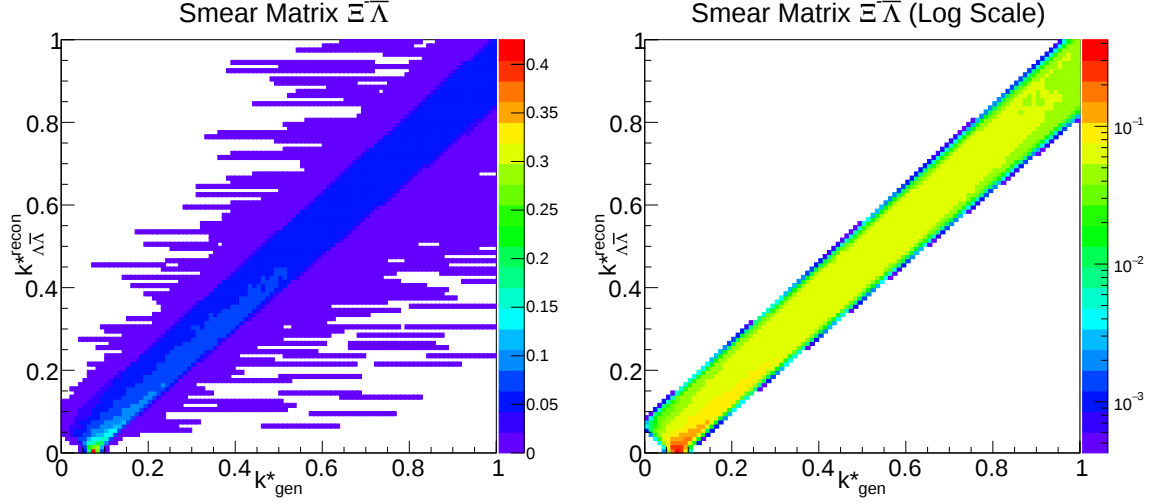


Figure 3.33: Smear matrix for $\Xi^-\bar{\Lambda}$ and $\Xi^+\Lambda$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

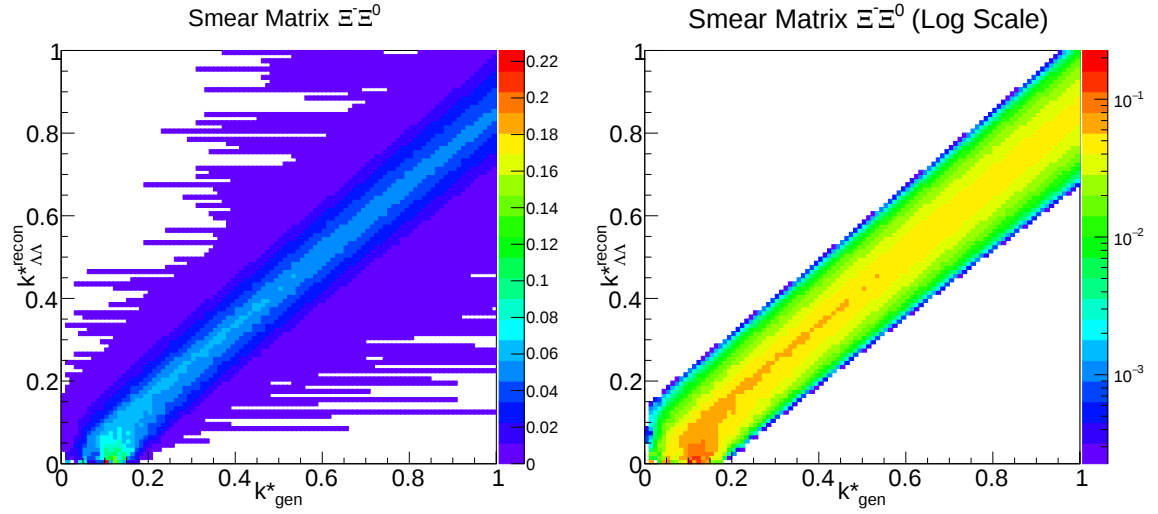


Figure 3.34: Smear matrix for $\Xi^-\Xi^0$ and $\Xi^+\bar{\Xi}^0$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

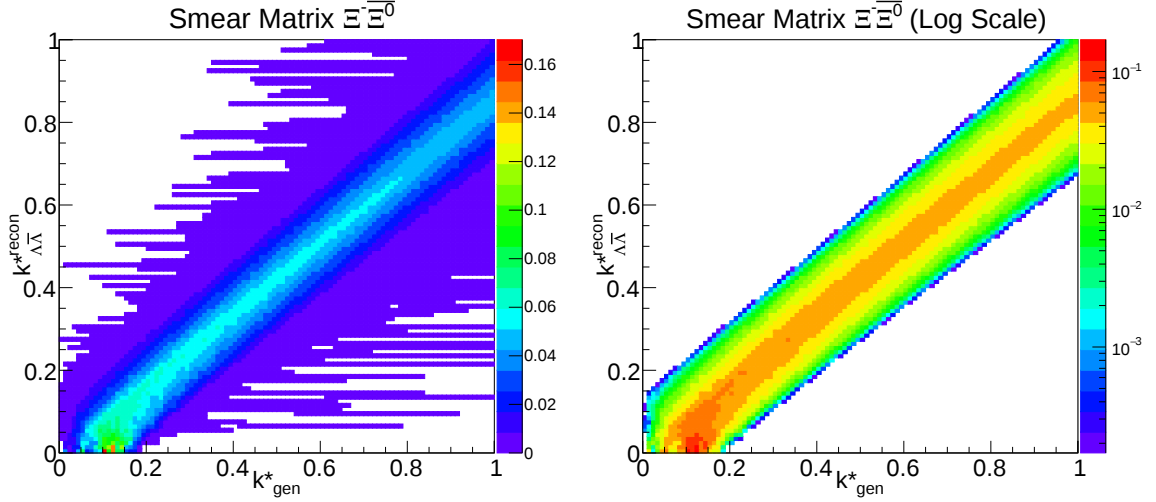


Figure 3.35: Smear matrix for $\Xi^-\Xi^0$ and $\Xi^+\Xi^0$, which accounts for both residual correlation and momentum resolution smearing. This matrix can be used to correct the theoretical correlation function so that the fit has the same smearing as the data.

3.4.2 Implementation in fit

When calculating the fit function, the contribution from the primary correlation function $C_{\Lambda\Lambda}(k_{\Lambda\Lambda}^*)$ and each residual correlation ($C_{ij}(k_{ij}^*)$) is first calculated according to Eqn. 3.18. Then each contribution is smeared with its own smear matrix via Eqn. 3.40. Finally, each contribution is combined using the λ weights according to Eqn. 3.20.

3.5 Fit procedure

In this section, we discuss the procedure by which we fit the data. We use a self-developed software package called FemtoFitter. At its heart, FemtoFitter utilizes the ROOT TMinuit implementation of the MINUIT [65] fitting package. In the following sections, we'll describe the functionality of MINUIT and how it fits into FemtoFitter, we'll discuss our usage of fit parameters, and we'll elaborate on general fitting details such as fit range and simultaneously fitting multiple correlation functions.

3.5.1 MINUIT

The MINUIT tool is designed to minimize the value of a multivariable function and to map out the shape of that function around the minimum. It is typically used to minimize a chi-square or log-likelihood function in order to find the best fit for and statistical uncertainty of the input parameters.

When using chi-squared minimization to fit the data, MINUIT varies the free parameters to minimize the function

$$\chi^2(\alpha, \beta, \dots) = \sum_{i=1}^n \frac{(f(x_i; \alpha, \beta, \dots) - m_i)^2}{\sigma_i^2}, \quad (3.42)$$

where α , β , etc. are the input parameters, m_i is a measurement with uncertainty σ_i , and $f(x_i; \alpha, \beta, \dots)$ is a functional prediction of the data that depends on the input parameters. For a femtoscopic analysis, the measured data m_i are the bin values of the combined correlation function described by Eqn. 2.4, and the summation is done over the $x_i = k^*$ bins. The function $f(x_i; \alpha, \beta, \dots)$ is the predicted femtoscopic correlation function given by Eqn. 3.19, and the variable parameters α , β , etc. include the scattering length f_0 , effective range d_0 , and femtoscopic radius R . Chi-square minimization is not necessarily optimal for fitting correlation functions, since the ratio of two Poisson distributions is not a Poisson distribution [13]. As a result, the statistical uncertainties generated via the chi-square method are less accurate than those from the log-likelihood method.

In the alternative log-likelihood method, one does not use the measured correlation function as the data. Instead, one utilizes the unscaled, Poisson-distributed numerator and denominator distributions $A(k^*)$ and $B(k^*)$. For a correlation function analysis, the principal of maximum likelihood (PML) function is given [66] by

$$\chi_{\text{PML}}^2 = \sum_{i=1}^n -2 \left[A_i \ln \left(\frac{f(x_i)(A_i + B_i)}{A_i(f(x_i) + 1)} \right) + B_i \ln \left(\frac{(A_i + B_i)}{B_i(f(x_i) + 1)} \right) \right]. \quad (3.43)$$

Here, $f(x_i)$ is the same theoretical prediction as in Eqn. 3.42. Since we use the unscaled numerator and denominator distributions rather than correlation functions, we cannot feed in a combined correlation function as we did in Eqn. 3.42. Instead, to combine data sets for an improved fit, we calculate $\chi_{\text{PML, total}}^2$ as the sum of the values found by applying Eqn. 3.43 to each of those data sets separately. As a side note, while log-likelihood minimization uses similar nomenclature to chi-square minimization (χ_{PML}^2 and χ^2), they represent different phenomena. Therefore chi-square heuristics, such as $\chi^2/\text{ndf} \sim 1$ being a decent fit, do not apply.

In this analysis, we tried fitting via both methods. Both yielded approximately the same results, but the log-likelihood minimization took about an order of magnitude longer to converge. Due to the run-time issues, we opted to use the chi-square method when making the systematic uncertainty estimates and measuring the final results.

3.5.2 FemtoFitter

The FemtoFitter package has several classes: Fitter, PairSystem, and LednickEqn. The Fitter class sets up runs the TMinuit object; accepts user inputs such the data to be fit, fit

parameter information, and fit range; and saves the results.

A PairSystem object is created for each data set that is being fit, and the measured correlation function for that system is passed in. For example, the measured $\Lambda\Lambda$ 0–10% centrality correlation function would get a PairSystem, as would $\Lambda\bar{\Lambda}$ 0–10%. When the Fitter receives new fit parameters from the TMinuit object, it passes the relevant femtoscopic radius R and scattering parameters f_0 and d_0 to each PairSystem. The PairSystem, in turn, passes those scattering parameters on to its set of LednickyEqn objects. The PairSystem has one LednickyEqn for the primary-primary correlation contribution and one for each residual correlation.

Each LednickyEqn object is given a name (e.g. “LambdaSigma”), a λ parameter (see Section 3.3.4 for values), a smear matrix (see Section 3.2.1), and a bool that denotes whether the pair is composed of identical particles. The LednickyEqn uses that information, along with the scattering parameters that it receives from the PairSystem, to calculate the value of the Lednicky and Lyuboshits equation (Eqn. 3.16) for a set of k^* values. The LednickyEqn then uses its smear matrix to transform that theoretical correlation function into the $k_{\Lambda\Lambda}^*$ basis.

To determine the χ^2 value of a fit, the PairSystem takes the smeared Lednicky equation from each of its LednickyEqn objects, combines those into a correlation function according to Eqn. 3.19, and calculates χ^2 according to Eqn. 3.42 or 3.43. The PairSystem passes that on to TMinuit by way of the Fitter, TMinuit uses that information to determine how to continue to vary the parameters, and the process repeats.

3.5.3 Fit parameters

We use the following parameters in our fit: femtoscopic radius R , real and imaginary part of the scattering length f_0 , effective range of interaction d_0 and a normalization constant. R , f_0 , and d_0 have already been described in preceding sections. A multiplicative normalization constant n is introduced to the theoretical correlation function via $C'(k^*) = (k^*)/n$ to help match the height of the background. In χ^2 fitting, the correlation function (data) is already normalized to one in the mid- k^* region, so this amounts to a fine-tuning with n usually in the range $0.98 < n < 1.01$. For log-likelihood fitting, the numerator and denominator distributions are not pre-normalized, so the normalization parameter falls in the range of $5 < n < 10$.

Each of the listed parameters can be left as a free parameter of the fit or fixed to a certain value. For the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$, $\Im f_0$ is fixed to 0, as there is not expected to be an imaginary scattering component for baryon-baryon scattering. Other parameters can be fixed for exploratory fitting or testing purposes.

The FemtoFitter can be used to fit a single PairSystem by itself, or it can fit multiple PairSystems at the same type and constrain them to share some of their parameters. For

example, we do a joint fit of $\Lambda\bar{\Lambda}$ 0–10%, 10–30%, and 30–50% centrality. We expect each of those systems to have a different femtoscopic radius R , so we introduce 3 radius parameters. In contrast, we use a single f_0 and a single d_0 to describe the nuclear scattering, since the nuclear interaction should be independent of event centrality.

3.5.4 Non-femtoscopic background

Equation 3.18 predicts a flat background at moderate k^* (~ 0.3 GeV) and above, but the measured correlation functions are decidedly non-unitary in that region. This may be due to flow, pair-production in minijets, and other non-femtoscopic effects that are not included in the model. This non-flat background at large k^* gives concern that the background at low k^* may not be flat either. Unfortunately, we do not know if the slope of the low- k^* non-femtoscopic effects is positive, negative or zero. Some femtoscopic analyses examine MC models to characterize the background. Unfortunately, the HIJING model fails to even qualitatively reproduce the large k^* region of these correlation functions. As we do not know the background of the low k^* region, we must take a guess at some sort of model to describe it. In this analysis, we assume that the background in the low k^* region follows the trend of the high k^* background. We therefore introduce into our fit a polynomial factor to correct for the background:

$$C(k^*)' = C(k^*) * (1 + ak^* + bk^{*2}). \quad (3.44)$$

We use this to quadratic form fit the correlation function in the region $k^* < 0.4\text{GeV}/c$.

In the high k^* region of $0.6\text{ GeV}/c < k^* < 0.9\text{ GeV}/c$, we instead fit with

$$C(k^*)'' = (1 + ak^* + bk^{*2}). \quad (3.45)$$

While the Lednicky equation should be unity in this region for reasonable scattering parameters, we fix it to unity to safeguard against the fitter trying to use unusual f_0 or d_0 values to over-fit the background.

3.5.5 Scattering parameters of residual pairs

So far, we have discussed the scattering parameters f_0 and d_0 for $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$ scattering, but we have omitted discussion of the scattering of the residual correlation pairs such as $\Lambda\Sigma$ or $\Sigma\Sigma$. In an ideal world, we would look up these values in a table of hadron-hadron interactions, plug those values into the respective Lednicky equations, and only have to contend with an unknown femtoscopic radius for these pairs. Unfortunately, little to nothing is known about the scattering lengths of these interactions. Therefore, we must find some other way to account for these residual correlation effects. One option would be to have f_0 , d_0 , and R as independent free parameters for each type of pair (e.g. $f_{0,\Sigma\Lambda}$,

$f_{0,\Sigma\Sigma}$, $f_{0,\Xi-\Lambda}$) and extract those values from the fit. Alas, this introduces more than a score of new parameters, and it would lead to far too unconstrained of a fit to extract anything meaningful. Another option—the one we utilized in this analysis—is to make the simplifying assumption that those pairs have the same scattering parameters and radius as the primary-primary pair. In that case, each `LednickEqn` belonging to a `PairSystem` will get passed the same parameters, and they will differ only by the smear matrices used to transform them.

3.6 Systematic uncertainties of fits

We have two types of systematic uncertainties that play into our fit results: uncertainties coming from our choice of topological and pair cuts and uncertainties related to our fit procedure. In the following sections, we'll discuss how we quantify those uncertainties.

3.6.1 Systematic uncertainties from topological and pair cuts

As discussed in Section 2.4, there are systematic uncertainties in our results arising from ambiguities in our choice of topological reconstruction cuts and pair cuts. We analyzed how varying those cuts varied the correlation functions, which resulted in systematic error bars on the correlation functions. We do not directly include the correlation function systematic error bars in our fits. Instead, we fit the correlation functions for each cut configuration and then see how the fit results vary.

As discussed earlier, for each type of cut we construct correlation functions using the nominal cut value as well as using $\pm 10\%$ of the cut value. We then fit each of these correlation functions and record the measured fit parameters. For each type of cut and each parameter, we find the standard deviation of the three measurements. Then we calculate the rms value of those standard deviations, and quote that as the systematic error. The results can be found in Table 3.1.

3.6.2 Systematic uncertainties from fit method

We have examined various aspects of the fit process for sources of systematic uncertainty. We varied the fit range, changed the order of the polynomial background fit, and explored different options of how to do joint fitting of multiple correlation functions.

Our exploratory analysis of joint correlation function fitting didn't contribute to the systematic error bars, but it did help us to determine that we shouldn't including the 30-50% centrality bin of $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ in the overall fit. Rather than increasing the statistical significance of the fit results, including that bin drastically changed the measured scattering parameters from what was seen with only the first two centrality bins. For example, the effective range went from $d_{0,\Lambda\Lambda} \sim 5$ fm to $d_{0,\Lambda\Lambda} \sim 500$ fm. In addition, we observed that

separate fits to the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$ radii for a given centrality bin returned radii that differed by about a factor of two ($R_{\Lambda\bar{\Lambda},0-10\%} = 2.9 \pm 0.3$ fm vs $R_{\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda},0-10\%} = 5.4 \pm 0.6$ fm). As one might expect, by forcing the two pair types to use the same radius for a given centrality bin, we found radii roughly halfway between the separate results, also with a reasonable χ^2 . There are arguments why those radii should be the same, such as how there should be no asymmetry between the Λ and $\bar{\Lambda}$ single-particle emission functions at LHC energies. Nonetheless, the separate results differed by such a statistically significant amount that we decided to remove the same-radius constraint and let the data tell their own story.

To model the non-femtoscopic background we used a polynomial fit of the high k^* region (see Section 3.5.4 for more details), which had the general form

$$f_{\text{back}}(k^*) = (1 + ak^* + bk^{*2}). \quad (3.46)$$

As the background fit is purely phenomenological, we experimented with several configurations of fit parameters. In addition to the quadratic fit above, we also fit with a linear function ($b = 0$) and with the background being unity ($b = a = 0$). For all three options, the fitter was able to converge to reasonable looking values. Because of the visibly downward trend of the background at large k^* , these three parameter configurations resulted in qualitatively different backgrounds in the low k^* region. In general, the quadratic fits resulted in the backgrounds going down below unity, and the linear fits made the backgrounds go up above unity.

We explored variations of the upper bound of the fit range: 0-300 MeV, 0-400 MeV, 0-500 MeV. There were no drastic differences to the fit results in any of the fit ranges.

The systematic uncertainty for each fit parameter was calculated as the standard deviation of the measured values. Changing the background fitting method contributed a larger portion of the systematic uncertainties ($\sim 80\%$), while varying the fit range had a smaller effect.

3.7 Fit results

Table 3.1 shows the measured values for the extracted radii and scattering parameters, along with statistical and systematic error bars. The “Fit Sys” and “Cut Sys” columns list the systematic uncertainties arising from variation in the fit method and in the topological and pair cut values, respectively. In almost all cases, the cut systematics are an order of magnitude smaller than the fit method systematics. The “Total Sys” column shows the fit method and cut systematics added in quadrature.

Parameter	Value (fm)	$\pm$ Stat	$\pm$ Total Sys	Fit Sys	Cut Sys
Radius $\Lambda\bar{\Lambda}$ 0-10%	2.9	0.3	0.3	0.3	0.02
Radius $\Lambda\bar{\Lambda}$ 10-30%	2.3	0.2	0.2	0.2	0.02
Radius $\Lambda\bar{\Lambda}$ 30-50%	1.8	0.2	0.2	0.2	0.01
Radius $\Lambda\Lambda$ 0-10%	5.4	0.6	0.2	0.2	0.1
Radius $\Lambda\Lambda$ 10-30%	4.2	0.5	0.2	0.2	0.04
$\Re f_0$ $\Lambda\bar{\Lambda}$	-0.5	0.1	0.1	0.1	0.01
$\Im f_0$ $\Lambda\bar{\Lambda}$	0.14	0.09	0.02	0.02	0.006
$\Re f_0$ $\Lambda\Lambda$	-0.6	0.3	0.05	0.04	0.02
d_0 $\Lambda\bar{\Lambda}$	1.7	0.2	0.2	0.2	0.01
d_0 $\Lambda\Lambda$	5.3	2.7	0.3	0.2	0.2

Table 3.1: Measured fit parameters with statistical and systematic errors. All values are measured in fm. The Fit Sys and Cut Sys columns show the systematic uncertainties arising from variations in the fit method and the topological and pair cut parameters, respectively. The Total Sys column comes from the fit and cut systematic errors added in quadrature.

3.7.1 $\Re f_0$

Figure 3.36 shows the measured $\Re f_0$ for $\Lambda\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$. A number of other measurements are included for comparison. From the left: $\Lambda\bar{\Lambda}$ and $\Lambda\bar{\Lambda} \oplus \bar{\Lambda}\bar{\Lambda}$ scattering lengths from this analysis, $\Lambda\bar{\Lambda} \oplus \bar{\Lambda}\bar{\Lambda}$ from Au–Au $\sqrt{s_{\text{NN}}} = 200$ GeV/c measured at STAR [49], $\Lambda\bar{\Lambda}$ result from a 4-body cluster model of hypernuclear interactions [37] applied to a measurement of ${}^6_{\Lambda\Lambda}\text{He}$ [48], $\Lambda\bar{\Lambda}$ result from the Nijmegen D soft-core interaction model [38] (also applied to [37]), $p\bar{\Lambda}$ extracted from a residual correlation analysis [46] of STAR data, $p\Lambda$ singlet and triplet scattering lengths [67], np singlet and triplet scattering lengths [62], and the nn scattering length [68].

The np triplet spin state is a known bound state — the ground state of the deuteron. Our measurement $\Lambda\bar{\Lambda}$ scattering length (~ -0.6 fm) is about an order of magnitude smaller than the deuteron scattering length. Since the deuteron is only weakly bound, the significantly smaller scattering length of $\Lambda\bar{\Lambda}$ suggests that it does not have a dibaryon bound state.

Figure 3.37 shows $\Re f_0$ again, but here it is zoomed in to better highlight the details of the various $\Lambda\bar{\Lambda}$ measurements. There are several features to note here. First, our measured real scattering length for $\Lambda\bar{\Lambda} \oplus \bar{\Lambda}\bar{\Lambda}$ is consistent with that for $\Lambda\bar{\Lambda}$. This suggests that their behavior under elastic scattering is the same; only the annihilation of the $\Lambda\bar{\Lambda}$ pair (accounted for by the imaginary part of the scattering length) sets them apart.

Second, we can see that the STAR measurement of $\Lambda\bar{\Lambda}$ agrees with our measurement within their systematic uncertainties. That said, neither our measurement nor STAR’s agrees with the results from the hypernuclear interaction models. While they are of a similar magnitude, the hypernuclear scattering results have a positive sign, suggesting an

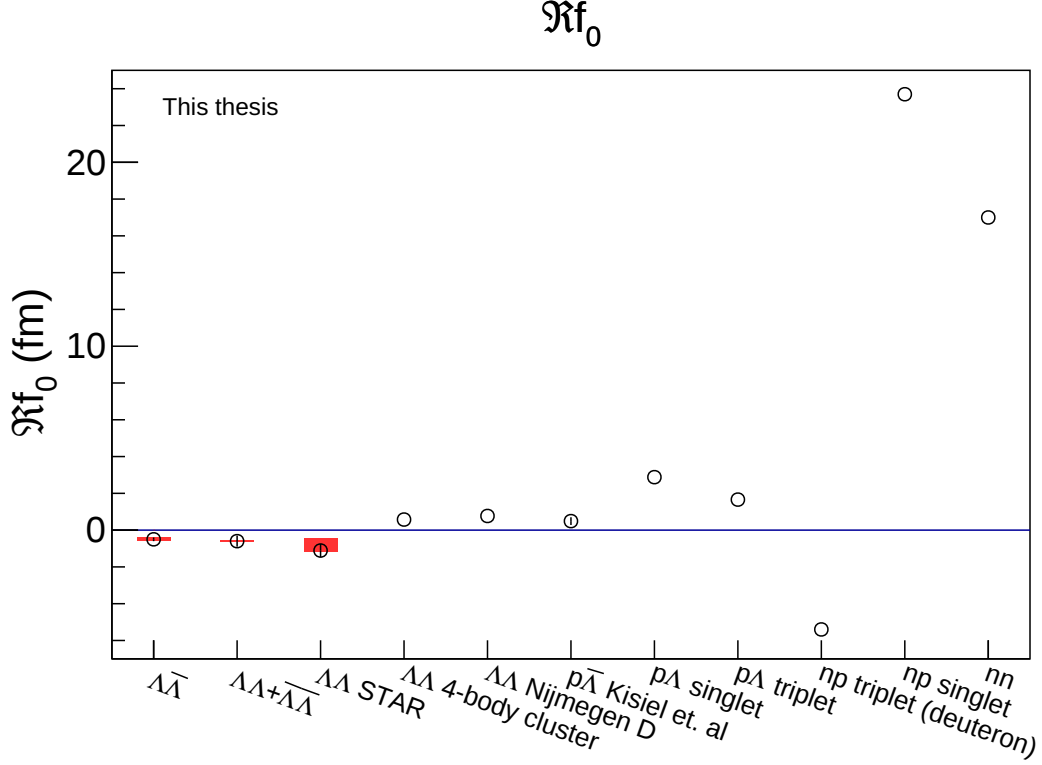


Figure 3.36: Several measurements of the real part of the $\Lambda\Lambda$ and $\Lambda\bar{\Lambda}$ scattering lengths. $p\Lambda$, $p\bar{\Lambda}$, np , and nn are included for comparison. The np spin-triplet state — the deuteron — is a weakly bound state. The $\Lambda\Lambda$ scattering length is an order of magnitude smaller than the np scattering length, which is evidence that there is no $\Lambda\Lambda$ dibaryon bound state.

attractive interaction rather than repulsive.

3.7.2 $\Im f_0$

The few measured baryon-antibaryon scattering lengths have imaginary components on the order of 1 fm. The imaginary part of the $p\bar{p}$ scattering length has been measured to be $\Im f_{0,p\bar{p}} = 0.95 \pm 0.12$ fm, and the value for $p\bar{n}$ is $\Im f_{0,p\bar{n}} = 0.83 \pm 0.07$ fm [69]. Analysis of the STAR $p\bar{\Lambda} \oplus \bar{p}\Lambda$ correlation functions in 200 GeV Au–Au [45] collisions with (without) accounting for residual correlations yields $\Im f_{0,p\bar{\Lambda}} = 0.82 \pm 0.28$ fm ($\Im f_{0,p\bar{\Lambda}} = 1.00 \pm 0.21$ fm) [46]. It has been hypothesized that baryon-antibaryon pairs in general may have similar annihilation cross sections at comparable relative momentum [46]. Indeed, as many $B\bar{B}$ interactions are unknown, the UrQMD model treats all annihilation cross-sections as being the same as the $p\bar{p}$ annihilation cross-section at the same $\sqrt{s}$ [36]. In contrast with these various measurements and models, our measured scattering length is

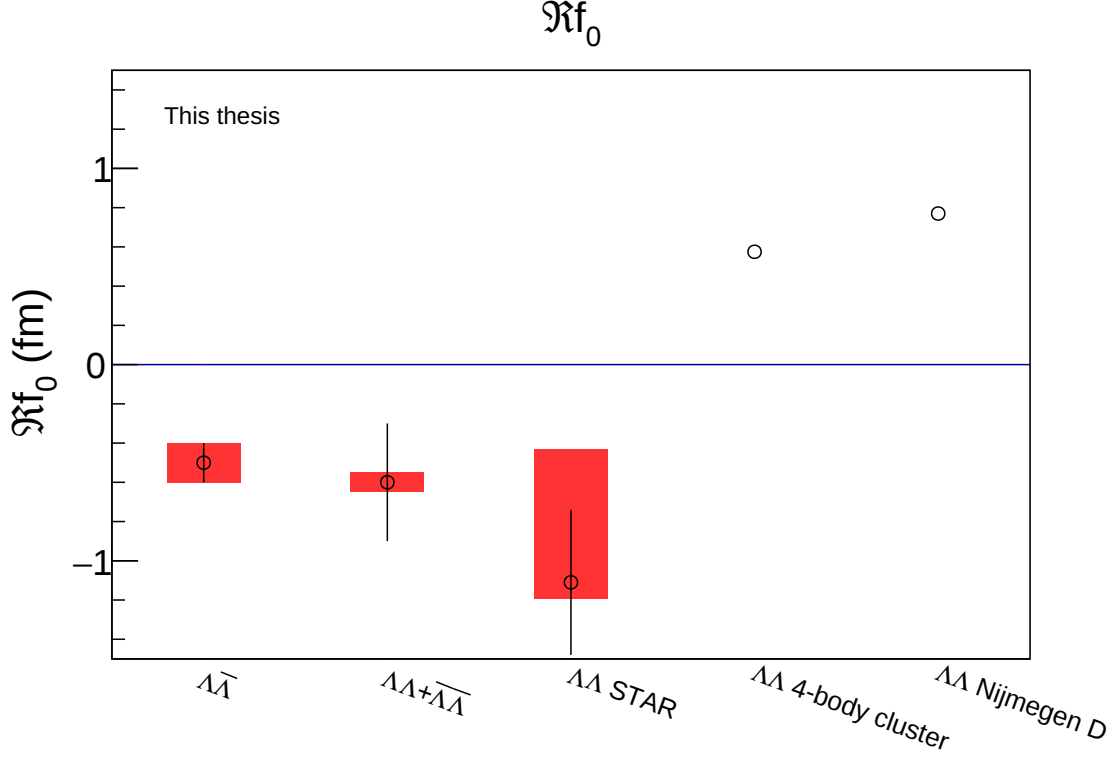


Figure 3.37: Several measurements of the real part of the $\Lambda\Lambda$ and $\Lambda\bar{\Lambda}$ scattering lengths. The STAR measurement [49] of the $\Lambda\Lambda$ scattering length agrees with our measurement within their systematic uncertainty. Our measurement is similar in magnitude but opposite the sign of $\Lambda\Lambda$ scattering lengths calculated from hypernuclear structure models [37, 38] applied to a measured ${}_{\Lambda\Lambda}{}^6\text{He}$ nucleus [48].

$\Im f_{0,\Lambda\bar{\Lambda}} = 0.14 \pm 0.09$ (stat) ± 0.02 (sys) fm, about a factor of 6 smaller than the other values.

3.7.3 d_0

d_0 enters into the wave function as a correction to the scattering amplitude, and in general it has a less profound effect on the correlation function than the scattering length. There are a few different treatments of d_0 across the field of research. When the value is known, such as for pp femtoscopy [29], it is implemented as a fixed value in the fits. In other cases, it has been set to zero to reduce the number of free fit parameters and avoid over-fitting [46, 64, 45]. In this analysis as well as the STAR analyses of $\bar{p}p$ and $\Lambda\Lambda$ correlations [44, 49], d_0 is left as a free parameter with the intention of letting nature tell whatever story it wants to tell.

The moral of that story seems to be that d_0 is not very well constrained. We obtained

$d_{0,\Lambda\Lambda} = 5.3 \text{ fm} \pm 2.7 \text{ (stat)} \pm 0.3 \text{ (sys)}$ — about 50% uncertainty. For $\Lambda\bar{\Lambda}$, we fared better with 17% uncertainty: $d_{0,\Lambda\bar{\Lambda}} = 1.7 \text{ fm} \pm 0.2 \text{ (stat)} \pm 0.2 \text{ (sys)}$. STAR measured a larger effective range $d_{0,\Lambda\Lambda} = 8.52 \text{ fm} \pm 2.56 \text{ (stat)}^{+2.09}_{-0.74} \text{ (sys)}$ — about 30–40% uncertainty. And for their recent measurement of the $\bar{p}p$ interaction, they found $d_{0,\bar{p}p} = 2.14 \text{ fm} \pm 0.27 \text{ (stat)} \pm 1.34 \text{ (sys)}$ — about 60% uncertainty.

While the uncertainties are all fairly large, the measured values are at least in the same ballpark as the effective range values measured for low-energy nucleon-nucleon scattering, all of which lie in the 2.5 to 3 fm range [70]. This suggests that it isn't hopeless to try to extract d_0 from femtoscopic analyses, but those analyses will need more statistics than are available for $\Lambda\Lambda$ at this time.

3.7.4 Radii

Figure 3.38 shows the extracted 1D radii R_{inv} for $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$ as a function of centrality and the average transverse mass m_T of the measured pairs. Also included are the radii for $\pi^\pm\pi^\pm$, $K^\pm K^\pm$, $K_S^0 K_S^0$, pp , $\bar{p}\bar{p}$ [29], and $p\Lambda \oplus \bar{p}\bar{\Lambda}$ (unpublished, [71]). With $R_{\Lambda\bar{\Lambda},0-10\%} = 2.9 \pm 0.3 \pm 0.3 \text{ fm}$, $R_{\Lambda\bar{\Lambda},10-30\%} = 2.3 \pm 0.2 \pm 0.2 \text{ fm}$, and $R_{\Lambda\bar{\Lambda},30-50\%} = 1.8 \pm 0.2 \pm 0.2 \text{ fm}$, the $\Lambda\bar{\Lambda}$ points seem to follow the m_T scaling of the pions, kaons, and protons. The $\Lambda\Lambda$ and $p\Lambda$ radii are quite large however, and while they are consistent with each other, they do not follow the m_T scaling trend. See Figure 3.39 for a qualitative illustration and interpretation of these results.

There are two peculiar details in these results that need explanation. What kind of effect would introduce an asymmetry between the particle-particle radius and the particle-antiparticle radius? Furthermore, why do the particle-particle radii deviate so much from the loose m_T trend observed for the pions, kaons, and protons? An ideal explanation would answer both of these questions with a single effect (e.g. perhaps there is some unaccounted for detector effect modifying the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ correlation function in such a way as to make the radius appear larger). If the measured results are correct, we would like to identify the underlying physics. If the measured results are incorrect, we want to determine where the problem lies in the experimental method or analysis process. In what follows, we will offer several hypotheses for why these results are seen, but first, let us discuss why the difference in radii is unusual.

If one considers the two-particle emission functions to be simple convolutions of the single-particle Λ and $\bar{\Lambda}$ emission functions, and one assumes that there is no asymmetry between those single-particle emission functions at the $\sqrt{s}$ and μ_b of the LHC, then in the absence of other confounding effects one would assume that $\Lambda\Lambda$, $\bar{\Lambda}\bar{\Lambda}$, and $\Lambda\bar{\Lambda}$ would all have the two-particle emission functions and same radii. While the $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$ correlation functions and radii are consistent with each other, we see that the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii are larger than the comparable centrality $\Lambda\bar{\Lambda}$ radii by almost a factor of two.

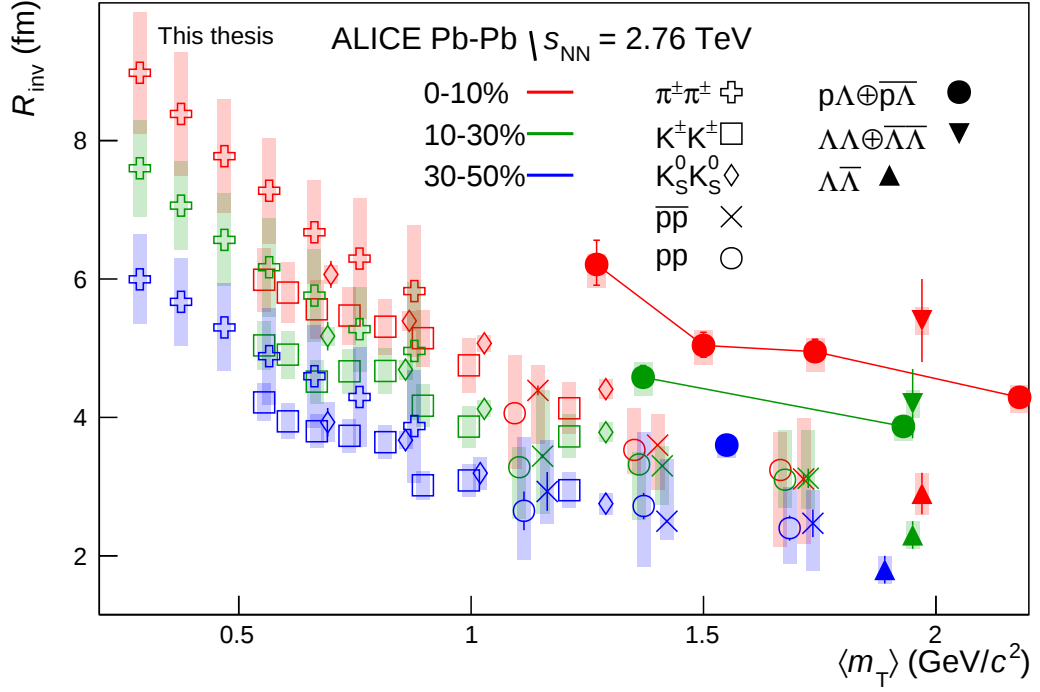


Figure 3.38: 1D radius R_{inv} vs transverse mass m_T for many different pairs measured by ALICE in $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collisions [29]. The $\Lambda\bar{\Lambda}$ radii appear to follow the trend of m_T -scaling [32, 13], but not $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ or $p\Lambda \oplus \bar{p}\bar{\Lambda}$ (unpublished, [71]).

It is possible that the anomalous differences between the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$ radii may be an error introduced by over- or under-corrected detector effects such as splitting and merging. If this is to blame, it would probably afflict the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ correlation functions. The $\Lambda\bar{\Lambda}$ correlation function is fairly robust against such effects. The correlation function is very wide, while splitting/merging affects the low- k^* range. In addition, for detector effects to occur with same-sign particles in $\Lambda\bar{\Lambda}$, a proton would have to be identified as a pion or vice versa.

So given the cuts implemented in Section 2.3.3, why might there still be an issue? In short, it is possible that real correlated physics effects might show up in our average separation correlation functions. In that case, when we set our pair cuts to remove merging, we may have accidentally cut out physics effects as well.

There is some correlation between the measured positions of charged tracks and the momenta of those tracks and their parents. Through that correlation, the detector effects of splitting and merging can show up as false signals in relative momentum correlation

Qualitative illustration of $\Lambda\Lambda + \bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$ radius results

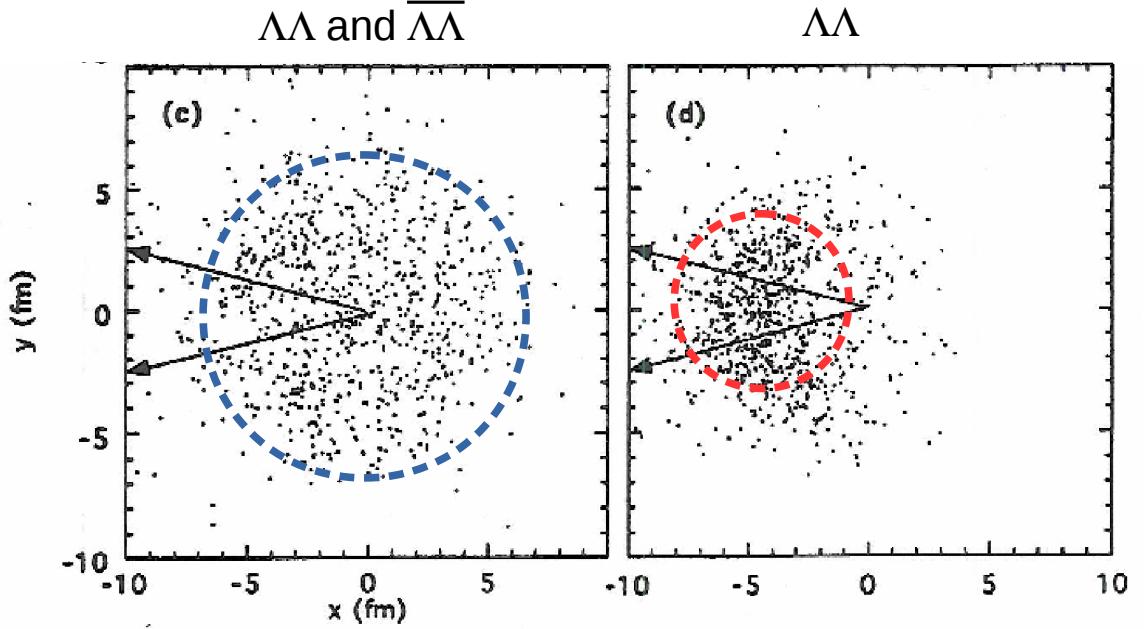


Figure 3.39: Qualitative representation of Λ femtoscopy radius results (dashed circles) inscribed on a plot of Monte Carlo-generated pion freezeout radii (original figure from [23]). $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$ have been measured to have a large homogeneity region of $R \sim 5.5 \pm 0.5$ fm for the most central collisions, while $\Lambda\bar{\Lambda}$ has a significantly smaller radius of $R \sim 3 \pm 0.3$ fm.

functions. It may be the case that true physics correlations can show up as false signals in average separation correlation functions, the purpose of which is to detect and remove splitting and merging. If true physics effects bleed into the average separation plots, it may lead us to impose strict cuts which unintentionally remove those physics effects from the femtoscopic correlation functions. Such a cut may lead to spurious fit results, such as the anomalously large $\Lambda\bar{\Lambda} \oplus \bar{\Lambda}\bar{\Lambda}$ radii. Future analyses may find merit in investigating alternative criteria to cut upon. For example, one might require the decay lengths of the V0s to differ by several centimeters. This would bias the daughter tracks to a larger separation, thus limiting splitting and merging effects, without introducing any relative momentum bias. Then it seems the risk of unintentionally cutting out true physics effects would be small, as it is unlikely that there is any correlation between the relative momentum of particles and their relative decay positions.

Of course, it is also possible that the radius results presented here are correct, that there truly is a physical difference between the emission functions of $\Lambda\Lambda/\bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$. One

way this might occur would be if the $\Lambda\bar{\Lambda}$ two-particle emission function contains a sizable number of pairs coming from pair creation.

For example, near-threshold pair creation of $\Lambda\bar{\Lambda}$ particles could cause low- k^* pairs to be emitted with only a small spatial separation. If there are a sizable number of these pairs, the two-particle emission function would have an additional component with a small radius, similar to the core-halo effect in pion femtoscopy (in addition to the usual Gaussian signal measured in pion correlation functions, there is also a very narrow enhancement in the lowest few k^* bins arising from a very large emission region, the pions that come from the decays of long-lived resonances). In the $\Lambda\bar{\Lambda}$ case, the two-particle emission function would have a component with a small radius from the pair-created particles, as well the expected “normal”-sized radius component arising from the convolution of the two single-particle emission functions.

The fit function we employ only has one radius parameter, not two, so the single radius radius may be something of an average of the size of these two regions. Because of that, the extracted $\Lambda\bar{\Lambda}$ radius would be smaller than the radius of either of the Λ or $\bar{\Lambda}$ single-particle emission functions, and thus smaller than the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radius. While this would account for a discrepancy between the particle-particle and particle-antiparticle radii, it seems like it would result in the $\Lambda\bar{\Lambda}$ radii looking too small, rather than the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii looking too big. When STAR was trying to account for the difference between the $p\Lambda$ and $p\bar{\Lambda}$ radii, they conjectured that pair creation might have something to do it [45]. But in their case, the particle-antiparticle radius looked too small, as would be expected from this effect.

These are a couple potential explanations for our curious results. It seems like the former issue, better treatment of detector effects, would be the best avenue to explore next, since it could account for the large $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii. The latter issue, pair creation giving rise to complicated emission functions, could possibly be explored by analyzing emission positions of particles in UrQMD.

Chapter 4

CONCLUSION

In this thesis, we presented centrality-dependent femtoscopy of Λ and $\bar{\Lambda}$ baryons measured in $\sqrt{s_{\text{NN}}} = 2.76$ TeV Pb-Pb collisions by the ALICE collaboration at the LHC. This is the first measurement of $\Lambda\bar{\Lambda}$ femtoscopy, and the first measurement of $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ to account for residual correlations using the same framework that is employed to calculate the primary correlations.

Our correlation function model follows the analytic framework laid out by Lednicky and Lyuboshits. It utilizes a one-dimensional Gaussian parametrization for the two-particle emission function that describes the relative separation of the emission points of particles. The particle interactions are accounted for with an s-wave scattering parameterization of the two-particle wave function, including antisymmetrization for the $\Lambda\Lambda$ and $\bar{\Lambda}\bar{\Lambda}$ pairs. For s-wave scattering, we used the asymptotic form of an incoming plane wave and an outgoing spherical wave, where the scattering amplitude depends on the scattering length f_0 and the effective range of interaction d_0 . When fitting, we constrained the same scattering parameters across each centrality bin to improve the precision of the measurement. In general, systematic uncertainties were dominated by the treatment of the non-femtoscopic background at low- k^* .

Using Monte Carlo HIJING events with ALICE reconstruction effects to determined particle efficiency, and THERMINATOR2 for thermal estimates of particle yields, we estimated the number of pairs where a Λ came from a Σ , Ξ^- , Ξ^0 , Ω , or false reconstruction. In doing so, we found that only 25% of pairs have two primary Λ and/or $\bar{\Lambda}$. The rest have at least one secondary Λ , and the physics describing the contribution of those pairs to the correlation function depends on the interactions between the parent particles. We estimated their correlation function contributions using the same analytic framework used to calculate the primary $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ and $\Lambda\bar{\Lambda}$ correlation functions. In that calculation, we assumed that the parent pairs interacted with the same strength (f_0 and d_0) and femtoscopic radius (R) as primary pairs, then we smeared out the relative momentum dependence according to the parent-to- Λ decay kinematics. We used the combined femtoscopic contribution of each pair

type to fit the measured correlation functions. In this way, we were able to turn the heavier hyperons' signal contamination into additional statistical power.

In this analysis, we extracted femtoscopic radii for three centralities of $\Lambda\bar{\Lambda}$ and two centralities of $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$. Hydrodynamic calculations suggest that there should be a rough scaling of the radii with the inverse of the transverse mass of the particles. The radii we have measured serve to complement other femtoscopic analyses by providing measurements at large transverse mass m_T . For $\Lambda\bar{\Lambda}$, we measured $R_{\Lambda\bar{\Lambda},0-10\%} = 2.9 \pm 0.3 \pm 0.3$ fm, $R_{\Lambda\bar{\Lambda},10-30\%} = 2.3 \pm 0.2 \pm 0.2$ fm, and $R_{\Lambda\bar{\Lambda},30-50\%} = 1.8 \pm 0.2 \pm 0.2$ fm. While these points do follow the trend set by pions, kaons, and protons, the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ break this scaling dramatically. Here we measured $R_{\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda},0-10\%} = 5.4 \pm 0.6 \pm 0.2$ fm and $R_{\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda},10-30\%} = 4.2 \pm 0.5 \pm 0.2$ fm. We did not have sufficient statistics to fit the 30-50% centrality bin. These $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii are nearly double their $\Lambda\bar{\Lambda}$ counterparts. It is unclear at this time what physics effects might lead to this separation. Pair creation might lead to $\Lambda\bar{\Lambda}$ radii being smaller than $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii, but would not account for $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii deviating significantly from the m_T trends. Another possibility is that the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ correlation function, which only extends through a few k^* bins, might have under- or overcompensated detector effects that have not been included within the systematic errors.

We measured the scattering parameters of these systems. This information can help constrain hadron cascade models, hypernuclear structure calculations, and neutron star equations of state. We found that the real part of the scattering length $\Re f_{0,\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}} = -0.6 \pm 0.3 \pm 0.05$ fm, which is within the systematic uncertainty of STAR's measurement. This interaction is small, about an order of magnitude smaller than the deuteron bound state, which supports claims that $\Lambda\Lambda$ has no dibaryon bound state. The effect of the nuclear final state interactions on the $\Lambda\Lambda$ pair appears to be small compared to the contribution from Fermi-Dirac suppression. The scattering length measured by femtoscopy experiments is approximately the same magnitude as the estimate arising from hypernuclear structure experiments, but they are opposite in sign. It is unclear how to rectify these differences at this time. The real parts of the scattering lengths of $\Lambda\bar{\Lambda}$ and $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ are comparable in size ($\Re f_{0,\Lambda\bar{\Lambda}} = -0.5 \pm 0.1 \pm 0.1$ fm), suggesting that the elastic interactions of these hyperons behave the same.

The imaginary part of the scattering length $\Im f_0$ accounts for the inelastic process of annihilation. As a baryon-baryon pair, $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ has no imaginary component to its scattering length, but we have measured it for $\Lambda\bar{\Lambda}$. Here, $\Im f_{0,\Lambda\bar{\Lambda}} = 0.14 \pm 0.09 \pm 0.2$ fm. This is considerably smaller than the measured values for pp and p $\bar{\Lambda}$, both of which appear to fall around 0.8 fm.

The effective range of interaction d_0 is a correction to the scattering amplitude, and its effect on the correlation function is small compared to that of the scattering length. Given that and the small FSI seen in $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$, it is not surprising that d_0 is not particularly

well constrained for $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$, with $d_{0,\Lambda\Lambda\oplus\bar{\Lambda}\bar{\Lambda}} = 5.3 \pm 2.7 \pm 0.3$ fm. The $\Lambda\bar{\Lambda}$ FSI effects extend to a much larger k^* range, so we have better resolving power for the effective range $d_{0,\Lambda\bar{\Lambda}} = 1.7 \pm 0.2 \pm 0.2$ fm.

Efforts are currently underway within ALICE to perform joint fitting of $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ ($\Lambda\bar{\Lambda}$) with other baryon-baryon (baryon-antibaryon) systems. The femtoscopic analyses of pp and p Λ have residual correlation effects from $\Lambda\Lambda$, as well as heavier pairs such as $\Lambda\Sigma$ and $\Lambda\Xi$. The underlying correlations from these parent pairs should be the same, regardless of the type of the primary correlation that is being measured. Whether it feeds down into pp, p Λ , or $\Lambda\Lambda$, the $\Lambda\Sigma$ correlation has one scattering length and one femtoscopic radius at a given event centrality and P_T . Likewise, $\Lambda\Lambda$ correlation function is the same whether it is a primary correlation as seen in this analysis, or a residual correlation in a p Λ analysis. Joint fitting—simultaneous χ^2 fitting of pp, p Λ , and $\Lambda\Lambda$ using shared scattering lengths and radii among the primary and residual correlations—has the potential to provide significantly stronger constraints on the interactions and radii of $\Lambda\Lambda$ than a solo fit.

We have used femtoscopy with residual correlation contributions to measure the interactions of $\Lambda\Lambda$ and $\Lambda\bar{\Lambda}$. Residual correlation contributions appear to be much larger in baryon-antibaryon pairs. Despite this, the measured scattering parameters and radii for $\Lambda\bar{\Lambda}$ appear to be robust. Nonetheless, some mysteries remain regarding both the $\Lambda\Lambda \oplus \bar{\Lambda}\bar{\Lambda}$ radii and scattering length. These may be resolved with further investigation of methods to remove detector effects. If issues such as these can be resolved, this approach may have applicability for measuring the poorly-known interactions of other pairs, such as $\Lambda\Xi^-$, ΛK^- , and $\Xi^- K^-$. The Pb-Pb data taken during the LHC's second run period (2015-2018) will afford a chance to examine these phenomena again at a higher energy, $\sqrt{s_{\text{NN}}} = 5.02$ TeV.

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