

Investigation Of Chiral Symmetry Restoration Using $\Xi(1820)$
Reconstruction From p-p, p-Pb, And Pb-Pb Collisions At ALICE

by
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DEDICATION/EPIGRAPH

This dissertation is dedicated to all of the wonderful and supportive people who encouraged and helped me reach this point. I would have never reached this point if not for the support of so many people that in some small way led me to this moment.

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ABSTRACT

The Large Hadron Collider (LHC) at CERN, Geneva, Switzerland, stands at the forefront of our current understanding of high energy physics. Several important measurements, including assisting in the discovery of the Quark Gluon Plasma (QGP), have been made in recent years thanks to the efforts of the LHC. Further study of the properties of matter in these systems proves to be instrumental in our understanding of the universe and the physics used to describe it.

Recent lattice quantum chromo dynamic (QCD) calculations seem to indicate the phenomenon of chiral symmetry restoration, which would result in changes in the masses of some hadron species under extreme conditions of pressure and temperature. These lattice QCD calculations highlight the idea of parity doubling, where the masses of negative parity particles decrease when the system approaches a pseudo-critical temperature and phase transition. Investigation into the modification of the affected particle's mass, width, and yield due to partial chiral symmetry restoration would prove instrumental to confirming some of the last remaining predictions of QCD.

Measurements of the $\Xi(1820)^-$ and its antiparticle were performed with the ALICE detector in p-p, p-Pb, and Pb-Pb collisions at LHC energies $\sqrt{s} = 13$ TeV for p-p and $\sqrt{s_{NN}} = 5.02$ TeV for p-Pb and Pb-Pb. The mass, width, and yield of $\Xi(1820)$ are obtained and compared in various collision systems. The yield ratios of $\Xi(1820)$ to $\Xi(1530)$ and Ξ are shown and discussed.

Analysis into the $\Xi(1820)^\mp$ in p-p, p-Pb, and Pb-Pb collisions at LHC shows several signatures of chiral symmetry restoration. Most notably, a 2.38σ difference between the width of $\Xi(1820)^\mp$ signals from minimum bias p-p $\sqrt{s} = 13$ TeV and central Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data is observed. While very impressive, more statistics are required to reach a conclusive statement.

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1 Introduction

1.1 Extreme conditions for matter

There is considerable evidence that shortly after the *Big Bang* the universe was vastly different compared to the universe as it is today. With such high temperatures and energies, the four fundamental forces of the universe could not be distinguished from each other until the universe cooled and the forces separated. Despite that, understanding the universe and how it has evolved is one of the main research topics across all of science.

In the first few microseconds after the Big Bang, the universe was in a state called the Quark Gluon Plasma (QGP). The QGP is a state of matter where the energy and/or density of the matter is so high that the individual quarks that make up all of the confined hadrons in the universe were free from each other. Due to the temperature of the universe being over 150 MeV, all of the quarks produced were extremely relativistic [1]. Furthermore, due to asymptotic freedom, the quarks would have barely interacted at all in this strange, hot, colorful system [2].

Despite the difficulties, modern science has been able to recreate the QGP using relativistic heavy ion collision accelerators such as those at the Large Hadron Collider (LHC) [3] and the Relativistic Heavy Ion Collider (RHIC) [4]. Though small, these QGP systems are instrumental in our understanding of the evolution of the universe and the strong nuclear force. Thus, experiments performed at these heavy ion collision accelerators focus a great deal of time and resources into discovering the properties of the QGP, and how these properties change as the system cools.

As seen in the QCD phase diagram, Figure 1.1, strongly interacting matter can enter a new state of matter after experiencing a phase transition. It should be noted that, as seen in Figure 1.1, LHC experiments are high energy collision experiments, so they can reach higher temperatures with very low chemical potential. RHIC experiments are also collision experiments, but have lower collision energy that leads to partial stopping of the heavy ions, which increases the baryon chemical potential. They therefore reach slightly lower temperatures with higher chemical potential in comparison to the LHC. A *Quark Gluon Plasma* occurs when the temperature of the system

rises to a point that a phase transition, called *chiral symmetry restoration*, can occur. When matter goes through this phase transition, the chiral symmetry of the system is restored. Chiral symmetry restoration is accompanied by de-confinement of the color neutral state of hadrons into its colored quark constituents. This phase transition of matter going from normal matter to the QGP matter due to chiral symmetry restoration is the main focus of this dissertation. Similarly, the phase transition of the QGP to normal matter is signaled by chiral symmetry breaking.

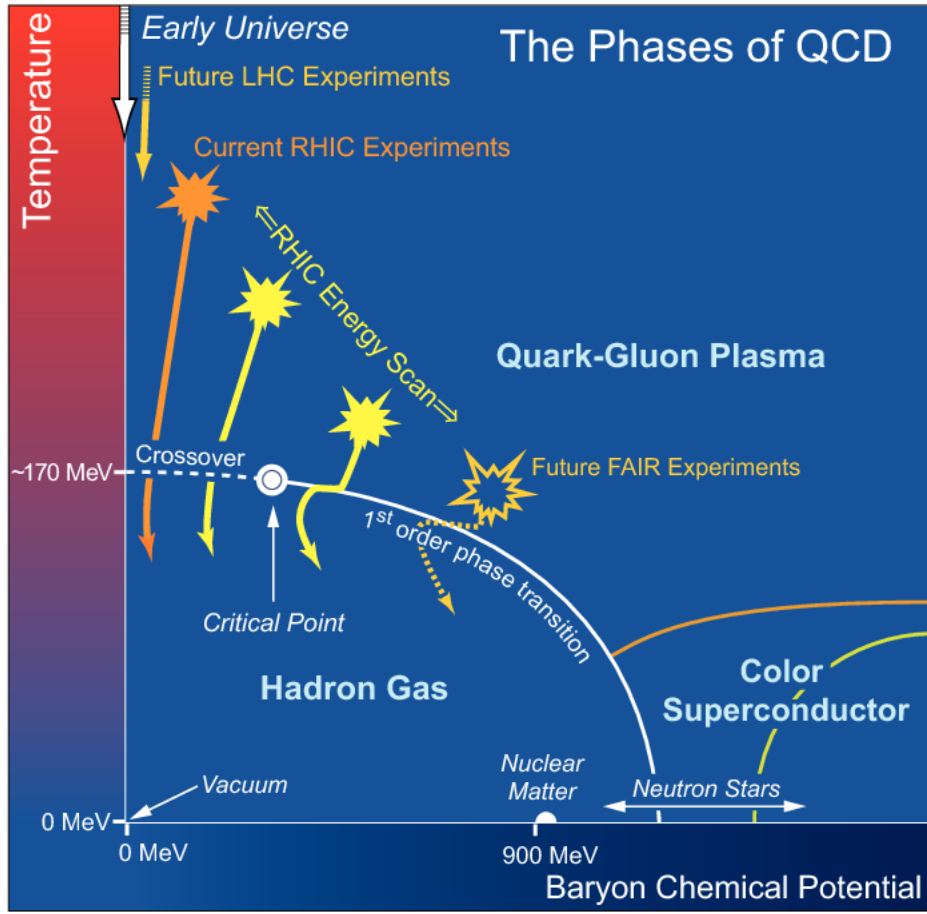


Figure 1.1: Illustration of the phase diagram of QCD matter [5].

1.2 Quark Gluon Plasma

After heavy ions that have been accelerated by the LHC to speeds close to the speed of light collide, they explode in a “fireball” that results in even the quarks and gluons being temporarily freed from the hadrons in which they are normally bound [6]. The speeds necessary to cause such a reaction are so great, that the Lorentz contraction can take the normally spherical shapes of the nuclei and reduce them to look more like thin disks just before the collision as seen in Figure 1.2 [7].

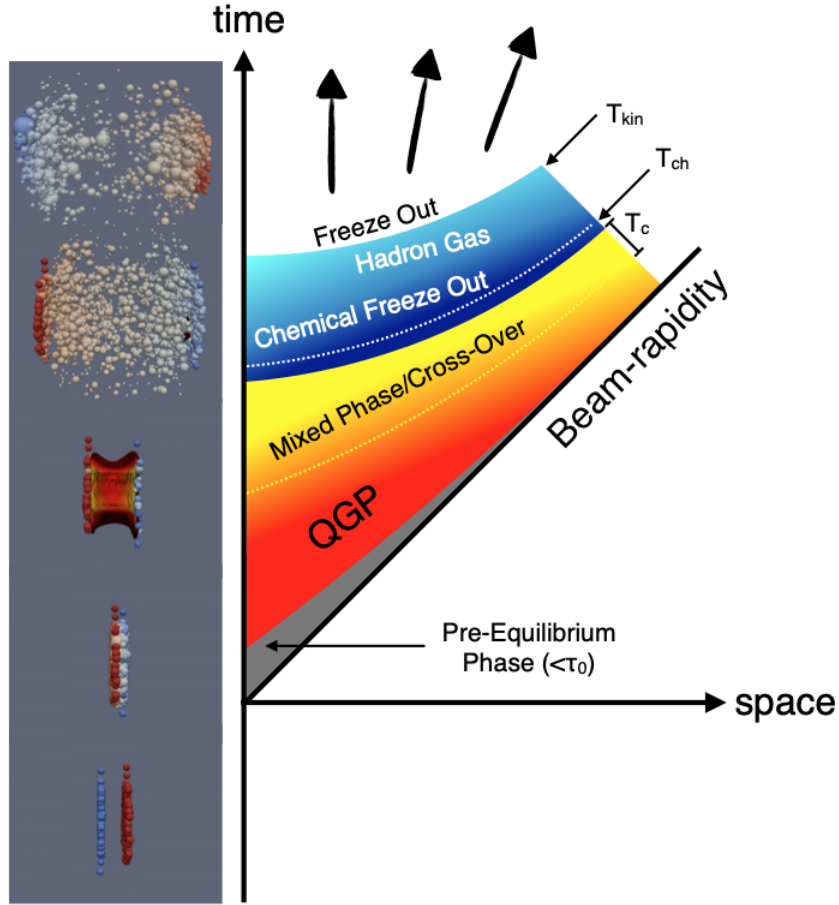


Figure 1.2: A schematic view of the evolution of an ultra-relativistic heavy-ion collision [8].

During the first fraction of 1 fm/c (about $3.4 \cdot 10^{-24}$ sec) after a heavy ion collision, quarks and gluons are produced in a quasi-free de-confined system referred to as the pre-equilibrium stage.

This de-confined system occurs when the quarks that are normally “confined” into hadrons of 2 or 3 quarks to remain color neutral, are no longer bound to this system and are free to interact with other quarks and gluons in the system. Though the “total color” of the QGP is still color neutral, locally this is not the case. The system reaches a state of equilibrium between the quarks and gluons at around $1 \text{ fm}/c$.

Between $1 \lesssim \tau \lesssim 10 \text{ fm}/c$, the system expands and reduces its temperature with elastic and inelastic collisions in what is referred to as the expansion stage. Due to the internal pressure of the QGP built up from the shockwave caused by the collision of two heavy ions, as it expands it also cools until it reaches the critical temperature T_c . At this point, the QGP will begin hadronization, the point at which quarks and gluons must form hadrons so that no “free” quark can remain. During this mixed-phase/cross-over phase transition, the system will approach the chemical freeze-out temperature, $T_{ch} \approx 155 \text{ MeV}$, when the number of hadrons is fixed and inelastic collisions stop [9]. While the number of long-lived hadrons produced is fixed at this point, the distances between the hadrons are still small enough for pseudo-elastic collisions to occur. These pseudo-elastic collisions arise when two longer lived hadrons scatter through a short lived resonance state, which then reverts to the original hadrons. As seen in Figure 1.3, there are two ways for the yields of particles with very short lifetimes, such as resonances, to be affected. Regeneration takes place when two stable hadrons collide and form a resonance, that then decays shortly afterwards, and this increases the yield of that resonance. Re-scattering transpires when the product of a resonance decay collides with other particles, thus making the reconstruction of the initial resonance difficult and decreasing the yield of that resonance. Purely elastic collisions will also happen that will change the particle’s momentum, but do not generate a resonance state.

Only when the system has reached the point that the kinematic distribution of the hadrons has been fixed and there are no more elastic collisions, known as the kinetic freeze-out temperature, $T_{kin} \approx 100 \text{ MeV}$, which occurs at around $15 \text{ fm}/c$, will the system no longer be considered under extreme conditions [9] [10]. These free particles are then measured by ALICE or other detectors.

With the discovery of the QGP, there have been several interesting properties observed that

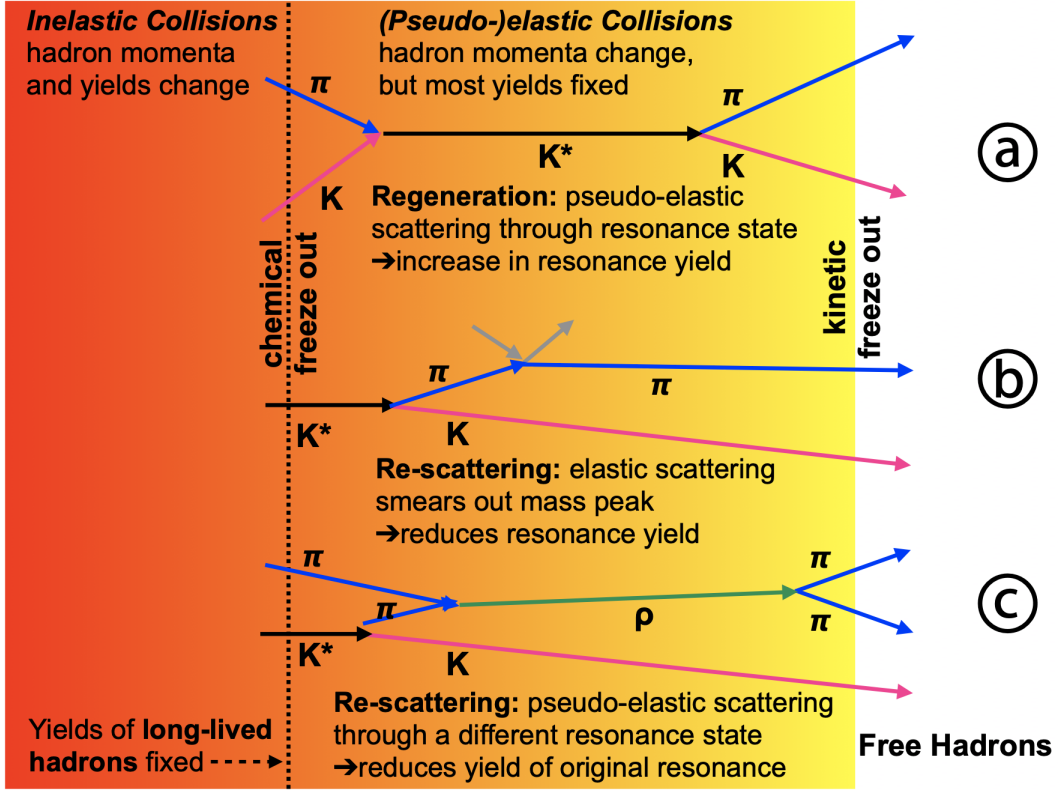


Figure 1.3: Possible interaction of resonances between chemical and kinetic freeze-out [8]. Going from left to right, time increases and temperature decreases.

are not seen in normal matter. The most notable property of the QGP is that it has a “perfect” liquid nature, meaning that it has almost no frictional resistance or viscosity [3]. A “perfect” liquid like the QGP could flow against itself without any frictional forces. In fact, because of the liquid nature of the QGP, hydrodynamics is used heavily to determine the properties of the “flow”, such as the second moment of the Fourier harmonics in the azimuthal distribution of particles with very strong interactions v_2 , and how it relates to the dynamic expansion of the QGP [3]. In particular, the even moments, such as v_2 , have been observed by RHIC experiments to indicate a strongly flowing fluid medium when compared to viscous hydrodynamic models [3]. However some of the most notable questions about the QGP are about the phase transition that occurs as the system cools and begins hadronization. Notably, the chiral symmetry breaking effects that are the focus

for this discussion prove very difficult to investigate.

1.3 Resonant States

To investigate the nature of chiral symmetry breaking, an observable is required that is not only measurable in the creation of the QGP and chiral symmetry breaking, but is also sensitive to the very small timescales of the QGP system. Thankfully, resonances satisfy these conditions quite well for our purpose.

Resonances are short-lived excited particles that have the same quark content as stable ground state particles, but different parity and/or angular momentum. Due to the fact that resonances decay strongly, they conserve quantum numbers across their decay channels. They normally have lifetimes of a few fm/c, comparable to the lifetime of the fireball created in heavy-ion collisions [11]. While this makes them ideal for studying the QGP as it undergoes hadronization and chemical and kinetic freeze-out, they can not be observed directly and must be observed via reconstruction of their decay products. Any information of the resonance decaying before chirality is completely broken, such as mass, width, and yield, is observed in the kinematic information of the daughter particles.

Resonances can be characterized by the Breit-Wigner formula:

$$f(E) \propto \frac{1}{(E - M_0)^2 + \Gamma^2/4} \quad (1)$$

$f(E)$: is the probability density formula

M_0 : is the resonance mass

Γ : is the full-width half-max

The Breit-Wigner formula shares some similarities to other distributions such as the Gaussian. Most notably, the resonance mass and Γ of the Breit-Wigner is very similar to the mean and σ of the Gaussian in their use and implementation. However, there are some differences that make the Breit-Wigner a bit unique. The most important feature of the Breit-Wigner for this discussion is

the Γ . The Γ is not only important for determining the distribution of the function around the peak resonance mass, but is also related to the properties of the resonance investigated, such as the lifetime as seen in Eq. 2:

$$\Gamma = \frac{\hbar}{\tau} \quad (2)$$

where Γ is the full-width half-max (FWHM) of the signal, $\hbar$ is Planck's reduced constant, and τ is the lifetime of the resonance.

The Γ 's representation of the full-width half-max is slightly different than the representation of the σ . Γ represents the range of the distribution (full-width) that the Breit-Wigner has when the distribution has decrease to half of the peak value (1 to 0.5, half-max). σ for the Gaussian is determined as the value needed for the distribution to decrease to 0.60653 ($e^{-1/2}$) of the initial distribution. Despite this, Γ can be related to σ using this simple numerical relation:

$$\sigma = \Gamma/2.35. \quad (3)$$

Eq. 2 is closely related to the Heisenberg uncertainty relation for energy and time:

$$\Delta E \Delta t \geq \hbar/2 \quad (4)$$

where ΔE is the uncertainty of energy and Δt is the uncertainty of time. Similarly, the uncertainty of energy can be related to the FWHM ($\Delta E \approx \Gamma$) just as the uncertainty of time can be related to the lifetime of the resonance ($\Delta t \approx \tau$).

Particle resonances have lifetimes of the order of about 10^{-23} s, thus making them perfect for investigating the QGP. However, not all resonances have the same lifetime. Some resonances, such as the $\phi(1020)$ resonance, have long lifetimes ($c\tau = 46$ fm/c) in comparison to the lifetime of the “fireball” [12]. Other resonances, such as the $\Lambda(1405)$ resonance, have such short lifetimes ($c\tau = 3.9$ fm/c) that they should completely decay inside the QGP [12]. Because different resonances decay with different lifetimes, it is possible to investigate the different temperatures of the QGP as

it cools by investigating the resonances that would decay in the same time frame that the system would be in for a corresponding temperature. One resonance may be useful for investigating the chemical freeze-out temperature, while another may be useful for investigating the kinetic freeze-out temperature.

As will be explained further in Section 3.1, a resonance particle can also have a “parity partner,” in that two resonances are similar to each other for all quantum numbers except for parity. Normally, the negative parity resonance particle will have a higher mass than the positive parity partner, thus leading to the possibility that the two resonances may decay via different decay channels [12].

1.4 Chirality

Chirality in the simplest definition can be defined as the “handedness” of a particle. More accurately, it is defined as the spin of a particle and direction the spin vector would point to using a right-handed or left-handed system. The most common use of chirality is in the definition of helicity, the projection of the spin vector onto the momentum vector. If a particle with right-handed chirality spins such that its spin vector is in the same direction of the momentum vector, the particle is said to have “positive” 1 helicity. If the same particle with the same direction of spin and momentum vector, but is in a left-handed system, then the spin vector would point in the opposite direction of the momentum vector, leading to “negative” 1 helicity.

However, helicity is well-defined only for massless particles traveling at the speed of light. For a massive particle, it is possible to seemingly reverse the particle’s helicity if an observer was to use a reference frame that was moving faster than the spinning particle in question. This change in reference frame would make the particle able to be passed up and appear to move backwards, thus causing the “apparent” momentum to reverse and as such the helicity.

It must be mentioned that chirality is not the same as helicity. Helicity is dependent on the angle between the spin vector and momentum vector, the spin of the particle, the handedness of the system, and the reference frame used. Chirality is the handedness of the particle and as such is not frame dependent.

Using QCD to further probe the complexities of chirality becomes instrumental at this point. Handedness of quark fields, as it is used in QCD, is defined as

$$q_L \equiv \frac{1 - \gamma^5}{2} q, q_R \equiv \frac{1 + \gamma^5}{2} q, \quad (5)$$

where q denotes the quark fields and γ^5 is the chirality operator. Normally, the four-vector γ^μ can be defined as any set of matrices where the components have properties of

$$(\gamma^0)^2 = 1, (\gamma^1)^2 = -1, (\gamma^2)^2 = -1, (\gamma^3)^2 = -1, \quad (6)$$

and the components have the anticommute property of

$$\gamma^\mu \gamma^\nu = -\gamma^\nu \gamma^\mu \quad (7)$$

if $\mu \neq \nu$ [13]. The chirality operator can be defined as

$$\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3, \quad (8)$$

which leads to the anticommute property ($\gamma^\mu \gamma^5 = -\gamma^5 \gamma^\mu$) and the Hermitian property ($\gamma^{5\dagger} = \gamma^5$ or $(\gamma^5)^2 = (-1)(1)(-1)(-1)(-1) = 1$). The notation of Eq. 5 is useful in the case of masses equal to zero since L and R would correspond to helicity -1 and +1 respectively. Similarly,

$$\bar{q}_L \equiv q_L^\dagger \gamma^0 = \bar{q} \left(\frac{1 + \gamma^5}{2} \right), \bar{q}_R \equiv \bar{q} \left(\frac{1 - \gamma^5}{2} \right). \quad (9)$$

Using the notation of the left- and right-handed quark fields, we can rewrite the normal Lagrangian density equation for QCD into one that shows the effects that chirality can have. The Lagrangian density for QCD is given by

$$L_{QCD} = \bar{q}(i\gamma^\mu D_\mu - M_q)q - \frac{1}{4}G_{\mu\nu}^a G_a^{\mu\nu}, D_\mu = \partial_\mu + ig_s \frac{\lambda_a}{2} A_\mu^a, \quad (10)$$

where A_μ^a denotes the gluon field, g_s is the strong coupling constant, and $M_q = \text{diag}(m_u, m_d, \dots)$ is the current quark mass matrix [14]. γ^μ and λ_a correspond to the Dirac and Gell-Mann matrices respectively [14]. While L_{QCD} has several global symmetries including the local $SU(3)$ color gauge symmetry, for our purposes, we will focus on the symmetry of chirality [14]. By replacing the normal quark fields q with the left- and right-handed quark fields from Eq. 5 and Eq. 9, we can rewrite L_{QCD} as

$$L_{QCD} = \bar{q}_L i \gamma^\mu D_\mu q_L + \bar{q}_R i \gamma^\mu D_\mu q_R - (\bar{q}_L M_q q_R + \bar{q}_R M_q q_L) - \frac{1}{4} G_{\mu\nu}^a G_a^{\mu\nu}. \quad (11)$$

In the case where the energies for QCD are larger than the light quark masses ($\Lambda_{QCD} \approx 200$ MeV $\gg m_{u,d} \approx 10$ MeV), we can ignore the mass terms and further simplify the previous equation. Finally, L_{QCD} is invariant under the separate transformations

$$q_L \rightarrow e^{-i\vec{\alpha}_L \cdot \vec{\tau}} q_L, q_R \rightarrow e^{-i\vec{\alpha}_R \cdot \vec{\tau}} q_R \quad (12)$$

where $\tau^a = \sigma^a/2$ are operators in (u-d) isospin space and $\alpha_{L,R}$ are three real angles [14]. The currents associated with the symmetries seen in Eq. 12 are respectively [14]

$$\vec{j}_L^\mu = \bar{q}_L \gamma^\mu \vec{\tau} q_L, \vec{j}_R^\mu = \bar{q}_R \gamma^\mu \vec{\tau} q_R. \quad (13)$$

According to Noether's theorem, a conserved quantity can be related to a continuous symmetry [15]. That means that if the chirality currents seen in Eq. 13 can satisfy the continuity equation

$$\partial_\mu \vec{j}^\mu = 0 \quad (14)$$

then chiral symmetry is satisfied. In the massless case, this symmetry means that left-handed and right-handed fermions are two different particles, also known as Weyl fermions [14]. In normal nuclear matter, they are the same particle. In fact, a Dirac fermion is composed of the left and

right handed components of the Weyl fermions [14].

In essence, this means that QCD is symmetric under chirality if the system can ignore the mass terms in the Lagrangian. Thus there are two ways for chirality to be satisfied. First, if the particles for the Lagrangian are massless, then the mass term is equal to zero. Second, if the energies of the system are much higher than the mass of the particles, the mass term will have such a low contribution that it can be ignored.

Chirality, much like other symmetries, is normally considered to be conserved. This means that a right-handed particle and left-handed particle are indistinguishable from each other. However, the complexities of chirality come from the fact that mass introduces a way to break this symmetry. Using the previous example, there is an important fact that the helicity of a particle changes if a reference frame is used that speeds past the particle in question. In the case of a massless particle, such as a photon, the particle is already moving at the speed of light, so no reference frame could be used to change the helicity of the particle. For these massless particles, regardless of the point of view of the observer or the reference frame used, the spin appears in the same direction along its momentum vector, thus helicity is the same as chirality in this case. As no massive particle can reach the speed of light, there will always be a reference frame that could change the helicity, thus preventing the relation between chirality and helicity to be as clear-cut as the one for massless particles.

The confusion of chirality can be summarized in this simple statement. Massless particles have chirality equal to helicity and are the same in all systems. Massive particles have chirality that “may not” be equal to helicity, thus meaning that right-handed and left-handed particle are different and can be distinguished from each other. If a massive particle could become indistinguishable from a similar particle of opposite chirality, either from reaching light speed or completing a phase transition, chiral symmetry would be “restored.”

According to thermal quantum field theory, at high energy or temperature the spontaneously broken chiral symmetry can be restored [16]. At high temperatures, the kinetic energy required for pair production increases to a point that the pairs evaporate and chiral symmetry is restored.

This phase transition causes massive, chiral symmetry broken particles to transition to massless, chiral symmetry restored particles. Interestingly, the temperature for this transition is at the same temperature needed for de-confinement. So by the time that temperatures are high enough for chiral symmetry to be restored, the particles have entered into a de-confined state, such that the “particle” would evaporate into separate quarks and gluons. While the exact temperature needed for de-confinement is still being investigated, most experiments and theories set the critical temperature at around 150 MeV [9].

It is also possible for chiral symmetry to be restored using high density instead of high temperature to create a “color superconductor” [14]. However, instead of the high kinetic energy from high temperatures, the high density causes the chemical potential of the pair production to increase. With this higher chemical potential, the energy required to induce pair production increases to the point that particles simply can not be produced.

The main question becomes, “What is a signature of chiral symmetry restoration?” As was seen in the previous equations, chiral symmetry is restored when particles are massless. Chiral symmetry is “broken” as soon as mass is introduced that can break the symmetry of left-handed and right-handed particles. For chiral symmetry to be “restored,” the mass of a particle should vanish. However, that corresponds to the case in which chiral symmetry is “fully” restored. In the case where chiral symmetry is “partially” restored, the mass of some particles would decrease and approach the mass of other particles whose mass did not change. This is explained in Section 1.5, but when chiral symmetry is partially restored, the mass difference between two parity partners disappears. While this may seem simple, actual observation of chiral symmetry restoration is very difficult.

Chiral symmetry breaking is most obvious in the mass of the proton and neutron. Both nucleons are formed of three elementary light quarks. The proton(neutron)’s mass of about 938(940) MeV/ c^2 is composed of two(one) up quarks of mass 2.3 MeV/ c^2 and one(two) down quark with mass 4.8 MeV/ c^2 , contributing only 9.4(11.9) MeV/ c^2 to the nucleon’s respective mass. The remaining 99% of mass is contributed by the binding energy between the quarks, which occurs because of chiral

symmetry breaking.

Few experiments have ever shown a measurable effect correlated to chiral symmetry restoration. In fact, the most well known paper on the potential effects of chiral symmetry restoration, the paper by the NA60 collaboration [17], refers to the $\rho \rightarrow \mu^- \mu^+$ invariant mass distribution as having a “broadened” width in In-In collisions as seen in Figure 1.4 [17]. This however was done in a system that is theorized to have no QGP forming. Because approximately restored chiral symmetry does not require a phase transition, as opposed to full chiral symmetry restoration that does require a phase transition, it is possible to start in a dense hadronic system and then be “completed” in a de-confined system. Analysis from Rapp has shown that hadronic many-body approaches, such as the meson-gas contributions to the self-energy that has the feature of strong broadening of the ρ spectral function, can describe the spectra in question [18]. This is mainly because the imaginary part of the retarded self-energy contains contributions that have the same sign due to the retardation condition, while the real parts contain both attractive (negative) and repulsive (positive) interactions, which tend to compensate each other given the large set of excitations [18].

This explanation from Rapp, while informative, does suffer from a few problems as it relates to our investigation into the potential effects of chiral symmetry restoration. First, while the Quark Gluon Plasma is filled with quarks and gluons, Rapp’s analysis focuses on hadronic many-body interactions [18]. Thus, Rapp’s theory would not be able to describe phase transitions, let alone the potential effect on quarks and gluons. Second, while the QGP as it is created evolves in a system from relatively low temperature (speed up of particles), to high temperature (collision), to finally low temperature (thermal expansion), Rapp’s analysis was focused on the density of the system [18]. This means that the evolution of the systems would not follow the same trajectories as seen on the QCD phase diagram from Figure 1.5. While Rapp’s analysis into the spectra of the NA60 paper can be considered a possible explanation of the signal, it may not be the actual effect. One must be careful with the approach used in determining the possible explanations of the effect in question, otherwise possible alternatives may have no bearing for our analysis. There will be hadronic effects due to increased hadronic density near the critical temperature T_c , according

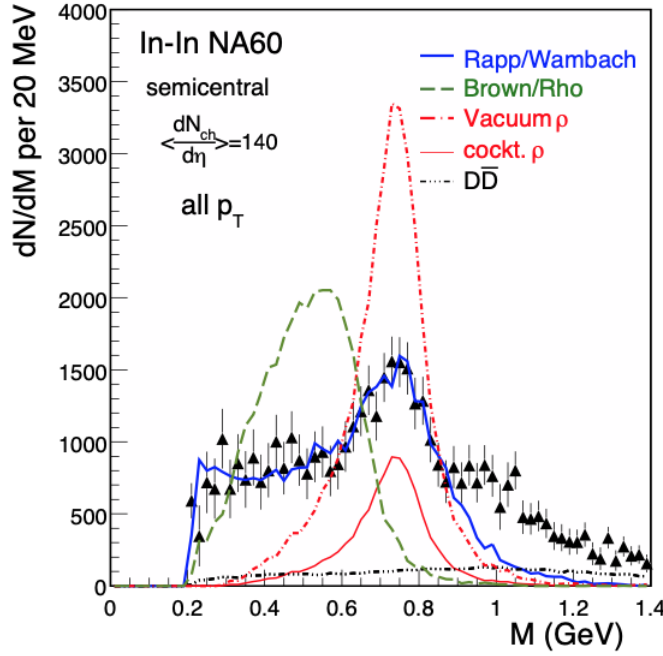


Figure 1.4: Excess di-muons above the known decay sources compared to theoretical predictions, renormalized to the data in the mass interval $M < 0.9$ GeV. No acceptance correction applied [19].

to Rapp, and there will be crossover transition effects such as parity doubling, according to lattice QCD in Section 1.5. These two effects may not be disentangled depending on the system and its evolution.

Nevertheless, the search for an experimentally observable chiral symmetry restoration signal continues. This is mostly driven by the fact that evidence of chiral symmetry restoration is the final piece of evidence needed to confirm QCD's ability to predict nuclear systems at this energy range.

1.5 Lattice QCD

Understanding the nature of how quantum objects interact with one another, let alone any collective features they may have as they form hadrons or baryons, is a very complicated subject. There are multiple ways to investigate the equations used to describe quantum phenomena, such as Dirac's

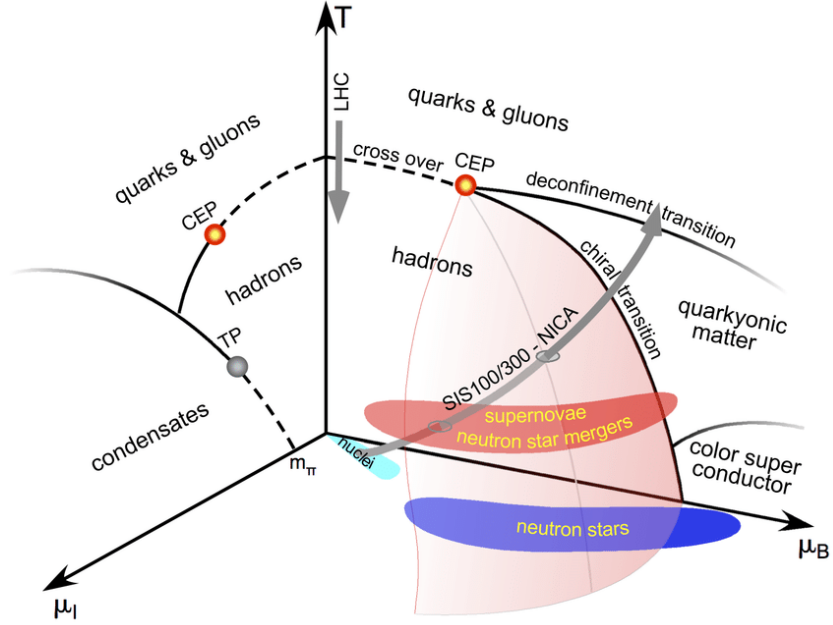


Figure 1.5: 3D representation of QCD phase diagram. μ_B corresponds to net-baryonic density. μ_I corresponds to isospin density, which can be related to the mass of π , but no deviation is used for this analysis [20].

wave equations or Heisenberg's matrices. However, one of the newest and potentially most powerful method to investigate these quantum features is the use of lattices, since the theory of strong interactions can only be solved numerically in the regime of interest for the experiments.

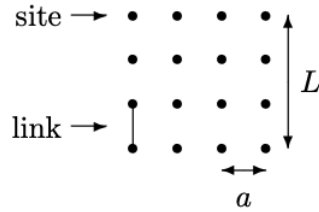


Figure 1.6: Basic approximation of lattice [21].

A lattice is a grid-like pattern of connecting points that has its sites spaced out in a periodic fashion as seen in Figure 1.6. A site can only interact with a neighboring site using the line segment,

or link, connecting the 2 sites. Thus, a single site can only interact with up to 4 other sites in a 2D lattice as seen in Figure 1.6, or 6 in a 3D lattice. A calculation using a lattice is limited by several factors.

First, the more sites or links between sites that are used in a simulation, the more processing power is required to simulate the lattice and any interactions.

Second, the complexities of the interactions between sites require more processing power to calculate and simulate the potential effect on the lattice. The complexities of using an Ising model to illustrate the interactions between sites is far less than using quantum field theory [22].

Third, lattice simulations can describe a system in thermodynamic equilibrium, while dynamical processes are not well captured by this technique.

Despite the potential problems of using lattices, they are fairly simple to construct and simulate using modern computer systems.

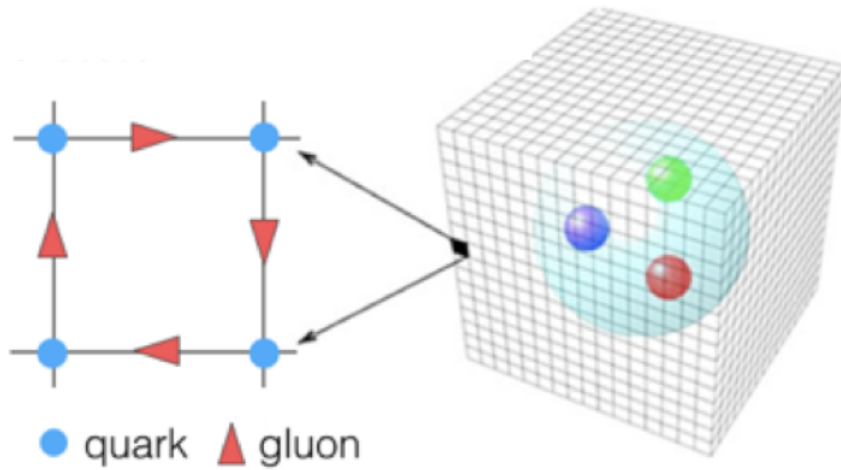


Figure 1.7: Physical representation of quarks and gluons on a lattice [23].

Using lattices to simulate QCD mainly comes down to the implementation of quarks and gluons into the sites and links of the lattice. Similar to Figure 1.7, we can construct a lattice that has its sites correspond to quarks, while the links are represented by gluons. This is very similar to the physical representation of quarks and gluons in a particle or in a QGP. However, because of the

nature of the lattice, the distance between the quarks (sites) are kept at a constant length (link) to each other. Thankfully, the “strength” of the links can be adjusted to approximate the effect that gluons would have on quarks. Furthermore, the lattice can be constructed so that the temporal dimension is equivalent to a temperature dimension.

As mentioned before, lattice simulations are very taxing for computers, even high performance super-computers. A formula to approximate the computational “cost” of the QCD calculation is

$$cost \approx \left(\frac{L}{a}\right)^4 \frac{1}{a} \frac{1}{m_\pi^2 a} \quad (15)$$

where a is the lattice spacing between two sites (the link), and L is the length of a side of the grid [21]. The first factor of Eq. 15, $(\frac{L}{a})^4$, corresponds to the number of lattice sites [21]. The more lattice sites for a given length, the more expensive the simulations. The second and third factors, $\frac{1}{a} \frac{1}{m_\pi^2 a}$, are due to the increase in computational effort needed when approaching critical points of a theory, simply called “critical-slowness”, of the algorithms used for the simulations [21]. From this formula, it is apparent that the most important factor of the cost of the simulation is the lattice spacing. Thus a balance must be made between having a small enough lattice spacing to work as close as possible to QCD in the continuum limit, while large enough to ensure that the simulation can finish in a reasonable amount of time. For example, a full mapping of a single order parameter on a sufficiently large sized volume takes about 100 million core hours on a super-computer facility.

Lattice QCD (LQCD) simulations of baryons under extreme temperature and/or density are fundamental to our understanding of matter in heavy-ion collisions. Due to QCD’s non-perturbative nature, relying on models using lattice simulations becomes necessary. LQCD simulations, such as those from the Swansea collaboration, seem to indicate a property of parity partner particles, known as *parity doubling* [24]. Parity doubling occurs when a baryonic system undergoes chiral symmetry restoration at high temperatures, at which positive and negative parity baryon partners appear degenerate. Since their masses have approached each other to the point that they are indistinguishable, this is referred to as parity doubling. As temperatures decrease, chiral symmetry

is broken, showing the positive-parity baryons to have a typically lower mass than their negative-parity partners.

It is important to note that the particles in these calculations still have a relatively large mass. This is different from the chiral symmetry restoration effect which causes the mass of some particles to reduce to 0 at high temperature. Instead, for these calculations, the mass of the negative-parity baryons decreases to almost match their positive-parity partners, but they still have a non-zero mass.

The FASTSUM collaboration used lattices that were specifically constructed for spectral studies of QCD at non-zero temperature [24]. To achieve this, lattices had to be constructed in a very anisotropic fashion, with $a_\tau/a_s \ll 1$ where a_τ is the lattice spacing in the temporal direction, while a_s is the one in the spatial direction [24]. Temperature is varied using the fixed-scale approach, namely, changing the number of time slices N_τ via the standard relation, $T = 1/(a_\tau N_\tau)$ [24]. The critical (crossover) temperature, T_c , is found to be $T_c = 185(4)$ MeV according to the inflection point of the renormalized Polyakov loop [24]. Due to the light quarks not having their physical masses ($m_\pi = 384(4)$ MeV), the critical temperature is higher than in nature [24]. Despite this, strange quarks have the mass at their current quark mass [24].

While these calculations are interesting, as illustrated in Figure 1.8, they come with two major drawbacks. First, because they used mass values for the light quarks that were much higher than what is found in nature, the main results of the manuscript might not be fully realistic [24]. Second, most of the particles listed in the study are very difficult if not impossible to study given current detector technology available at ALICE [25]. Either the decay products of the various particles are neutral and can not be detected by the ALICE tracking detectors, or they are produced with such a low yield that they can not be used for such a detailed analysis. With that said, the only viable particle that could be used for analysis would be the negative parity Ξ resonance, also called the $\Xi(1820)$ resonance particle.

One of the major innovations of these lattice QCD calculations is the introduction of the quasi-order parameter R , which is a ratio of the correlators ($G_\pm$) for both parity particles over their

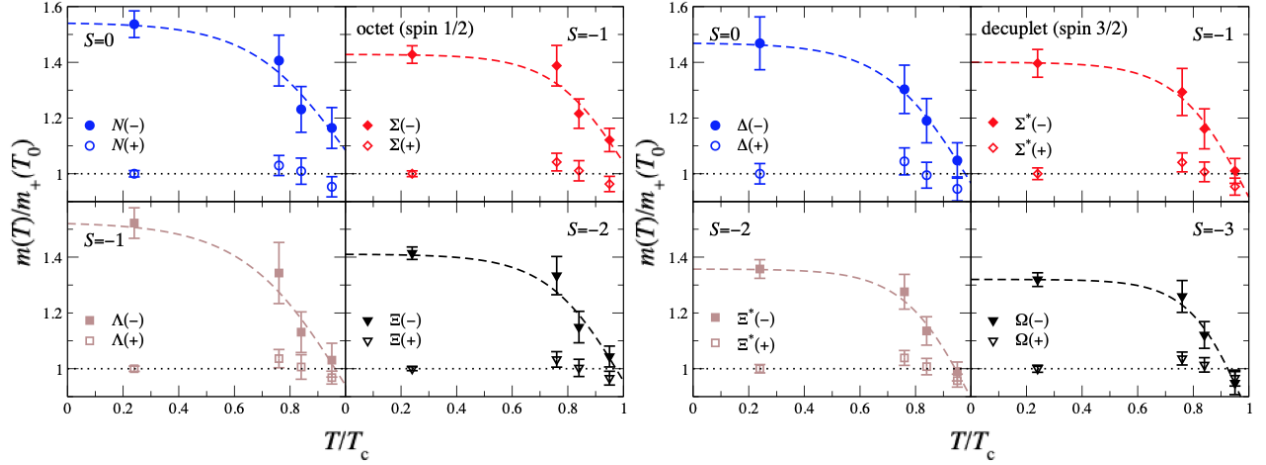


Figure 1.8: Temperature dependence of masses based on lattice QCD calculations. Masses are normalized using m_+ at the lowest temperature, $m_{\pm}(T)/m_+(T_0)$. Baryons with various spin states, octet (left) and decuplet (right), and strangeness content are calculated in the hadronic phase. Positive- (negative-) parity masses are indicated with open (closed) symbols [24].

corresponding errors [26][27]

$$R(\tau) = \frac{G_+(\tau) - G_-(\tau)}{G_+(\tau) + G_-(\tau)}, R = \frac{\sum_n R(\tau_n)/\sigma^2(\tau_n)}{\sum_n 1/\sigma^2(\tau_n)}. \quad (16)$$

R , as it is defined in these calculations, would correspond to 0 in the case of parity doubling and 1 in the case of the single states satisfying $m_- \gg m_+$ [24]. As seen in Figure 1.9, as the temperature increases (or more specifically the ratio of temperature over critical temperature), the value of R decreases from 1 to 0 in a fashion similar to phase transitions of matter. Furthermore, this behavior coincides with the crossover behavior of the de-confinement transition [24]. Interestingly, R is still slightly higher than 0 at the higher temperatures studied for baryons with larger strangeness, such as the Ω ($S=-3$) [24]. Because the high mass of the strange quark (in comparison to the light quarks) can break chiral symmetry, the projected effect is expected to vanish as $m_s/T \rightarrow 0$ [24]. This has lead the Swansea collaboration to conclude that the interpretation of parity doubling due to chiral symmetry restoration in the baryon sector is valid [24].

Despite the experimental challenges, the lattice QCD calculations show great expectations for

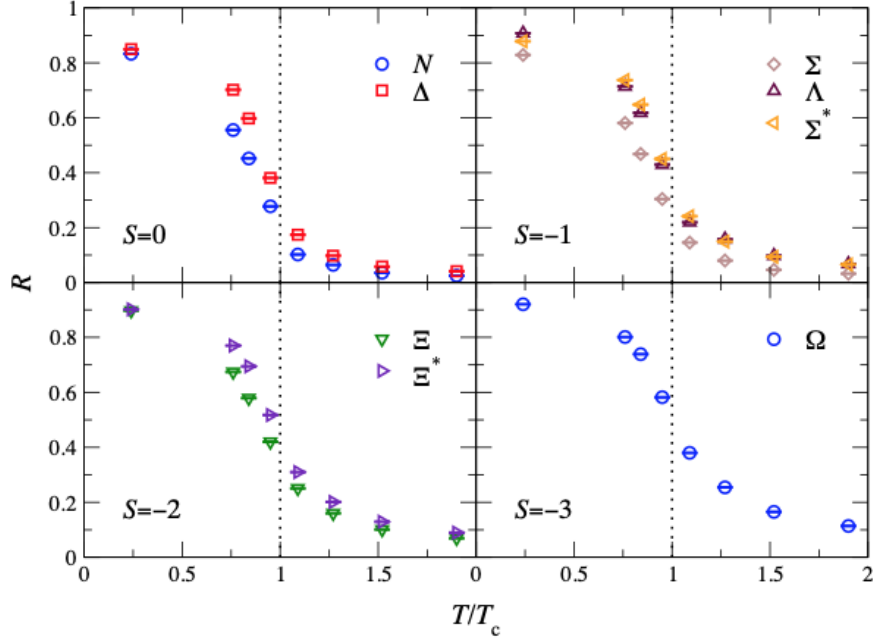


Figure 1.9: Temperature dependence of the ratio R , see Eq. 16. [24]

potential parity doubling results. As can be seen in the Figure 1.8, the ratios of mass of a negative parity particle over the mass of the positive parity particle at T_0 show that negative parity particles decrease in mass as they approach a critical temperature (ratio of T/T_c). This is seen for different spin states, octet states with spin 1/2 on the left and decuplet states with spin 3/2 on the right, and different strange quark contents, $S=0$ to $S=-3$. This parity doubling result is crucial to chiral symmetry restoration. Just as parity partners become indistinguishable from each other during this phase transition, also particles with different chirality become indistinguishable from each other during the phase transition. Thus, if a particle can be found that lowers its mass as it reaches the phase transition, then this would prove parity doubling. Furthermore, this would prove the first major observable evidence of chiral symmetry restoration, further leading to acceptance of QCD calculations and theories.

2 A Large Ion Collider Experiment and the Large Hadron Collider

A Large Ion Collider Experiment (ALICE) is one of the major experiments at the Large Hadron Collider (LHC) at CERN. At 27 km in circumference and 100 m underground, as seen in Figure 2.1, the LHC is the world's largest and most powerful particle accelerator and collider.

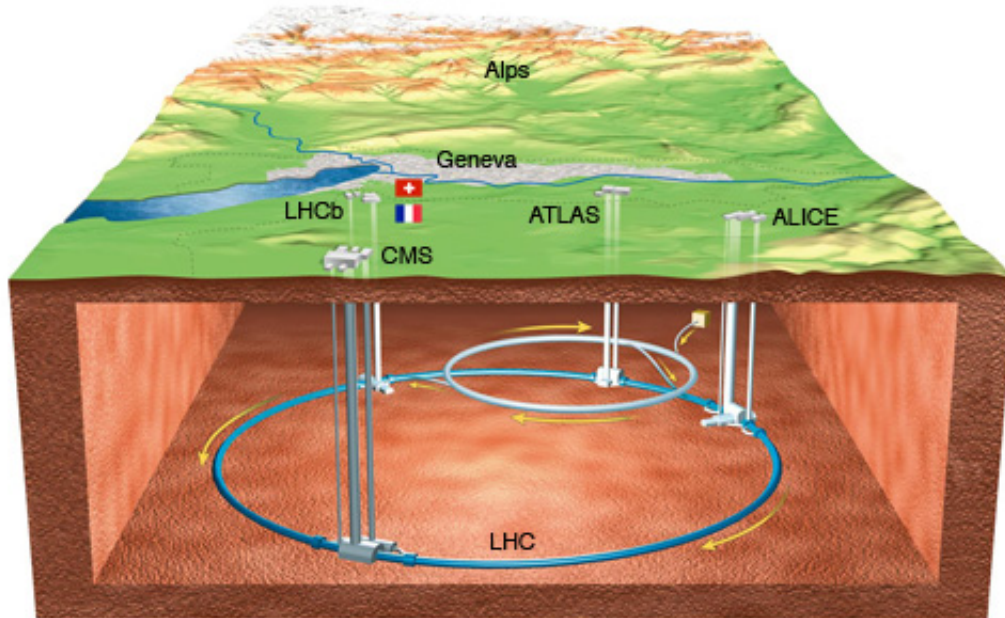


Figure 2.1: View of CERN and multiple experiments above and below ground.

ALICE's purpose is to understand the nature of the QGP and to further investigate its phase transition to hadronic matter. Thus ALICE is designed to be able to handle the large particle multiplicities associated with Pb-Pb collisions (which can reach up to 2000), while also observing as many QGP-related observables as possible. When compared to other experiments such as ATLAS, ALICE is able to provide excellent Particle Identification (PID) performance at low transverse momenta and tracking with very low material budget at mid-rapidity [8] [28]. The PID for several particles can be obtained by using a combination of various detectors from ALICE, some of which use different techniques and have been optimized for different momentum regions [8] [28].

As seen in Figure 2.2, ALICE is composed of 18 detector subsystems that can be categorized

into three main types of detectors based on their position or function.

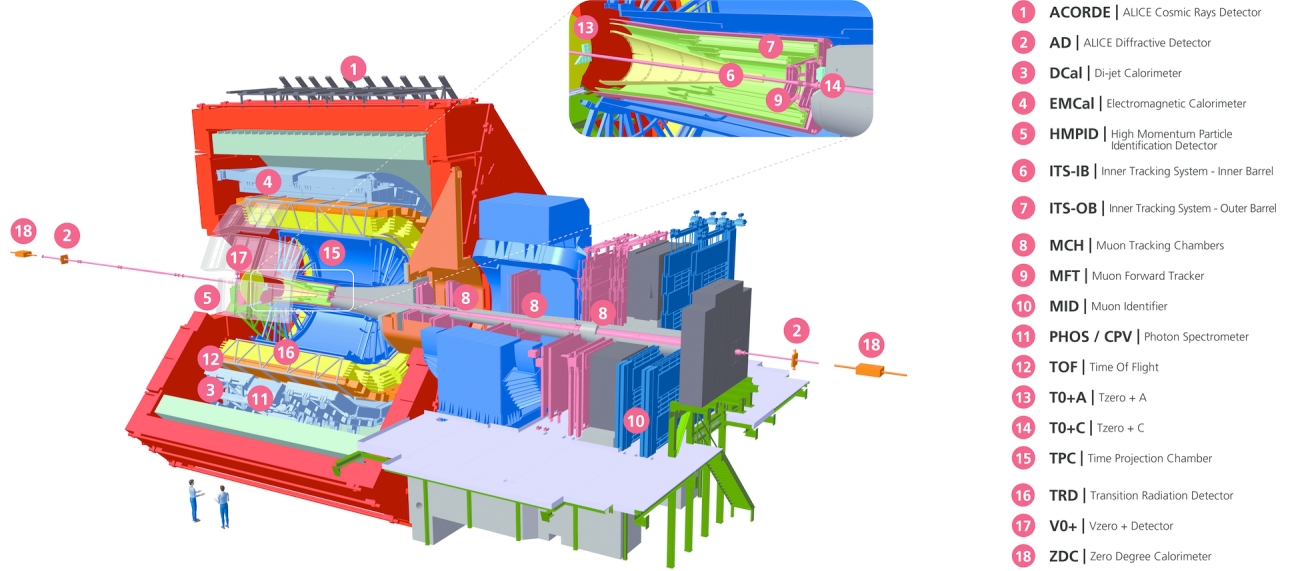


Figure 2.2: The ALICE experimental setup and detectors [29].

First, the subsystems located in the center of ALICE are referred to as central-barrel detectors. Central-barrel detectors are positioned inside of a large solenoid magnet that provides a constant 0.5 T magnetic field parallel to the beam and covers a pseudo-rapidity range of $-0.9 < \eta < 0.9$ [28]. Pseudo-rapidity (η), as it is used in particle physics, is a coordinate system used to describe the angle of a particle relative to the beam axis using the equation:

$$\eta = -\ln\left[\tan\left(\frac{\theta}{2}\right)\right] \quad (17)$$

where θ is the angle between the momentum of the particle and the positive direction of the beam axis. The more perpendicular the angle of the particle's momentum is to the beam axis ($\theta \sim 90^\circ$), the smaller the pseudo-rapidity ($\eta \sim 0$). The primary purposes of these central-barrel detectors are mainly particle identification, tracking, vertex reconstruction, and momentum measurement. Several detectors are listed below in order from the interaction region to the outer region of the detector:

1. Inner Tracking System (ITS)
2. Time Projection Chamber (TPC)
3. Transition Radiation Detector (TRD)
4. Time Of Flight (TOF)
5. High Momentum Particle Identification Detector (HMPID)
6. PHOton Spectrometer (PHOS)
7. ElectroMagnetic CALorimeter (EMCAL)

Do note that the HMPID, PHOS, and EMCAL have limited azimuthal range in the mid-rapidity region [28]. This is either due to budget limitations or the infrastructure of ALICE limiting the space available for the detectors.

Second, the muon spectrometer is located in a forward psuedo-rapidity range of $-4.0 < \eta < -2.5$ and is made up of a dipole magnet and tracking/trigger chambers [8]. The muon spectrometer can be seen as the components to the right of ALICE in Figure 2.2. The muon spectrometer has been optimized and configured to measure muons and to reconstruct heavy quark resonances such as the J/Ψ through their $\mu^+\mu^-$ decay channels [8] [28].

Third, forward detectors are placed in the high pseudo-rapidity area of the beam pipe, some of which can be seen in the small area on the top right of Figure 2.2. Forward detectors are used to measure and to trigger on global event characteristics [8].

1. Time Zero (T0) measures the start time of events, see Section 2.2, with a precision of the order of tens of picoseconds, as needed by the TOF.
2. VZERO (V0) is used to measure charged-particle multiplicities, see Section 2.5, at forward rapidities, trigger minimum bias events, see Section 2.3, and reject backgrounds coming from the beam-gas interaction.

3. Forward Multiplicity Detector (FMD) gives multiplicity information, see Section 2.5, and covers a large fraction of the solid angle ($-3.4 < \eta < -1.7$ and $1.7 < \eta < 5$).
4. Photon Multiplicity Detector (PMD) measures the spatial distribution of photons on an event-by-event basis in the $2.3 < \eta < 3.7$ region.
5. Zero Degree Calorimeter (ZDC) is used to measure and trigger on the impact parameter, see Section 2.5.

The TPC and TOF are the main detectors used for this analysis due to their efficiency and importance in reconstruction and tracking of charged particles. Some of the other detectors used for this analysis, such as the ITS and V0, will be discussed in Section 2.3.

2.1 Time Projection Chamber

The ALICE Time Projection Chamber (TPC), as shown in Figure 2.3, is the largest of its kind, with a volume just under 90 m^3 [25]. Its main purpose is the measurement of charged particle momentum, along with particle identification and vertex determination. In order to perform efficiently for high multiplicity Pb-Pb collisions, it was constructed in a cylindrical shape with a length of 5 m, an inner radius of 85 cm, and outer radius of 250 cm. The gas volume of the TPC is filled with 90% Ne and 10% CO_2 [25].

When a charged particle enters the TPC and ionizes the gas, the released electrons drift to the endplates of the detector under a 400 V cm^{-1} electric field applied uniformly in the z (beam) direction. The design of the endplates is based on the Multi-Wire Proportional Chamber (MWPC) technique with cathode pad readout. Electrons interacting with the grid of charged wires will cause the electrons to avalanche, resulting in a charge as an electronic signal [25] [31].

The central-barrel detectors at ALICE are inside a 0.5 T uniform solenoidal magnetic field [28]. Due to this large over-arching magnetic field and the large electric field of the TPC, charged particles will travel in a path known as helical motion. While the direction of the charged particle is determined by the initial kinematics and the electric field, the circular motion is dictated by the

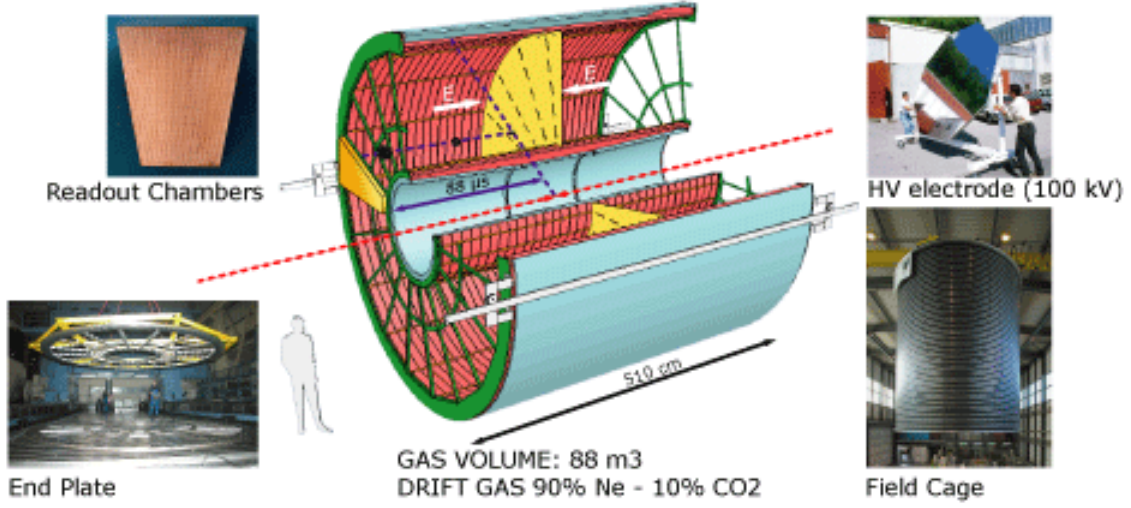


Figure 2.3: Schematic view of TPC with several components [30].

magnetic field as seen in Eq. 18

$$\vec{F}_{total} = q\vec{E} + q\vec{v} \times \vec{B}. \quad (18)$$

Assuming in the best case that the magnetic field is perpendicular to the momentum of the charged particle, the magnetic component of Eq. 18 can be simplified to

$$qvB = \frac{\gamma mv^2}{r}. \quad (19)$$

By replacing the mass and velocity terms with the relativistic transverse momentum ($p_T = \gamma mv_T$), Eq. 19 can be further simplified to

$$qB = \frac{p_T}{r}. \quad (20)$$

Using Eq. 20, it is possible to determine some features of charged particles in a magnetic field. Charged particles with higher momentum will have a larger radius of curvature for the track, thus

being less impacted by the magnetic field. Charged particles with lower momentum will have a smaller radius of curvature for the track, thus being more impacted by the magnetic field. The magnetic field is so strong that low momentum charged particles, normally around $p_T \leq 100 \text{ MeV}/c$, are curved to the point that they do not reach most of the central-barrel detectors [32].

To identify a particle using the TPC, the energy loss per unit length within the matter crossed by the charged particle (simply called dE/dx) must be observed. Energy loss in this case can be described by Bethe-Bloch parameterization (see Eq. 21) since it is key for the identification technique used. According the Bethe-Bloch parameterization, the mean specific energy loss is [8]:

$$-\langle \frac{dE}{dx} \rangle = k_1 \cdot z^2 \frac{Z}{A} \cdot \frac{1}{\beta^2} \left[\frac{1}{2} \ln(k_2 \cdot m_e c^2 \cdot \beta^2 \gamma^2) - \beta^2 + k_3 \right] \quad (21)$$

where

β : the speed of the particle as a fraction of the speed of light, c

γ : the Lorentz factor $1/\sqrt{1-\beta^2}$ (note that $\beta\gamma = p/Mc$)

p : ionizing particle momentum

M : ionizing particle mass

Z : atomic number of the ionized gas (Ne/CO₂)

A : mass of the ionized gas (g/mol)

m_e : electron mass

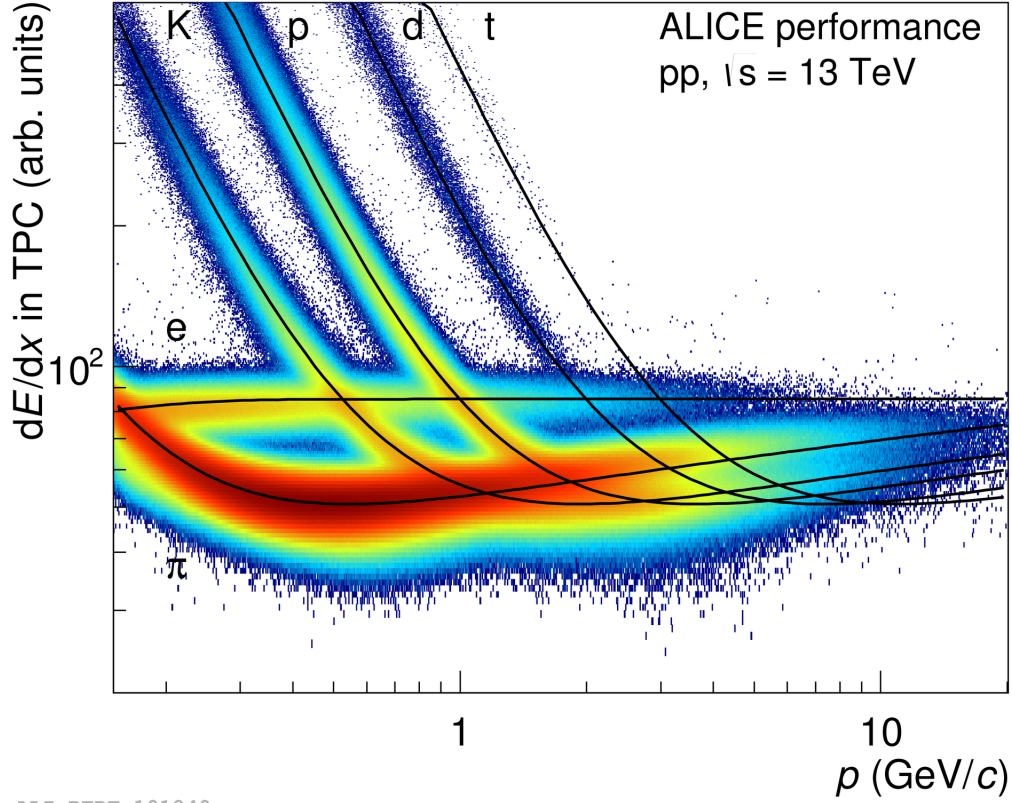
z : electric charge of the ionizing particle in units of electron charge e

k_1, k_2, k_3 : constants depending on the ionized medium

In order to match a corresponding track within the TPC to its particle signature, a parameter known as n_σ (nSigma) is calculated as the difference between the measured value from the TPC and the theoretical value from the Bethe-Bloch parameterization divided by the resolution of the TPC (σ_{TPC}).

$$n_\sigma = \frac{(dE/dx)_{measured} - (dE/dx)_{Bethe-Bloch}}{\sigma_{TPC}} \quad (22)$$

In short, the closer n_σ is to 0, the closer a track is to a corresponding Bethe-Bloch value and the



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Figure 2.4: TPC dE/dx for 0.2T run in Run2. Solid lines correspond to signal calculated from Bethe-Bloch formula. Intensity corresponds to how many signals of dE/dx for specific momentum p are measured.

higher the certainty of the track being correctly identified as a particular type of particle. This does lead to some complications when the Bethe-Bloch parameterizations of different particles intersect with each other as seen in Figure 2.4, but limiting the momentum range of particle identification or using strict n_σ cuts can reduce the probability of false identification.

Despite the complexities, the TPC is still very accurate in its ability to identify particles on a track-by-track basis, especially in the low momentum range with a 3σ separation from other particles. These criteria allow for the identification of charged pions, charged kaons, and protons within the transverse momentum ranges of 0.25-0.70 GeV/ c , 0.25-0.45 GeV/ c , and 0.45-0.90 GeV/ c , respectively [33]. However, due to the intersection of dE/dx for the “mid” transverse momentum,

the TPC is not able to accurately separate these particles in the approximate momentum ranges of 1.0-3.0 GeV/c [33]. Thankfully, in the higher transverse momentum ranges the TPC can again separate the pions, kaons, and protons with a 1-3 σ separation from other particles in the transverse momentum ranges of 3.0-20.0 GeV/c, 4.0-20.0 GeV/c, and 4.0-20.0 GeV/c respectively [33]. This region of the dE/dx curve, which is partially shown in Figure 2.4, is called the relativistic rise region due to the “rise” of the Bethe-Bloch values for the particles.

2.2 Time Of Flight

The Time Of Flight (TOF) system was designed to improve the particle identification capability of the ALICE experiment. Due to the inability of the TPC and ITS to successfully identify particles in the intermediate momentum range, the TOF is constructed to cover the missing momentum range as shown in Figure 2.5. This allows for the entire ALICE detector to cover a wide range of transverse momentum, from 150 MeV/c² to 20 GeV/c², for PID.

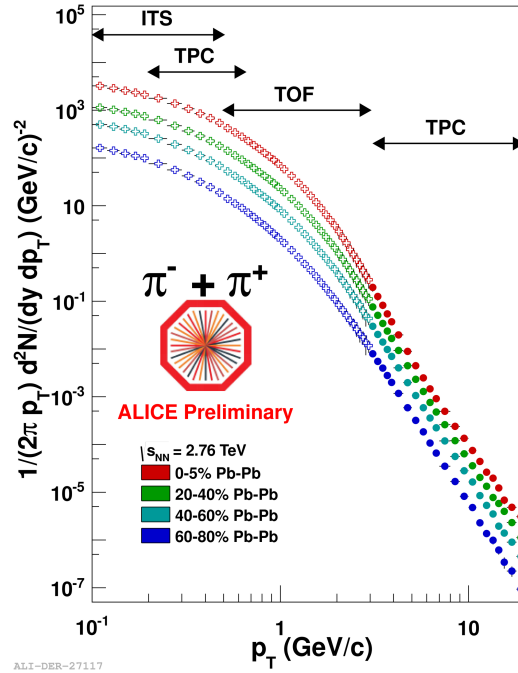


Figure 2.5: Invariant yield for charged pions, $\pi^+ + \pi^-$, as a function of p_T for different centrality classes measured with 3 different PID detectors in 4 different p_T regimes.

The TOF is constructed as a cylindrical array at 3.7 m radius from the beam axis. This range allows for it to cover polar angles (θ) between 45 degrees and 135 degrees over the full azimuth ($2\pi\phi$) [34]. As seen in Figure 2.2, the TOF is positioned between the TRD and EMCAL.

The TOF contains 1638 Multigap Resistive Plate Chamber (MRPC) strips in a modular structure of 18 sections in ϕ and 5 modules in each section along the beam axis [34]. The MRPC is a stack of resistive glass plates with a high voltage applied to the external surface. When a charged particle ionizes the gas between the gaps of the plates, the resulting electron signal is amplified by the high electric field. This causes an electron avalanche, where the initially freed electrons are accelerated by the electric field and collide with other atoms in the gas and cause more ionization, thus causing the signal to amplify. While the resistive plates stop the avalanche development in each gap, they are transparent to the signal produced by the electrodes. Thus, as shown in Figure 2.6, by determining the time it takes for a particle to cross from the collision vertex to the TOF detector, it is possible to determine the mass of the particle using an equation such as

$$m = \frac{p}{c} \sqrt{\frac{c^2 t^2}{l^2} - 1} \quad (23)$$

where m is the mass of the particle, p is its momentum (measured via the TPC and ITS detectors), t is the time of flight, and l is the track length. It should be noted that the time that the particle starts at from the primary vertex is determined by the T0 detector (which is an array of two Cherenkov counters called T0-A and T0-C described further in Section 2.3) that is a part of the forward detectors [35]. With this in mind, the more glass plates that are available in the system would correspond to more signal and thus higher efficiency, and a smaller gap width between these plates would correspond to less uncertainty in time and thus better time resolution [34]. The TOF was optimized to find a balance between the number of MRPC strips needed to obtain a signal with high efficiency, while having a small enough gap between the plates to reach a time resolution needed for identifying particles.

Using the previous equation of mass, it is possible to determine the difference in signals between

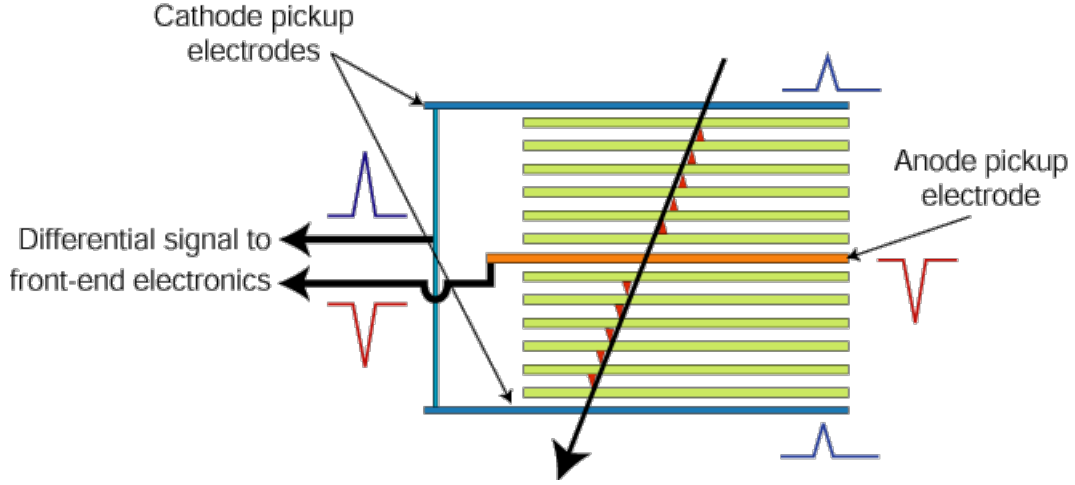


Figure 2.6: Illustration of a signal passing through the TOF [36].

two particles of different masses using the TOF. To determine the difference between two time of flight signals (δt) from two particles that are relativistic ($p \gg mc$), have the same momentum ($p_1 = p_2 = p$), the same track length ($l_1 = l_2 = l$), but different masses, the equation

$$\delta t = \frac{l}{2c} \frac{(m_1^2 + m_2^2)}{p^2} \quad (24)$$

can be used. Do note that momenta are in units of MeV/c , masses are in MeV/c^2 , the speed of light c is in m/s , track lengths are in m , and δt is in units of seconds. Furthermore, by including the

time resolution of the detector (dt), the equation to distinguish two particles in the TOF system can be approximated by:

$$n_{d,1-2} = \frac{\delta t}{dt} = \frac{l(m_1^2 + m_2^2)}{2cp^2dt}. \quad (25)$$

This Eq. 25 is also known as the “separation power” of two particles in the TOF and can be illustrated using Figure 2.7. This separation power is very similar to the n_σ discussed in Section 2.1 for the TPC, even having some of the same notation. However, there are a few key differences. Most notably, the n_σ for the TPC is a comparison between a model and the actual measurement of a signal, while the separation power of the TOF is a comparison of two particles with no model. Do note that there is a version of n_σ for the TOF that is calculated the same way as n_σ for the TPC, and is used in a similar fashion. This n_σ will be discussed in more detail in Section 6.2.

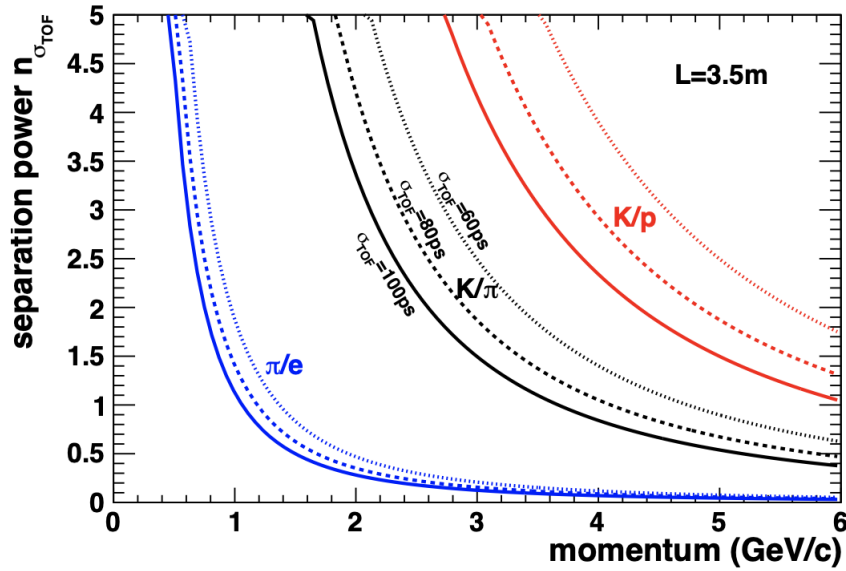


Figure 2.7: Separation power (n_σ) of particles vs. momentum for TOF with track length $L = 3.5$ m and three separate time resolutions (σ_{TOF} (dt) = 60, 80, and 100 ps) [37]. Assume infinite precision on momentum and track length measurement.

Even though the time difference (δt) between two particles is the main component used in separating particles, the velocity $\beta = v/c$ is the value determined for each particle and signal from the TOF. This is shown in Figure 2.8 where most particles can be separated at the lower momentum

range, while at higher momentum ranges the signals become indistinguishable from each other.

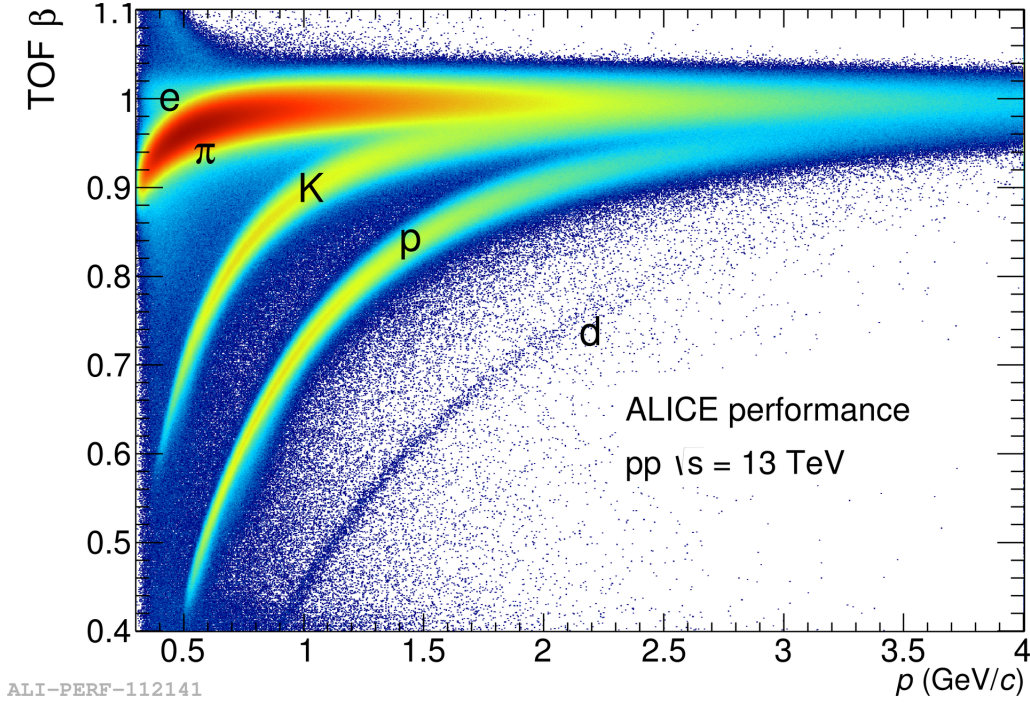


Figure 2.8: TOF β vs. momentum performance plot in p-p collisions at $\sqrt{s} = 13$ TeV.

Because the primary purpose of the TOF is to determine the difference between the signals of $\pi - K$ and $K - p$, the TOF is optimized so that for a 4 meter particle track, a 3σ separation between π and K is achievable at the higher end of the expected momentum range of the TOF. Using a momentum of 2.5 GeV/c (and with the mass of the π and K known), the time resolution of $dt < 100$ ps is achieved using the current technology available at ALICE [34].

Misidentification of particles occur at the higher momentum range of the TOF. This is because the time difference between two particles will be comparable to the time resolution of the TOF. The lower momentum threshold is limited by the curvature of tracks in the magnetic field. With a magnetic field of 0.5 T and assuming a path length of $L = 3.5$ m, particles that do not have momentum larger than 300 MeV/c can not even reach the TOF [37].

2.3 Other Detectors

As mentioned before, the TOF and TPC are the main detectors used for this analysis. However, some other detectors are also used in conjunction with these detectors.

The Inner Tracking System (ITS) is a series of six cylindrical layers of silicon detectors, located at a distance between 2 and 43 cm from the center of the ALICE detector [25]. The number, position, and segmentation of these layers have been optimized for efficient track finding and high impact-parameter resolution (see Section 2.5 for more details on impact parameter) [25]. For our purposes, it should be mentioned that the outer radius of the ITS is determined by the need to match tracks with those from the TPC [25]. This allows for further efficiency in calculating the tracks of charged particles and better resolution for determination of parameters such as the primary vertex, also known as PV. The primary vertex is simply the position of the collision from which all primary particles originate. Using the ITS as a vertex finder is important in determining the starting point of all tracks, this determines if a track did come from the primary vertex, or if it came from a secondary decay vertex such as the decay vertex of the Λ as discussed in Section 4.1.4.

The V0 detector is a small angle forward detector consisting of two arrays of scintillator counters [25]. The arrays, called V0A and V0C, are installed on either side of the ALICE intersection point in pseudo-rapidity ranges $2.8 \leq \eta \leq 5.1$ and $-3.7 \leq \eta \leq -1.7$ respectively [25]. The primary purpose of the V0 detector is to provide minimum-bias triggers for the central-barrel detectors in p-p and A-A collisions [25]. These triggers are activated by particles originating from the initial collisions and from secondary interactions in the vacuum chamber elements [25]. Because the dependence between the number of registered particles on the V0 arrays and the number of primary emitted particles is linear, the V0 serves as an indicator of the centrality of a Pb-Pb collision via the multiplicity recorded in the event [25]. See Section 2.5 for more detail on centrality.

The purpose of these triggers is to categorize events according to certain criteria that these triggers are primed to look for during the experiment. The minimum bias (MB) trigger is the simplest trigger in that if an event produces a signal at the V0A or V0C, the event is categorized

as being a part of the minimum bias data set. Because of these triggers, it is possible to determine if a collision did occur, or if the signal observed in the detectors came from beam-gas interactions. Other triggers used for Pb-Pb data include those that require a signal above a specific threshold in the V0 detector to record preferential central collisions [8]. The central trigger (kCent) corresponds to events with centrality between 0-10 %, the semi-central trigger (kSemiCent) corresponds to events with centrality between 30-50 %, and the minimum bias trigger (kINT7) corresponds to event with centrality between 0-90 %. See Section 2.5 for more detail on centrality.

The T0 detector is two arrays of Cherenkov counters that have two main purposes. First, the T0 detector is used to determine the start time (T0) of the TOF detector, indicating the time that the collision occurred [25]. Second, the T0 detector is used to determine the vertex position for each collision [25]. The two Cherenkov counters, called T0-A and T0-C, are positioned in high pseudo-rapidity ranges so as to determine the start time of the collision independent of the position of the vertex. The T0-C is placed 72.7 cm from the center of the ALICE detector in a pseudo-rapidity range of $-3.18 \leq \eta \leq -2.97$, while the T0-A is placed 375 cm from the center of the ALICE detector in a pseudo-rapidity range of $4.61 \leq \eta \leq 4.92$ [25].

2.4 Data Sets to be Analyzed

Between the end of 2009 and the beginning of 2013 (referred to as Run1), ALICE was able to accumulate data from the entire first phase of LHC operation. This corresponds to p-p collisions at center of mass energies corresponding to $\sqrt{s} = 0.9, 2.76, 7, \text{ and } 8 \text{ TeV}$, p-Pb collisions at center of mass energies per nucleon of $\sqrt{s_{NN}} = 5.02 \text{ TeV}$, and Pb-Pb collisions at center of mass energy per nucleon of $\sqrt{s_{NN}} = 2.76 \text{ TeV}$. Having successfully carried out data collection for each of these energy systems, ALICE was ready for data collection at higher energies. Between 2015 and the end of 2018 (Run2), beam energies were increased for collision systems such as, but not limited to, p-p collisions at $\sqrt{s} = 13 \text{ TeV}$, p-Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$, and Pb-Pb at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$. These three energies and systems specifically are the main focus of this dissertation mostly because of two factors. First, the higher energies of these collision systems contained higher statistics due

to more efficient data acquisition systems and a higher collision rate at the LHC, thus making the observation of $\Xi(1820)$ more feasible. Second, the higher energies of Pb-Pb would correspond to higher particle multiplicity per event, thus higher chance of chiral symmetry restoration being noticeable in the analysis.

Data from p-p collisions at $\sqrt{s} = 13$ TeV were selected for two major criteria. First, mainly due to the fact that the $\Xi(1820)$ is a very rare particle to find, approximately 1 $\Xi(1820)$ is formed in every 31000 p-p $\sqrt{s} = 13$ TeV events (see Section 4.4). Thus a large number of events is necessary in order to find and analyze it. With that in mind, the p-p $\sqrt{s} = 13$ TeV data corresponds to approximately 1.97 billion events, the largest number of events for any system and energy available at ALICE.

Second, a baseline is required to determine the properties of the $\Xi(1820)$ and determine if any feature changes in heavy ion collisions where the formation of a QGP and thus the restoration of chiral symmetry is more likely. According to the particles listed in the annually updated summary publication booklet of the international Particle Data Group (PDG), the last analysis of the $\Xi(1820)$ was performed in 1999, while previous measurements ended in 1987 [12]. Furthermore, these calculations have very large errors for the mean and width such that the $\Xi(1820)$ is marked as a 3 out of 4 star particle in the PDG, representing the fact that while the $\Xi(1820)$ is known to exist with some approximated features, some of the features such as the branching ratios are not well measured [12]. Since the p-p $\sqrt{s} = 13$ TeV analysis will provide smaller errors than those stated in the PDG, it is reasonable to use these numbers for the baseline calculations instead of the PDG value. Furthermore, measuring the baseline using ALICE allows us to determine if there are any detector effects that would affect our measurement in the absence of parity doubling.

The p-Pb $\sqrt{s_{NN}} = 5.02$ TeV and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data were selected for similar reasons as the p-p $\sqrt{s} = 13$ TeV data. Of the p-Pb and Pb-Pb data from Run2, $\sqrt{s_{NN}} = 5.02$ TeV data had sufficient statistics collected for each system to assume that a $\Xi(1820)$ signal could be extracted. Furthermore, unlike the p-p $\sqrt{s} = 13$ TeV data, Pb-Pb data does have the likelihood of producing the QGP and chiral symmetry restoration in the most central collisions. If a mass shift

or width broadening signal is found in Pb-Pb data that is significantly higher than the signal from p-p $\sqrt{s} = 13$ TeV, this information would be instrumental in proving chiral symmetry restoration in the QGP.

It should be noted that the p-Pb $\sqrt{s_{NN}} = 5.02$ TeV data only contains the minimum bias trigger for this analysis. The Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data contains 3 triggers (kCent, kSemiCent, and kINT7) that can be used to correspond to different charged particle multiplicities and centralities. Each centrality trigger for the Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data contains similar statistics for each analysis. The kCent triggered data for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV contains $9.6 \cdot 10^7$ events, the kSemiCent trigger contains $8.4 \cdot 10^7$ events, and the kINT7 trigger contains $1.34 \cdot 10^8$ events. See Section 2.5 for more detail on centrality.

2.5 Centrality

One of the most important parameters for this analysis is the measurement of the multiplicity of the events being analyzed. Multiplicity can be simply defined as the number of particles produced in a collision. The more particles that are being produced in a collision, the more central or violent the system that created them should be. Two protons that hit each other head on should produce more particles than two protons that only graze each other and do not produce any particles.

For our purposes, a value known as the “multiplicity percentile” is calculated for this analysis. For a given multiplicity X , the corresponding multiplicity percentile is the fraction of events that have a multiplicity greater than or equal to X . For example, in the case of 2 protons colliding at about $\sqrt{s} = 13$ TeV energy, an event that produces 3 particles (3 multiplicity) can be considered in the 70 % multiplicity percentile. This means that 70 % of events for this collision system and energy produce 3 particles or more, while the remaining 30 % of events produce less than 3 particles. Similarly, an event that produces 45 particles can be considered in the 0.1 % multiplicity percentile. This means that 0.1 % of events for this collision system and energy produces 45 particles or more, while the remaining 99.9 % of events produce less than 45 particles. The higher the multiplicity percentile (the closer to 0 %), the higher the number of particles that are produced in the event.

This notation is useful for determining how “violent” an event is in comparison to other events. The more violent the collision, the more particles produced, the higher the multiplicity percentile. This allows for further subdivision of datasets in accordance to how high the multiplicity percentile is. More violent, head on collisions between particles that produce a high multiplicity of particles, can be categorized as having a higher multiplicity percentile. Near miss collisions between particles that produce a low multiplicity of particles, can be categorized as having a lower multiplicity percentile. Similarly, events that produce higher multiplicity and are part of the higher multiplicity percentile correspond to collisions that are more head on and violent in their collision.

Centrality is defined as how “central” two opposing nuclei, normally Pb and Pb, overlap with each other. The more “central” a collision is, the more the two particles overlap each other. The less “central” a collision is, the less the two particles overlap. To better analyze centrality, an impact parameter b can be used to further define centrality. b is the length of a 2D vector, connecting the center of the two nuclei which provides a geometric scale of the overlapping region [8]. This impact parameter is illustrated in Figure 2.9.

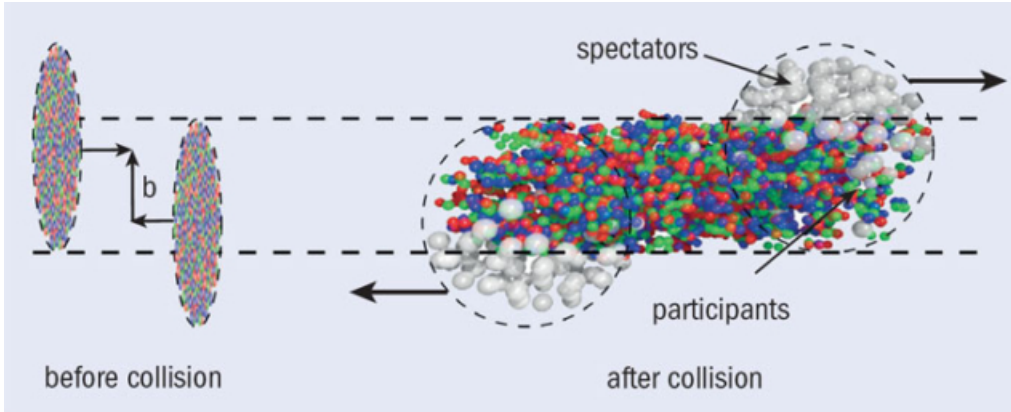


Figure 2.9: Left: Two Lorentz-contracted heavy ions before collision with impact parameter b . Right: The spectator nucleons continue unaffected while particle production takes place in the participant zone [38].

A “central” collision has a small impact parameter (b closer to 0), produces a large volume of QGP, and produces a large charged-particle multiplicity. The “peripheral” collision has a large

impact parameter (b closer to value of 2 radii), produces a small volume of QGP, and produces a small charged-particle multiplicity.

It must be noted that, while multiplicity percentile and centrality may seem to be the same observable, they are not. Multiplicity percentile refers to the number of particles being produced in a collision. Centrality refers to how much overlap two nuclei have over each other in a collision. This can be illustrated using Figure 2.10. Do note that the Negative Binomial Distribution (NBD)-Glauber fit from Figure 2.10 is a fit from Monte Carlo implemented simulations of Glauber models, which assumes that the number of particle producing sources is given by the formula

$$f \times N_{part} + (1 - f) \times N_{coll} \quad (26)$$

where N_{part} is the number of participating nucleons, N_{coll} is the number of nucleon-nucleon collisions, and $f \sim 0.8$ quantifies their relative contributions [39]. The number of particles produced by each source is distributed according to a Negative Binomial Distribution (NBD), parametrized with μ and k , where μ is the mean multiplicity per source and k controls the contribution at high multiplicity [39]. While it is true that the closer each value is to zero, the higher the number of particles produced in a collision, the more central the collision is, there are two main examples that can show the differences.

First, centrality can not be used for p-p collisions in the same way that it is used for Pb-Pb collisions. This is due to the fact that protons are smaller and simpler than Pb nuclei. As stated in Section 1.2, when Pb nuclei reach the highest speeds of the LHC, the Lorentz contraction has caused the normally spherical nuclei to flatten to disk-like circles. The 208 nucleons of the Pb nuclei are now flattened into circles that allow for much simpler, but still complex, collision physics. When protons reach the highest speeds of the LHC, the three quarks in the proton will be flattened along the direction of the beam (momentum) vector. Furthermore, the impact parameter b would be almost impossible to calculate since the protons are so much smaller than the Pb nuclei.

Second, multiplicity percentile and centrality are not exactly dependent on each other. As

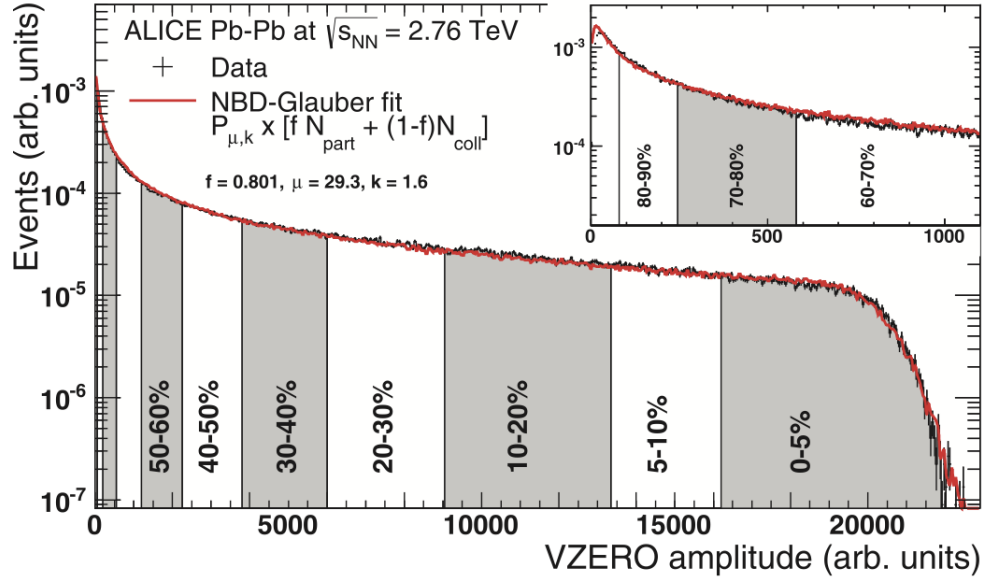


Figure 2.10: Distribution of the sum of amplitudes in the V0M scintillators. Centrality classes indicated by shaded regions. Distribution is fitted with the NBD-Glauber fit shown as red line [39].

stated before, the closer each value is to zero, the higher the number of particles produced, and the more central the collision is. Similarly, the further away from zero each is, the lower the number of particles produced, and the less central the collision is. However, there is no reason for the multiplicity percentile and centrality to have the same percentage value for a given collision. It is entirely possible for a Pb-Pb collision to have a 50 % centrality collision, meaning that 50 % of the area of the two nuclei have overlapped, and for the collision to produce enough particles to be consider in the 40 % multiplicity percentile, meaning that 40 % of events produce the same or more particles.

Furthermore, the impact parameter, which is used to determine the centrality, can not be measured directly. So in practice, the multiplicity percentile is used as a proxy to estimate the centrality percentile.

The centrality estimation and multiplicity calculation used in this thesis is based on the measurements of the signals from the V0 scintillators. To categorize the events, information from the sum of the amplitude in the 2 V0 detectors (called V0A and V0C separately or V0M if summed

together) is used to determine the centrality, as seen in Figure 2.10 for Pb-Pb, and multiplicity percentile, as seen in Figure 2.11 for p-p, of the collisions. Do note the differences between Figure 2.10, which show the centrality classes for Pb-Pb data using V0M, and Figure 2.11, which shows the multiplicity classes for p-p data using V0M. It should be noted that while V0A and V0C are used for p-p and Pb-Pb collisions, only V0A is used for p-Pb collisions due to the “lopsided” nature of the collision. This is because when the proton collides with the Pb nuclei, it will not completely cover the area of the Pb nuclei. Thus, a portion of the Pb nuclei will continue its trajectory seemingly uninfluenced by the proton collision, so only the V0A is needed to determine the multiplicity percentile.

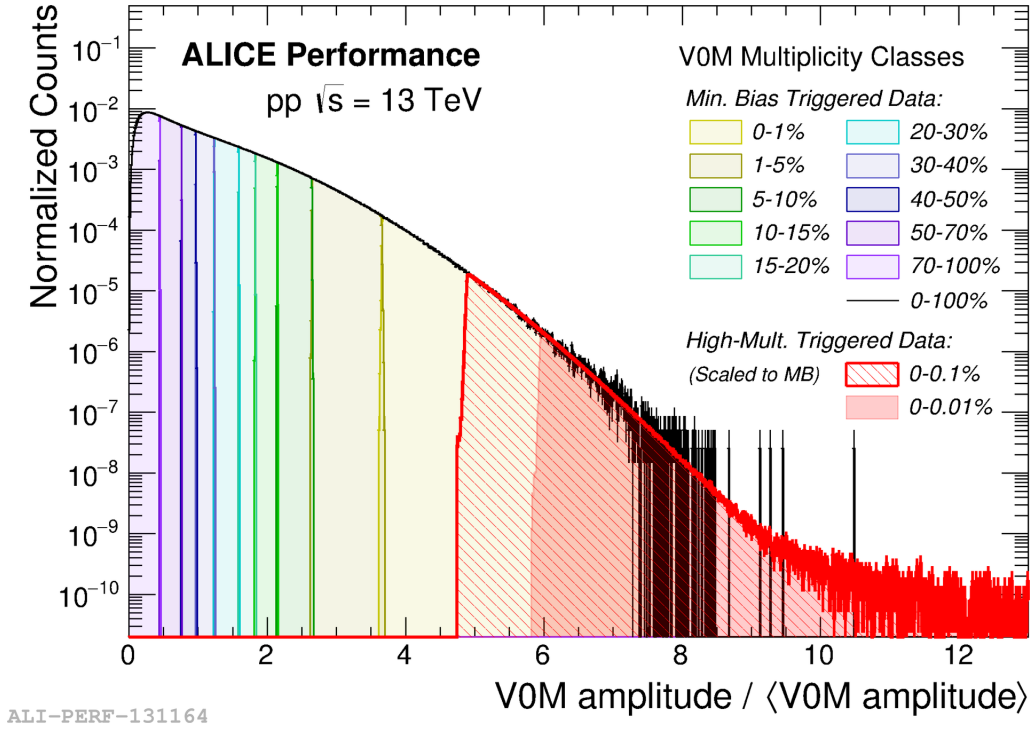


Figure 2.11: V0M/ $\langle V0M \rangle$ normalized distribution for V0M multiplicity classes for minimum-bias (MB) and high-multiplicity (VHM) triggers.

As mentioned in Section 2.4, Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data contains 3 triggers that can be used to correspond to different charged particle multiplicities and centralities. The kINT7 trigger corresponds to the “minimum-bias” trigger, and is used to mark an event as being part of the

“minimum-bias” if the signal from the V0 detectors correspond to the event having a centrality between 0-90 %. The kCent trigger is the “central” trigger, used to mark an event as having a “central” collision with a centrality between 0-10 %. The kSemiCent trigger is the “semi-central” trigger used to mark an event as having a “semi-central” or “peripheral” collision with a centrality between 30-50 %. Each trigger is used to further divide the Pb-Pb data into easily accessible data sets that allow for further analysis.

The multiplicity percentile and centrality classes used in this analysis are described in Table 1. These centrality and multiplicity percentiles correspond to mean charged-particle multiplicity densities ($\langle dN_{ch}/d\eta_{lab} \rangle$) measured at mid-rapidity ($|\eta_{lab}| < 0.5$) [40] [41] [42]. Do note that, because the multiplicity classes and centrality classes do not completely match up with the classes used in this analysis, the weighted average is calculated to determine the mean charged-particle multiplicity densities and the uncertainty from the sources [40] [41] [42].

Multiplicity percentile of p-p $\sqrt{s} = 13$ TeV	$\langle dN_{ch}/d\eta_{lab} \rangle$	sys. uncertainty
0-100 %	7.60	0.14
0-10 %	19.2	0.5
10-30 %	12.0	0.3
30-100 %	4.68	0.16
Multiplicity percentile of p-Pb $\sqrt{s} = 5.02$ TeV	$\langle dN_{ch}/d\eta_{lab} \rangle$	sys. uncertainty
0-100 %	17.4	0.7
0-30 %	32.0	0.8
Centrality of Pb-Pb $\sqrt{s} = 5.02$ TeV	$\langle dN_{ch}/d\eta_{lab} \rangle$	sys. uncertainty
0-90 %	545	6
30-50 %	415	10
0-10 %	1765	36

Table 1: Multiplicity percentile and centrality classes used for p-p collisions at $\sqrt{s} = 13$ TeV, p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [40] [41] [42] [43].

3 Specific Measurements to Address Chiral Symmetry Restoration

As mentioned in the introduction, measuring any signature of chiral symmetry restoration is very challenging. Examination of Figure 1.8 seems to show several possible candidates that could be studied to see if there is any parity doubling that could be measured. However, the experimental limitations in measuring negative parity particles mean that only one set of particles could be measured. The particles labeled $\Xi^*(-)$ and $\Xi^*(+)$, which correspond to the $\Xi(1820)$ and $\Xi(1530)$ resonant states respectively, are the only particles that do not have neutral particles at the end of their decay chains or have difficulty being studied due to low yield. Thus, the focus of this discussion is the only measurable negative parity baryon that was calculated by the Swansea collaboration: the negative parity partner of the resonant Ξ , simply called the $\Xi(1820)$.

3.1 Chiral Partners

Every particle in the standard model of the universe also has an anti-particle with the same mass but opposite quantum numbers. However, resonance particles have the unique feature that, while they have some of the same quantum numbers as their ground state, such as strangeness or isospin, they can have different values for other quantum numbers, such as total angular momentum or parity. Nevertheless, any resonance particle, i.e. excited state, even if it decays shortly after forming, must have a unique set of quantum numbers to be considered an actual state. Thus, while resonance particles are normally identified by their mass and width, the rest of their quantum numbers help further differentiate them from other excited states. Normally, this is seen in an increase of the orbital angular momentum quantum number or a change in the parity of the particle and a difference in mass.

As was seen in Figure 1.8, two particles can be considered “Chiral Partners” if they share all of their quantum numbers except for parity, and are similar to each other in mass. For example, as is seen in Figure 1.8, the particles listed as $\Xi^*(-)$ and $\Xi^*(+)$ (which correspond to the $\Xi(1820)$

and $\Xi(1530)$ particles, respectively) have the same strangeness ($S=-2$), the same isospin ($I_z=1/2$), the same spin ($J=3/2$, decuplet), but they have different masses and different parity. Thus, they can be considered parity partners. Furthermore, in the event that some features between the two particles change according to a phase transition or symmetry breaking, the two particles can be considered chiral partners. For example, if the mass of the negative parity particles decreases to become degenerate with the positive parity particle as temperature increases, the particles are chiral partners. This is because, in the event of chiral symmetry restoration, the two chiral partners will become indistinguishable from each other.

3.2 Relevant Measurements

Chiral symmetry restoration and its effects can be very difficult to detect, let alone analyze. As stated in Section 1.4, normal matter is chirally broken because of its mass.

During the phase transition of normal matter to the QGP, chiral symmetry is restored, thus the differences between left- and right-handed chiral particles and anti-particles vanish. While at first, treating left-handed particles no different than right-handed particles may seem like a minor adjustment to one's equations, the main issue is what this symmetry means for other properties of matter. As stated in Section 1.4, because chiral symmetry is broken due to the introduction of mass, it follows that, for chiral symmetry to be restored, mass must be removed. Thus, if chiral symmetry is restored, the mass of a particle should decrease. Conversely, if a system does have chiral symmetry, that symmetry may be broken due to the introduction of mass.

What are the consequences of particles without mass? While inside the QGP, particles that have similar quantum numbers, such as Ξ and its resonances, will have the same mass and as such will be degenerate when compared to each other. However, because these particles are in the QGP, they will be in a de-confined state such that the quarks and gluons are no longer bound to the colorless groups they normally try to form.

One obvious question about chiral symmetry restoration is whether the particles have zero mass when the symmetry is restored. The phase transition from normal matter to QGP at the LHC is

not a 1st order phase transition, where at one temperature the order parameter changes instantly. This would be similar to how water that is boiling at 100°C goes from a liquid to a gas without any in-between state of matter. Instead, the phase transition is a cross-over, meaning that over a range of temperature, the state of matter can be in a mixed-phase as it transitions from one phase to another. This is somewhat seen in Figure 1.8 where, before the phase transition at the pseudo-critical temperature, only the negative parity particles decrease in mass while the positive parity particles remain independent. Furthermore, Figure 1.9 shows that R does not reach 0 at the phase transition at the pseudo-critical temperature, but continues to drop even until twice the pseudo-critical temperature is reached. Even then, not all values of R reach 0. Despite the fact that the masses of all particles are not zero, the observation that the mass of some particles have decreased to match their chiral partners means that chiral symmetry has at least been “partially” restored. This means that any notable change in mass of the particles would be a signature of chiral symmetry starting to be restored, even if not completely restored. It should be noted that chiral symmetry is still valid even at non-zero mass as long as the mass is negligible compared to the energy of the system.

In the event that the mass of a particle changes, there needs to be a way to observe such a change. However, observing a change that can only be seen for a few fm/c after a collision is no easy task.

Despite the usefulness of resonances, there is one major flaw with using them to determine if chiral symmetry is partially restored. Resonances, and all particles that decay, follow an exponential decay law. This means that, out of all the resonances that could show an effect of chiral symmetry restoration, only some will decay at a lower mass, referred to as “off-shell,” while some will remain until chiral symmetry is completely broken and decay at their regular mass, referred to as “on-shell.” The lifetimes and decay lengths that are used in the discussions about decaying particles only refer to the amount of time, or length if traveling at the speed of light, that the radioactive or exponential decay law needed for the initial distribution of particles to be reduced to 36.79 % (e^{-1}), while the other 63.21 % ($1-e^{-1}$) of particles decay into their daughter particles. Furthermore,

given the changing mass of a resonance when compared to the changing temperature of the system, as will be discussed in Section 3.4.1, it is possible for off-shell resonances to decay at a range of different masses, thus making the observation of a complete mass shift to be impossible.

Even if it is not possible to observe a complete mass shift of a resonance signal, there are still other observables that can be used to determine if chiral symmetry has been restored. The two main observables for this discussion are width broadening and changes in yield ratios.

3.2.1 Width

The change in the mass of a particle will be hard to notice given that the QGP will hadronize after about $\sim 10 \text{ fm}/c$, thus leaving very little time for a particle to form and decay. Thankfully, resonances have the very useful property of having lifetimes on the same scale of the QGP lifetime. Furthermore, instead of the Dirac-like decay reconstruction of a more stable particle, which is only limited by the resolution of the detectors used to reconstruct them, resonances have a signal that can be characterized by the Breit-Wigner formula of Eq. 1. This is very useful because of the Γ , FWHM, that is used to determine the width of the signal.

While the mass of the resonance that is reconstructed is important for the Breit-Wigner formula, it is the Γ and its relation to the lifetime of the resonance that will prove to be most useful for our purposes. Similar to a Gaussian mean and σ , the mean and Γ of the Breit-Wigner formula contain most of the important information for this analysis.

If there is a signal of chiral symmetry restoration, it will be more noticeable in a change of the width of the resonance signal than a change in the mass of the signal. This change is explained in more detail in Section 3.4.1.

3.2.2 Yield

While the mass and width of a resonance that has its chiral symmetry restored is the most noticeable feature of chiral symmetry restoration, another important feature is the yield of the particle in question. If the mass of a particle decreases to a point that it drops below the kinematic limit of

a specific decay channel, i.e. the rest mass of the two decay particles, then the particle will not decay into the decay channel of interest. Thus the yield of the resonance, for the decay channel of interest, will decrease.

This, however, presents a new problem. Yields cannot be compared between different systems and energies due to the different multiplicities of each system. Yield is normally only compared between different particles in the same system at the same energy. Given that a comparison would need to be made between a system without QGP and one with QGP, such as p-p $\sqrt{s} = 13$ TeV and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV, to see if there is indeed a change in the yield, it would seem that using the yield to determine chiral symmetry restoration is simply not possible.

Fortunately, comparing yield ratios between different systems or energies would allow for the different multiplicities to be ignored. Simply by taking the ratio of the yields of two particles, one whose yield may change when chiral symmetry is restored and one whose yield will not change, one can determine if the yield has changed when chiral symmetry is restored.

The yield ratio would best be constructed between the chiral partners shown in Figure 1.8, since they contain the same spin and strangeness content, thus limiting the amount of potential other effects that could affect the yields. If the yield ratio is constructed as such, in the event that chiral symmetry remains broken, the ratio will remain constant. In the event that chiral symmetry is restored and the yield of one particle decreases, the ratio will increase or decrease depending on the ratio.

Similarly, the yield can also be constructed between the resonance particle and the stable ground state particle. Instead of relying on the property of positive parity particles to have a stable mass, the ground state is stable, and as such will have a much longer lifetime. Thus, the yield of the ground state will not be affected by chiral symmetry breaking. So, just like the ratio between the chiral partners, if chiral symmetry is restored and the yield of one particle decreases, the ratio will increase or decrease depending on how the ratio is defined. Thankfully, ALICE has already conducted measurements of the ground-state Ξ and $\Xi(1530)$ in the collision systems used in this analysis, so there is no need to search for these particles while also searching for the $\Xi(1820)$.

3.3 $\Xi(1820)$

The $\Xi(1820)$ is a resonance state that still has several features that have not been accurately measured over the years. However the positive parity partner of the $\Xi(1820)$, the $\Xi(1530)$, has been investigated in enough detail to determine the different masses and width for the charged and neutral state [12]. As seen in Table 2, the mass of the $\Xi(1820)$ is about $1823 \text{ MeV}/c^2$ with a full-width half-max of $24 \text{ MeV}/c^2$. The branching ratios (BR) for the $\Xi(1820)$ have not been well studied, thus the ratios only contain one measurement with large errors.

Parameter	Value
Charged States	$2 (\Xi(1820)^0 \text{ \& } \Xi(1820)^\mp)$
Quark Content	$\text{uss}(\Xi(1820)^0) \text{ \& } \text{dss}(\Xi(1820)^\mp)$
$I(J^P)$	$\frac{1}{2}(\frac{3}{2}^-)$
Mass	$1823 \pm 5 \text{ MeV}/c^2 \text{ (est.)}$ $1823.4 \pm 1.4 \text{ MeV}/c^2 \text{ (ave.)}$
Width	$24 \pm 15(10) \text{ MeV}/c^2 \text{ (est.)}$ $24 \pm 6 \text{ MeV}/c^2 \text{ (ave.)}$
Decay Mode (BR)	$\Lambda \bar{K}$ (large) $\Sigma \bar{K}$ (small) $\Xi \pi$ (small) $\Xi(1530)\pi$ (small)

Table 2: Quark content, quantum numbers, mass, width, decay modes, and branching ratios of the $\Xi(1820)$ [12]

Do note that the values for mean and width listed in Table 2 were calculated using two different methods in the PDG [12]. The estimate (est.) was based on the observed range of the data, not from a formal statistical procedure [12]. The average (ave.) was based on a weighted average of selected data [12]. Thus the average values are used for this analysis.

It should be mentioned that, while the $\Xi(1820)^0$ and $\Xi(1820)^\mp$ are treated in this case as having the same properties, that is most likely not the case. Similar resonances such as the $\Xi(1530)$ do have different masses and widths for the charged and neutral state. The $\Xi(1820)$ most likely does have a similar feature, but this requires a detailed study of both charged states. Regardless, for this analysis, both charged states are treated as having the same properties unless specifically

mentioned.

Furthermore, the branching ratios used for yield calculations for this analysis into the $\Xi(1820) \rightarrow \Lambda \bar{K}$ will be assumed to be 100 %. The only two measurements into the branching ratio of the $\Lambda \bar{K}$ listed in the PDG have been averaged to a value of 0.25 ± 0.05 , but these values are very poorly determined [12].

3.4 Estimation of Chiral Symmetry Restoration Effect Based on Calculations

3.4.1 Width Broadening

As discussed in Section 1.3, the signal normally produced from reconstructing resonance particles from their daughter particles contains several important characteristics of the resonance particle. The mean is related to the mass of the particle. The FWHM is related to the lifetime of the particle as seen in Eq. 2.

However, with the possibility of observing chiral symmetry restoration taken into account, the normally reliable Breit-Wigner function develops a small problem. Namely, the function can not account for the potential effects of chiral symmetry restoration, at least not in its idealized form.

Similar to how a double Gaussian can be approximated to a single Gaussian, the potential “double” (or even “multiple”) Breit-Wigner functions can still be approximated by one Breit-Wigner function. Though, like any approximation, there will be an inherent error with such an application. In the case of a double Gaussian being approximated by a single Gaussian, there would be a noticeable effect on the mean and sigma of the single Gaussian. The significance of the effect will depend on the parameters of the double Gaussians such as their means, their sigmas, their amplitudes, and their distance away from each other. Similarly, using one Breit-Wigner to approximate multiple potential signals would show similar effects.

In the idealized case of the $\Xi(1820)$ signal in a vacuum without chiral symmetry restoration as in the PDG, the signal would appear like the red line shown in Figure 3.1 [12]. To approximate the potential effects of chiral symmetry restoration on the signal, a few key points need to be addressed.

First, as will be mentioned in the following sections, the $\Xi(1820)$ signal is reconstructed from

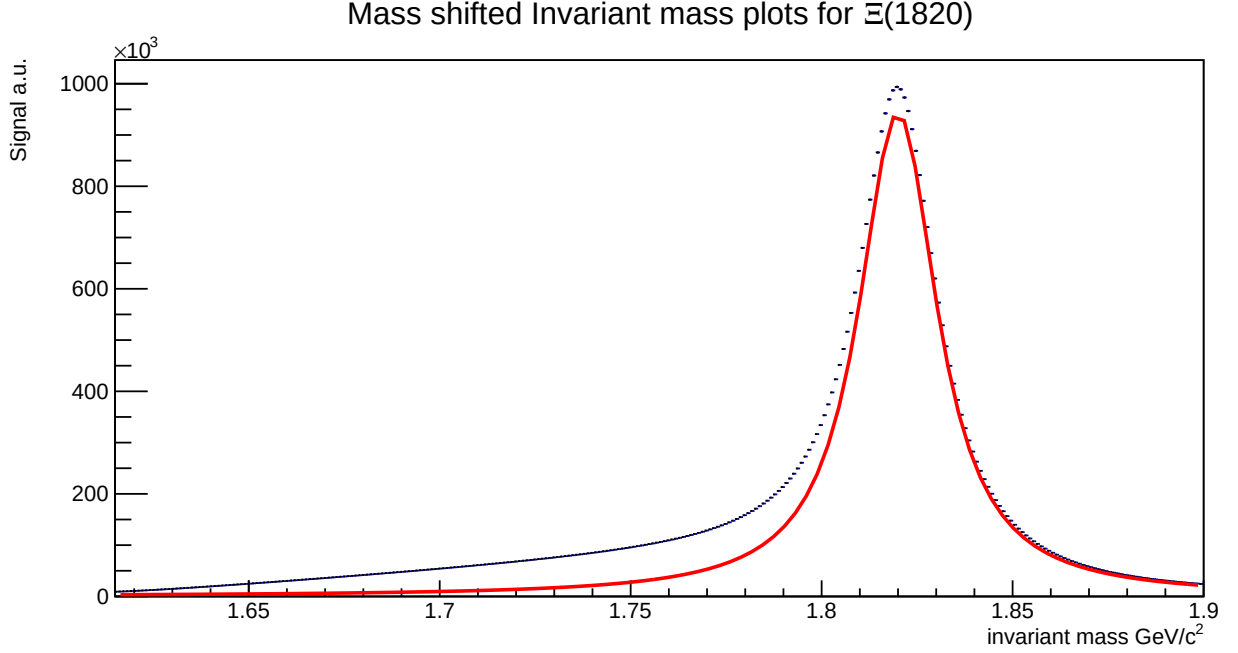


Figure 3.1: Theoretical distribution of $\Xi(1820)^\mp$ masses. Red: Distribution of $\Xi(1820)^\mp$ for a selection of particles that were not subject to possible in-medium modification. Black data points: Distribution of $\Xi(1820)^\mp$ that contains contributions from in-medium modification effects.

its decays to Λ and K ($K^\pm$ in this case). Thus, this limits the potential “decrease” of the $\Xi(1820)$ mass, in that it can not have a mass lower than the sum of the masses of the particles it will decay into. With the mass of the Λ at $1115 \text{ MeV}/c^2$ and the $K^\pm$ at $494 \text{ MeV}/c^2$, the “minimum” mass that the $\Xi(1820)^{(\mp)}$ could drop to would be about $1609 \text{ MeV}/c^2$ [12]. While the $\Xi(1820)$ could drop to an even lower mass, it would not decay into the decay channel we are investigating, so it is ignored for the purpose of this discussion.

Second, due to the de-confined nature of the QGP, it is not possible to form a $\Xi(1820)$ before hadronization has begun. At first, this may seem counter-intuitive. For a $\Xi(1820)$ to have a lower mass, it needs to form in the QGP. However, in the QGP a $\Xi(1820)$ can not form. The solution to this problem comes from evidence that the phase transition from the QGP to hadronic matter is not an instantaneous first-order phase transition, but instead a crossover transition [44]. Lattice QCD calculations such as the one shown in Figure 3.2 show the difference between light

and strange chiral condensates. The chiral condensate can be considered an order parameter of chiral symmetry. If $\langle \bar{q}q \rangle$ is 1, then chiral symmetry is broken, while if $\langle \bar{q}q \rangle$ is 0, then chiral symmetry is restored. Furthermore, because of the fact that strange quarks have a much higher mass than light quarks ($m_s = 96 \text{ MeV}/c^2 > m_{u,d} \sim 10 \text{ MeV}/c^2$) the difference between the quark masses and chiral condensates will decrease as chiral symmetry is restored [12]. Figure 3.2 seems to indicate that the transition is more of a mixed-phase transition; thus it is possible (though rare) for a particle to form at a temperature hot enough for chiral symmetry to be partially broken. In this tight window of mixed-phase transition, lies the possibility for a $\Xi(1820)$ to form.

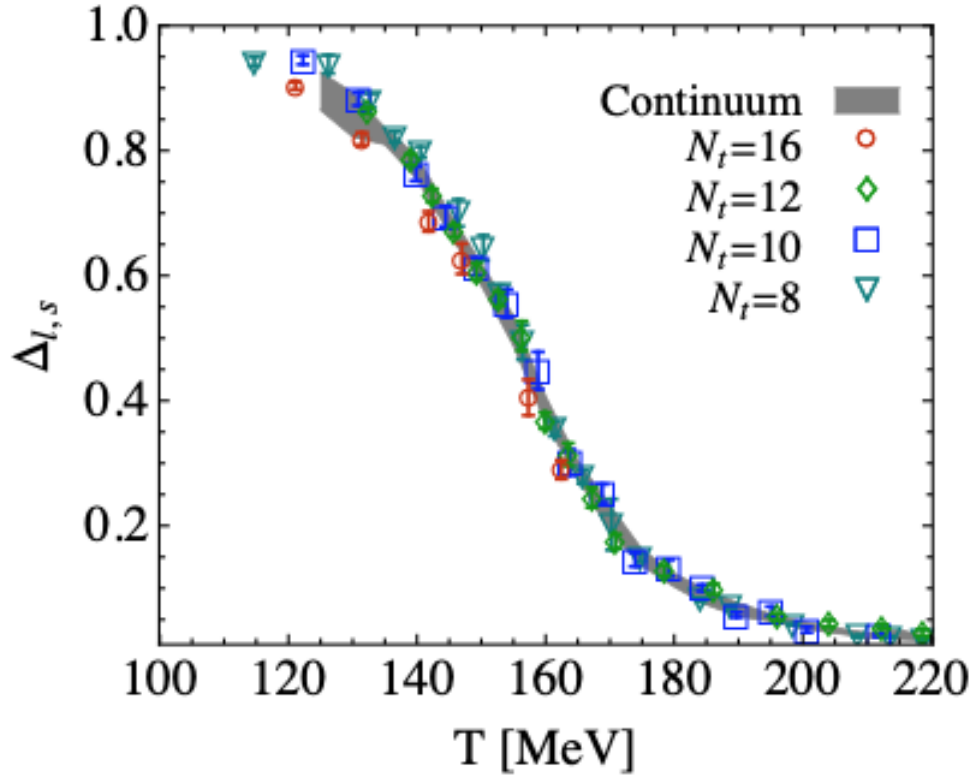


Figure 3.2: Difference in chiral condensate between light and strange quarks ($\Delta_{l,s}$) as a function of the temperature [44].

Third, the timeframe of the sequence of events needed for the signature of chiral symmetry restoration to be seen is very tight. A $\Xi(1820)$ needs to form in the phase transition from QGP

(chiral symmetry restored) to hadronic matter (chiral symmetry broken), decay in the time frame such that its mass is between the mass of the vacuum mass of $\Xi(1820)$ and the minimum mass of the $\Lambda + K^\pm$, and be detectable among the other $\Xi(1820)$ resonances that decay at normal mass. This is further compounded by the low yield of the $\Xi(1820)$.

With all of the previous key points considered, it is still possible to approximate the potential effects of chiral symmetry restoration on the $\Xi(1820)$ signal. This can be done using a number of approximations.

First, an approximation to the potential temperature ranges in which the $\Xi(1820)$ could decay, needs to be taken into account. In Figure 1.8, the potential temperature ranges can be interpreted by determining the temperature at which the $\Xi(1820)$ mass reaches its minimum mass of about $1610 \text{ MeV}/c^2$. Because this is a lattice QCD calculation with overestimated masses of the particles in question, temperature ratios with higher than expected critical temperature, and noticeable uncertainties, this estimate is far from ideal. However, by normalizing the critical temperature in the simulation to the critical temperature found in experiments, and readjusting the mass of the $\Xi(1820)$ to its correct values, an approximate value can be inferred. Furthermore, an approximate function of the mass of the $\Xi(1820)$ vs. temperature can be inferred from the distribution.

Second, an approximation of the probability for a $\Xi(1820)$ forming must be formulated. Figure 3.2 can be interpreted as a function of the probability for a particle to form at a certain temperature. While this is not the primary use of such a figure, it is useful for our purposes. With this, an equation can be derived for the probability of a particle forming during the phase transition vs. temperature.

Third, an approximation of the time evolution of the temperature in the QGP as it undergoes a phase transition is necessary. While this has been done using figures such as Figure 1.2 or Figure 3.3, the main aspect of the function is to determine at what temperature the system is, as a function of time. Using the approximations of $T = 180 \text{ MeV}$ at $5 \text{ fm}/c$, $T = 155 \text{ MeV}$ at $10 \text{ fm}/c$, and $T = 130 \text{ MeV}$ at $15 \text{ fm}/c$, a simple linear relation can be produced. Note that the reason that temperature vs. time is needed is to determine the time necessary for the $\Xi(1820)$ to decay after forming.

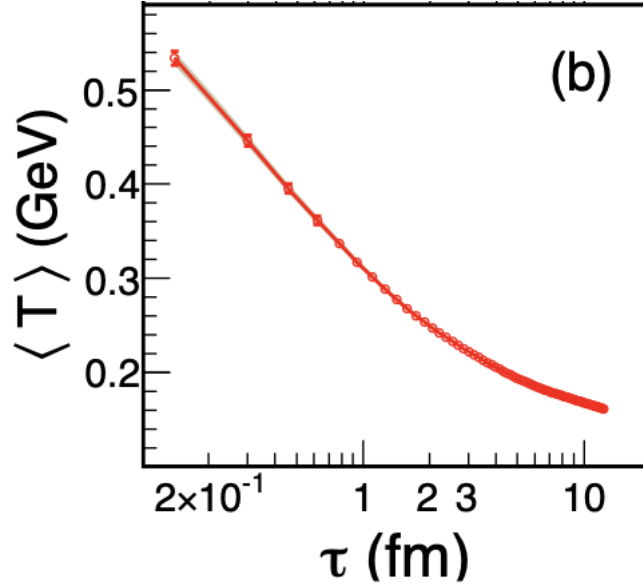


Figure 3.3: Temporal evolution of average temperature. Obtained from hydro-dynamic calculations for central Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. [45]

Finally, by convoluting the probability of a $\Xi(1820)$ forming with the probability of the $\Xi(1820)$ decaying in the mass range that is of interest to us, we can approximate the potential effects of chiral symmetry restoration. Figure 3.4 shows the potential signal we would expect from a $\Xi(1820)$ (blue) that also includes resonances formed and decayed before chirality is completely broken. Notice the elevated section of the distribution below the $\Xi(1820)$ mass peak is significantly higher than the signal that does not take into account such particles. This makes the normally Breit-Wigner shaped signal appear stretched to the left. Furthermore, Figure 3.4 shows that the normal Breit-Wigner fit (red) would instead be approximated by the slightly wider Breit-Wigner fit (blue) to the left. Thus, if $\Xi(1820)$ data did include a signature of chiral symmetry restoration, it would be seen in a slight shift of the peak of the signal and an increase in the width. The increase in the width is referred to as “width broadening”, and is the major focus of this dissertation.

Do note that Figure 3.4 is simply the Voigtian, a convolution of a Breit-Wigner with Gaussian function discussed in Section 4.3, fits of the data from Figure 3.1. The Voigtian fit of the $\Xi(1820)$ distribution from Figure 3.1 with no in-medium modification was constructed from a Voigtian

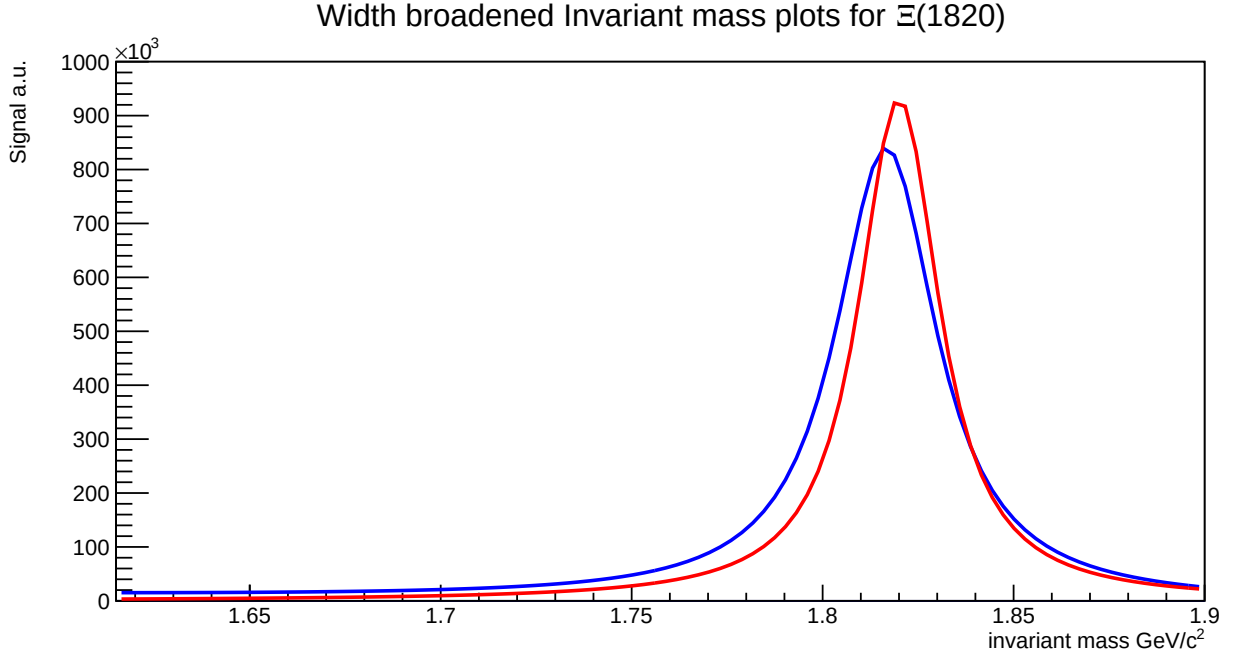


Figure 3.4: Voigtian fits (convolution of Breit-Wigner with Gaussian function, discussed in Section 4.3) of distributions seen in Figure 3.1. Red: Fit of Voigtian function to $\Xi(1820)$ signal that has had no in-medium modification (Figure 3.1 red distribution). Blue: Fit of Voigtian to $\Xi(1820)$ signal that includes in-medium modification (Figure 3.1 black data points). Note that the mean of the two distributions has shifted by approximately 3 MeV/ c^2 (Red = 1820.00 MeV/ c^2 , Black = 1816.64 MeV/ c^2) and width has increased between the two distributions by about 7 MeV/ c^2 (Red = 24.0000 MeV/ c^2 , Blue = 30.8952 MeV/ c^2). σ fixed to 2 MeV/ c^2 for this analysis, see Section 4.3.

distribution, so the Voigtian will fit perfectly to the distribution as seen in Figure 3.4. The Voigtian fit of the $\Xi(1820)$ distribution from Figure 3.1 that contains in-medium modification effects will not have a perfect Voigtian distribution, so the fit will be different than the previous distribution as seen in Figure 3.4.

While these calculations and plots are highly informative and useful in estimating the potential effects of chiral symmetry restoration, please do note that the previous discussion and figures were all produced using multiple assumptions, approximations, and applications that are not rigorous enough to be conclusive. Assumptions such as the lifetime of the $\Xi(1820)$ remaining constant or the probability of the formation of $\Xi(1820)$ forming for example, greatly limit any true predictive

power of these statements. With all of that in mind, it is still informative to have any calculation on such a rare and difficult to detect signature, even if the idealized form may not be present in reality. These calculations allow us to gauge the expectations and the necessary statistics in order to obtain a measurable signal.

3.4.2 Yield Ratios

Another important signature of chiral symmetry restoration would be the potential effects on the yield of the $\Xi(1820)$. More specifically, the yield ratio between the $\Xi(1820)$ and $\Xi(1530)$. As seen in Figure 1.8, while the mass of the $\Xi(1820)$ would decrease as it approaches the critical temperature, the mass of the $\Xi(1530)$ would remain relatively constant. This would mean that, in the event that the $\Xi(1820)$ mass decreased to a point that the total measurable yield of $\Xi(1820)$ decreased while the yield of the $\Xi(1530)$ does not, this is an indicator of a signature of chiral symmetry restoration. Such an effect could be propagated through 2 similar possibilities.

First, the $\Xi(1820)$ signal could decrease so much that the Breit-Wigner fit used to fit the signal can not catch all of the potential signal outside the mass region used to determine the yields. Similar to Figure 3.4, the potential extra signal from the lower mass $\Xi(1820)$ particles might not be integrated into the fit and yield calculation.

Second, the $\Xi(1820)$ signal could drop below the minimum mass threshold of the Λ and $K^\pm$ masses, thus it would not decay into the decay channel investigated. Though this would be rare, accounting for the forming and decaying below the minimum mass threshold would be possible. If such a $\Xi(1820)$ did form, it could decay below the minimum mass threshold, most likely into some decay channel such as $\Xi(1820) \rightarrow \Xi + \pi$ or some similar decay channel.

Regardless, if a yield ratio between the $\Xi(1820)$ and the $\Xi(1530)$ can be formed, this would also prove useful for confirming another signature of chiral symmetry restoration. The disadvantage is the need for sufficient statistics to construct a p_T spectrum. The p_T spectrum is needed because of the fact that the yield is only measured in the acceptance given. To calculate the total yield, the yields must be extrapolated using the p_T spectrum to account for p_T -ranges that are not measured.

Furthermore, at low p_T some of the yield is lost because the magnetic field causes low momentum charged particles to not be detected. At high p_T , some of the yield is lost because of the lack of resolution since tracks do not curve enough in the magnetic field and all momenta become indistinguishable. In dealing with a particle of such low yield already, the required statistics for such an analysis could prove insufficient for the amount of data currently available. Nevertheless we will show our current result on the yield ratios in Section 7.4.

4 Analysis Method

4.1 Reconstruction

In order to study chiral symmetry restoration and the phase transition of the QGP to hadronic matter, the $\Xi(1820)$ resonance was measured in p-p collisions at $\sqrt{s} = 13$ TeV, p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV by reconstruction of its hadronic decay into ΛK using data collected at ALICE. In this chapter, detailed steps used in the $\Xi(1820)$ reconstruction will be provided.

4.1.1 Event Selection

For an event containing a $\Xi(1820)$ to be analyzed, it must pass the following criteria:

- Events must have a z-position of the primary vertex (V_z) within ± 10 cm of the center of the ITS/TPC (center of the detector). This is to make sure that the collision did occur in the center of the detector, and not outside the expected region of collision. This also ensures that the tracks produced in the collision lie within the acceptance of the central barrel detectors, even for larger pseudorapidities ($|\eta|$ near 0.8).
- Reject pile-up events. These are events that contain particles or tracks from other events and as such can “contaminate” one event with the data from other events. One of the main ways to reject such events is to require that an event must contain a minimum bias signal from the V0 detectors to be considered for the analysis.

4.1.2 Signal Reconstruction

The $\Xi(1820)$ signal was reconstructed using invariant-mass analysis of decay products such as Λ and K particles. Please note that the use of the symbol $\Xi(1820)^\mp$ is a bit unique in this discussion. There is technically no $\Xi(1820)^+$ particle, only the $\Xi(1820)^-$ particle and the $\bar{\Xi}(1820)^+$ anti-particle (which is also assumed) case. However, to avoid confusion with how to represent the charged $\Xi(1820)$ for this discussion, $\Xi(1820)^\mp$ is used so that the (-) can represent the “normal” negative

charged state, and the (+) can represent the “anti-particle” positive charged state. The charged and neutral states of the $\Xi(1820)$ follow a simple decay process as shown here:

$$\Xi(1820)^{\mp} \rightarrow \Lambda(\bar{\Lambda})K^{\mp} \quad (27)$$

$$\Xi(1820)^0 \rightarrow \Lambda(\bar{\Lambda})K_s^0 \quad (28)$$

One should note that, due to the neutral charge of the Λ and K_s^0 , they will not be directly detected by the TPC and TOF. Furthermore, given the time dilation and average decay lengths of the Λ and K_s^0 (7.89 and 2.68 cm, respectively), these particles will need to decay before they can be detected by the TPC and TOF, though some may be able to decay inside the TPC. While some Λ and K_s^0 particles will decay at or before their corresponding decay length, about 36.79 % will decay later using the exponential decay law. These particles decay in the decay process as shown here:

$$\Lambda(\bar{\Lambda}) \rightarrow p(\bar{p})\pi^{\mp} \quad (29)$$

$$K_s^0 \rightarrow \pi^{\pm}\pi^{\mp} \quad (30)$$

Thus the complete decay process of both charged and neutral states of the $\Xi(1820)$ would appear as:

$$\Xi(1820)^{\mp} \rightarrow p(\bar{p})\pi^{\mp}K^{\mp} \quad (31)$$

$$\Xi(1820)^0 \rightarrow p(\bar{p})\pi^{\mp}\pi^{\pm}\pi^{\mp} \quad (32)$$

Due to the fact that the neutral channel of the $\Xi(1820)$ requires a 4-body reconstruction, as opposed to the charged $\Xi(1820)$ signal’s 3-body reconstruction, and the efficiency to reconstruct a K_s^0 is lower than for $K^{\pm}$, thus leading to lower efficiency of the $\Xi(1820)^0$ when compared to the $\Xi(1820)^{\mp}$, the neutral channel is much more difficult to analyze. While still noticeable in some sections of the p-p $\sqrt{s} = 13$ TeV data as seen in Figure 4.1 for certain p_T and multiplicity distributions, the signal is much smaller in comparison to the charged channel’s signal in the same

distribution. It is for this reason that, while some information about the neutral $\Xi(1820)$ can be interpreted, it is not the main focus in this thesis.

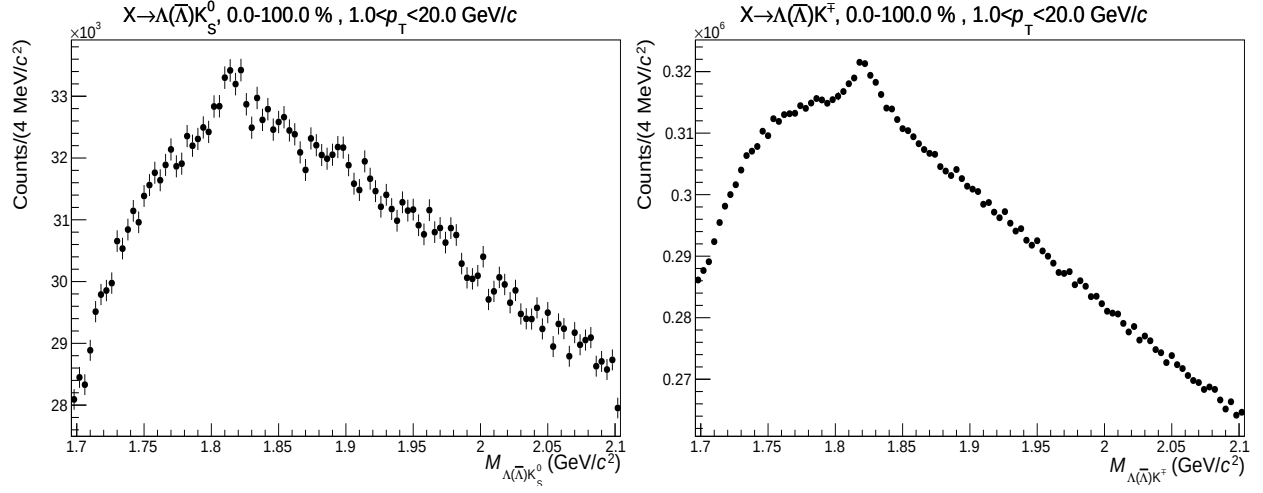


Figure 4.1: Left: Invariant mass plot for $\Lambda(\bar{\Lambda})K_s^0$. Right: Invariant mass plot for $\Lambda(\bar{\Lambda})K^\pm$.

4.1.3 Selection of Primary $K^\pm$

The cuts used to identify the decay products of the $\Xi(1820)$ have been selected so as to line up with other analyses of similar resonances. These primarily consist of determining the properties of the decay products and the potential signal that can be identified.

Due to fact that $K^\pm$ can be detected directly from the TOF and TPC, the procedure to identify and use $K^\pm$ for reconstruction of the $\Xi(1820)$ is much simpler than the process to identify the Λ or K_s^0 . The candidates for $K^\pm$ must satisfy the standard set of cuts used for many ALICE studies of resonances with the following conditions as seen in Table 3 [33].

In order to make sure that the $K^\pm$ selected can match with a Λ , an autocorrelation check is performed. This autocorrelation check makes sure that the $K^\pm$ selected is not on the same track as the $\pi^\pm$ or proton that comes from the Λ decay. Do note that the $K^\pm$ would be a misidentified π or proton from a Λ decay, which would then by chance be paired with another track to create a fake “ Λ ”. Furthermore, tracks from Λ decays can still be considered in general, just not when they

Track cuts	
p_T	$> 0.15 \text{ GeV}/c$
$ \eta_{lab} $	< 0.8
Number of hits in the inner two layers of the ITS for $K^\pm$	≥ 1
Number of crossed rows in TPC	≥ 70
Ratio of crossed rows in TPC over findable clusters	≥ 0.8
DCA of $K^\pm$ to Primary Vertex (PV) in transverse direction	$< 7 \sigma_{xy}$
DCA of $K^\pm$ to PV in longitudinal direction	$< 2 \text{ cm}$
PID of $K^\pm$ using TPC	$< 2 \sigma_{TPC}$
PID of $K^\pm$ using TOF	$< 3 \sigma_{TOF}$

Table 3: Primary track cuts

are paired with the Λ that produced them.

Tracks must have a $p_T > 0.15 \text{ GeV}/c$ and $|\eta_{lab}| < 0.8$ [46]. The magnetic field used for ALICE causes charged particles with transverse momentum lower than $150 \text{ MeV}/c$ to spiral and not even reach the detectors [28]. Primary tracks for the $K^\pm$ candidates require at least one hit in the inner two layers of the ITS, but this is not true for the V^0 daughter tracks as discussed in Section 4.1.4.

Track cuts such as the number of crossed rows in the TPC and the ratio of crossed rows in the TPC over findable clusters ensure that tracks are well enough defined. Due to the limited efficiency of the TPC, not all tracks will give a signal even if it passes through a row. By requiring that the number of crossed rows in the TPC be greater than or equal to 70, out of the 159 possible rows, and higher ratios, this leads to more certainty and available data to determine the properties of the tracks [28].

Track cuts such as the distance of closest approach (DCA) of $K^\pm$ to the longitudinal direction (z axis) to be less than 2 cm and the DCA of $K^\pm$ to the transverse direction (xy axis) to be less than $7 \sigma_{xy}$ ensure that the primary tracks are close to the primary vertex of the collision [46]. The resolution of the DCA in the transverse plane, σ_{xy} , is very tight, lower than $100 \mu\text{m}$ for $p_T > 0.5 \text{ GeV}/c$, but also strongly p_T -dependent [46].

PID cuts for the $K^\pm$ in the TPC and TOF of 2 and 3 σ respectively are part of the standard set of cuts used for identifying $K^\pm$. While the TPC and TOF n_σ cuts are not part of the standard set of cuts used by other ALICE studies, these method are still commonly used. Normally, given

the resolution of the TPC on dE/dx of about 6%, the TPC would allow for a $2 \sigma_{TPC}$ separation between $K^\pm$ and other particles up to $p_T \sim 0.8 \text{ GeV}/c$ in the low momentum region, and above $3 \text{ GeV}/c$ in the relativistic rise region of the dE/dx curve [46]. The n_σ is increased slightly for low momentum from this value in order to have a higher number of $K^\pm$ identified for this analysis. Similarly, given the resolution of the TOF and the T0 detector that is used to measure the start time of the event, the TOF is able to show a separation of $2 \sigma_{TOF}$ between π and K in the momentum range of $0.7\text{-}3 \text{ GeV}/c$, with a similar level of separation between K and p up to $5 \text{ GeV}/c$ [46].

4.1.4 V^0

V^0 , not to be confused with the V0 detector or the primary vertex, are neutral, unstable subatomic particles (mesons and baryons) which can decay into more stable lighter charged hadrons which curve in different directions in a magnetic field and therefore form the shape of a ‘V’ with their decay trajectories. Particles such as the Λ and the K_s^0 are V^0 particles since they have short lifetimes (not as short as resonance lifetimes) and are neutral. While resonances decay strongly, leading to their short lifetimes of 10^{-23} seconds and conservation of quantum numbers like strangeness, V^0 particles decay weakly, thus leading to their longer lifetimes and the lack of quark flavor conservation. Due to their neutrality, they would normally never be detected directly using the ALICE detectors. However, it is possible to detect their daughter particles ($\Lambda(\bar{\Lambda}) \rightarrow p(\bar{p})\pi^\mp$ and $K_s^0 \rightarrow \pi^\pm\pi^\mp$), which are charged and are stable enough to hit the detectors.

As described in Section 4.1.2, the Λ and K_s^0 must be reconstructed from their daughter particles that are detectable using the TPC. While this is slightly similar to the process discussed in Table 3 for the $K^\pm$, because this is a reconstruction of the V^0 mother particles, the process is much more complicated. This is due to the added criteria that the identified p and $\pi^\pm$ must come from the V^0 mother particles, and not from the primary vertex where most other particles are produced.

Since pions from the weak decay of the K_s^0 ($c\tau = 2.68 \text{ cm}$) and pions and protons from the weak decay of the Λ ($c\tau = 7.89 \text{ cm}$) are being produced away from the PV, separate topological cuts and track selection criteria are applied as seen in Table 4 [12].

Topological cuts	
DCA of V^0 decay products to PV in XY axis	> 0.06 cm
$\cos\theta_\Lambda$	> 0.99
$r(\Lambda)$	$0.5 < r(\Lambda) < 200$ cm
DCA between Λ decay products	< 1.0 cm
DCA of Λ to PV	> 0.4 cm
PID of π (decayed from Λ)	$< 5 \sigma_{TPC}$
PID of proton (decayed from Λ)	$< 5 \sigma_{TPC}$
$ M_{p\pi} - m_\Lambda $	< 6 MeV/ c^2
$\cos\theta_{K_s^0}$	> 0.97
$r(K_s^0)$	$0.5 < r(K_s^0) < 200$ cm
DCA between K_s^0 decay products	< 1.0 cm
DCA of K_s^0 to PV	> 0.3 cm
PID of π (decayed from K_s^0)	$< 5 \sigma_{TPC}$
$ M_{\pi\pi} - m_{K_s^0} $	< 30 MeV/ c^2

Table 4: Topological and track selection criteria

While most of these cuts are the standard cuts used for charged kaon and other particles, they are still fairly complex in their use. Thus, I will discuss each cut and selection used and explain the reasoning for each. Some further details about which cuts are changed to determine the systematic uncertainty will be discussed in Section 6.2.

Track selection cuts such as the DCA of V^0 decay products to the PV in xy axis to be greater than 0.06 cm ensure that the decay products' tracks are not too close to the primary vertex of the collision [46]. If they were too close, then it may be possible that the decay particle candidate came from the primary vertex instead of the V^0 .

Topological cuts such as the n_σ cuts from the TPC on π and p have been increased from their normal values so as to allow for more statistics in their reconstruction of their mother particles. As discussed in Section 4.1.3, given the resolution of the TPC, the TPC would allow for a $2 \sigma_{TPC}$ separation between π , p , and other particles up to $p_T \sim 0.8$ GeV/ c in the low momentum region, and above 3 GeV/ c in the relativistic rise region of the dE/dx [46]. Similarly, the TOF would allow for a $2 \sigma_{TOF}$ separation from π and p in the range of 0.7-3 GeV/ c and up to 5 GeV/ c respectively [46]. This was increased to 5σ so that more π and p from the decays of Λ and K_s^0 can be identified.

The topological cuts of $\cos\theta$ will be discussed in more detail in Section 6.2. However, the basic

definition of this cut is that it refers to the angle between the reconstructed particle's direction of momentum and the vector connecting the primary vertex to the V^0 decay point, simply known as the pointing angle. This pointing angle is illustrated in Figure 4.2. The smaller the angle, the more likely it is that the particle is coming from the primary vertex. Instead of using a value of the angle, the cosine of the angle is calculated. Do note that this cut is used on all reconstructed particles.

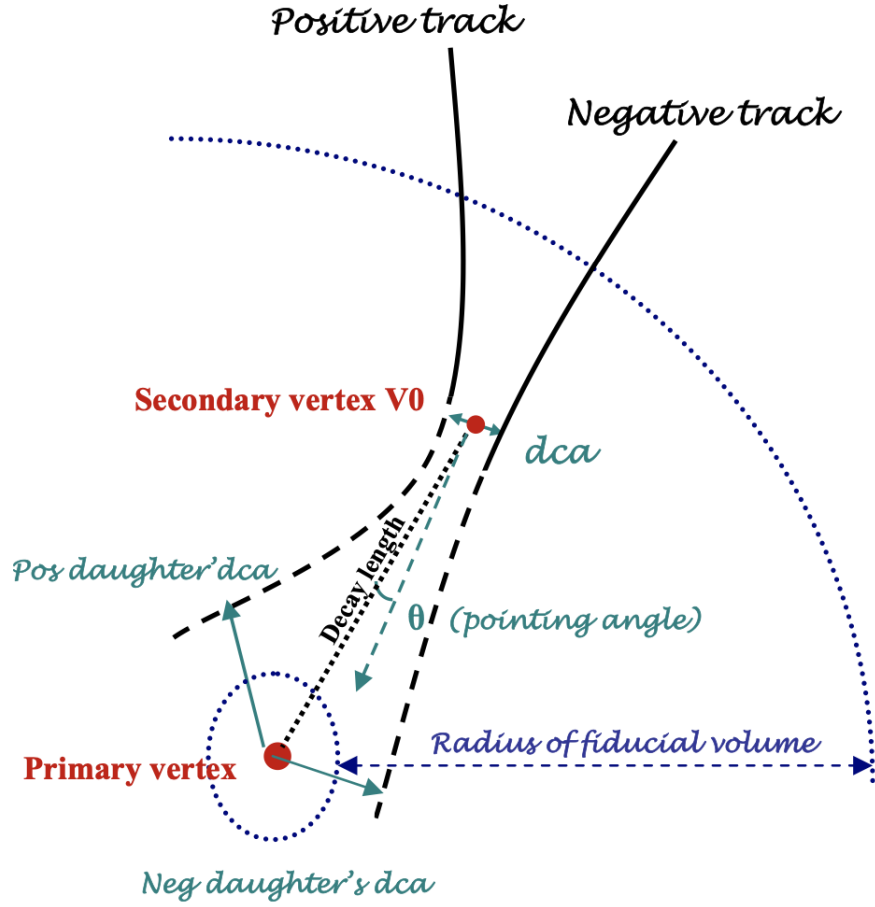


Figure 4.2: Figure of a V^0 decay [47].

The topological cuts of the DCA of the decay particles between each other is very important in determining if two particles can be used for the V^0 analysis. This DCA is also illustrated in Figure 4.2. After all secondary tracks of opposite charges have been examined, the DCA between each pair

of tracks is calculated to determine if the two tracks may have come from the same reconstructed mother particle. So, the closer two tracks of opposite charges are to each other, the higher the likelihood that they came from the same decay. However, given the resolution of the detectors used to identify the particles, and the calculations needed to reconstruct the helical path of the particles, it will not be possible to trace two particles back to the point where they decayed from their mother particle with infinite precision. Thus, the distance of closest approach (DCA) is used to determine if two particles are close enough to each other.

The topological cut of r refers to the decay radius. As stated in Section 2.1, charged tracks move in a helical motion that is determined by their kinematics and the electric and magnetic field they travel in. Furthermore, as discussed with the DCA of the decay particles, the particles will not be able to be matched “perfectly” to each other with infinite precision. This does lead to the question of where did the V^0 actually decay in 3D space. The simplest assumption would be that the V^0 is at the midpoint between where the daughter particles appear closest (DCA) to each other. However, because the V^0 can not be physically detected, a line connecting the positions of closest approach between the two daughter particles is estimated using a minimization program, with the V^0 being one point of this line, as shown in Figure 4.2 in 2D representation with the label “Secondary vertex V^0 ”. Because the distance from each daughter track is taken to be proportional to the precision of the track parameter estimations, once their positions are determined, only the V^0 candidates located inside the fiducial volume are kept [48]. The decay radius is also determined by the decay length of the V^0 particles. Because the decay of particles is still probabilistic, it is entirely possible for particles to decay before their associated decay length. While similar to the concept of the half-life, the decay length corresponds to the length a decaying particle will travel, moving at the speed of light, before 63.21 % ($1-e^{-1}$) of the initial decaying particles decay. After traveling two decay lengths, 86.47 % of particles will have decayed. The longer the decay length of the particles, the longer the inner boundary of the fiducial volume from the primary vertex to the V^0 . The inner boundary of this fiducial volume, like a cylinder with its z-axis extending into the beam axis, is at a radius of 0.5 cm from the primary vertex, while the outer limit is set at 200 cm,

thus covering most of the TPC volume [48].

The DCA of the reconstructed particles to the primary vertex is the simplest cut among those listed here. Similar to DCA of the decay particles to each other and the $\cos\theta$, the DCA from the V^0 to the PV is determined by extrapolating the V^0 's momentum vector back towards the primary vertex. The DCA then measures how close that trajectory would get to the primary vertex.

The track cut selection of the masses, such as $|M_{p\pi} - m_\Lambda|$, is very instrumental to the verification that the reconstructed particle does in fact reconstruct to the particle of interest. Even with all of the previous cuts and criteria, if the simple fact is that two daughter candidates for a V^0 mother particle do not reconstruct to the proper mass, then it is given that the two candidate tracks did not come from a V^0 decay.

4.1.5 Invariant Mass Techniques

Invariant mass plots are one of the most important tools for particle physics when it comes to reconstruction of particles. When a particle decays, the information about its mass, momentum, and even quark content (depending on the decay strength) will be split across the daughter particles. It stands to reason that if the information of the daughter particles is collected, then the initial mother particles can be reconstructed. In fact, the derivation and equation to determine the mass of the mother particle from two daughter particles is

$$E^2 = m^2 + \vec{p}^2 \quad (33)$$

$$E^2 = (E_1 + E_2)^2 = E_1^2 + E_2^2 + 2E_1E_2 \quad (34)$$

$$M^2 + (\vec{p}_1 + \vec{p}_2)^2 = m_1^2 + \vec{p}_1^2 + m_2^2 + \vec{p}_2^2 + 2E_1E_2 \quad (35)$$

$$M^2 = m_1^2 + m_2^2 + 2(E_1E_2 - \vec{p}_1 \cdot \vec{p}_2), \quad (36)$$

where the 1 and 2 subscripts represent the two daughter particles, M is the mass of the mother particle, E is the total energy of each daughter particle, $\vec{p}$ is the vector momentum of each daughter

particle, and the speed of light c is set to natural units ($c = 1$).

However, given that this analysis is for a collider experiment so principally the only momentum calculated is the transverse momentum, p_T , and most of the particles with high p_T are so relativistic that all of the energy of the particle is in the momentum of the particles, the previous equation can be redefined so that

$$M^2 = 2p_{T1}p_{T2}(\cosh(\eta_1 - \eta_2) - \cos(\phi_1 - \phi_2)) \quad (37)$$

can be used, with η corresponding to the pseudo rapidity and ϕ corresponding to the azimuthal angle. This would not be the case for some particles with higher mass and lower momentum, such as the $K^\pm$ which can be measured down to $p_T \sim 150 \text{ MeV}/c$.

A simple invariant mass plot of the reconstruction of $\Lambda + K^\mp$ is shown in Figure 4.3. Do note that seen among the black data points, there does appear to be a bump around $1.820 \text{ GeV}/c^2$. This is because, while some of the Λ and $K^\mp$ particles that are used for reconstruction of the mother particle actually correspond to the $\Xi(1820)$, there are large amounts of reconstructed $\Lambda + K^\mp$ that do not reconstruct to the $\Xi(1820)$, but instead correspond to a background of unrelated particle pairs.

Due to the peak structure of the signal in Figure 4.3, it is much easier to notice any actual reconstructed particles as opposed to a background distribution.

Invariant mass plots are important to the discovery and analysis of reconstructed particles, especially those that have not been studied much, such as the $\Xi(1820)$.

4.1.6 Discussion of Combinatorial Backgrounds

As stated in the previous section, not all daughter particle candidates will be reconstructed to match a potential mother particle. In fact, as seen in Figure 4.3, a large portion of $\Lambda(\bar{\Lambda}) + K^\mp$ do not reconstruct to the $\Xi(1820)$, but instead to a smooth curve that is called the combinatorial background.

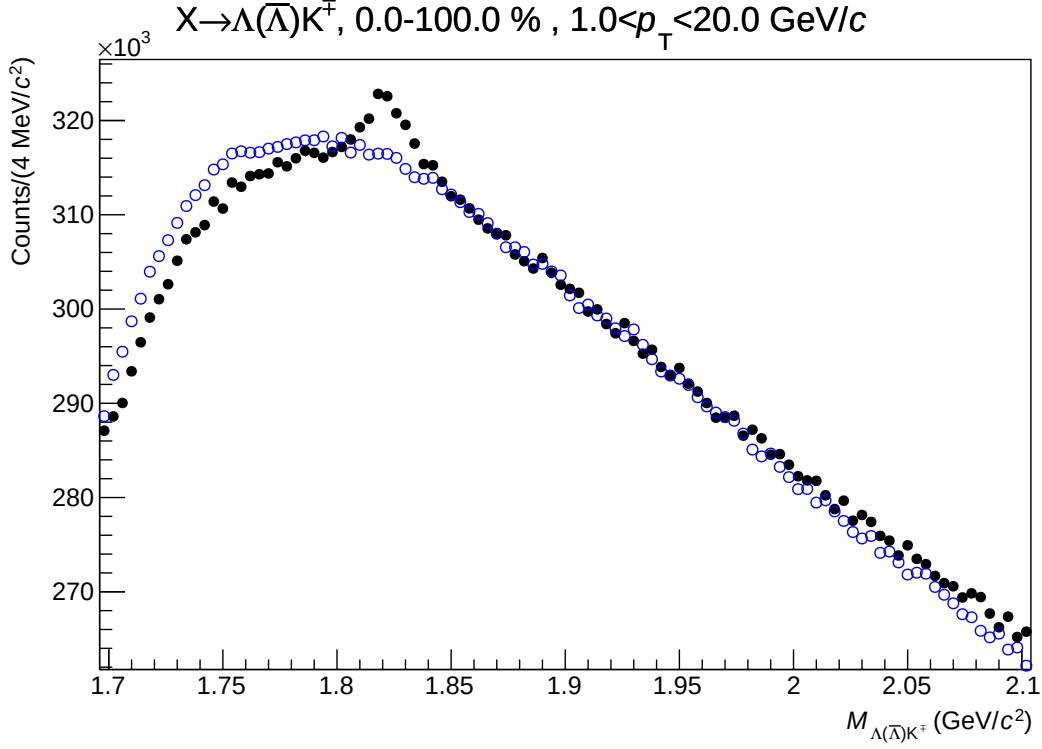


Figure 4.3: Invariant mass plot for $\Lambda(\bar{\Lambda})K^\mp$ for p-p collisions at $\sqrt{s} = 13$ TeV. Black data points represent raw signal. Blue data points represent estimated background, which will be discussed in more detail in Section 4.2.

The combinatorial background represents the amount of daughter particles that, while still matching most of the criteria used for reconstruction of the particle of interest, simply do not reconstruct to a true particle. Do not forget that the $\Xi(1820)$ decays very close to the primary vertex. So it becomes almost impossible to distinguish a Λ or K that came from a $\Xi(1820)$ decay, from all of the other Λ and K particles that are also produced in the initial collision. Furthermore, other resonances may decay into Λ or K particles, thus increasing the total number of Λ and K particles released into the sea of other particles that are being produced from the initial collision. Given all of these sources for Λ and K particles and the possibility of particles being misidentified, it is no wonder that a sizable background is created for $\Lambda+K$ invariant mass plots.

Despite this, a signal of the $\Xi(1820)$ can still be seen in Figure 4.3. Normally, the signal would simply be fit with a function used to represent the signal, such as the Breit-Wigner function.

However, given the very large combinatorial background and the small signal, it is simply not possible to represent the signal with a Breit-Wigner function. Furthermore, the combinatorial background would cause any calculation of yield to also contain a large amount of “non-signal” particles, those that correspond to the background and not the signal of interest.

Due to the significant amount of combinatorial background seen in the invariant-mass plots of the $\Xi(1820)$, it is difficult to extract a reliable signal to be used for analysis. To better estimate the background distribution, two methods to determine the background distribution were applied.

This is done so that an estimation of the “raw data’s” background can be calculated using similar candidates as used in the normal distribution. For example, to create a “mixed-event” background, one of the candidates for either a Λ or K particle is matched to the corresponding other particle from a different event in order to reconstruct the $\Xi(1820)$. This allows for a way to create the background of the distribution for $\Lambda+K^\mp$, without any of the signal from the $\Xi(1820)$ or sources of correlated background.

The mixed-event background was generated by combining uncorrelated decay products from 5 different events in p-p, p-Pb, or Pb-Pb collisions. To ensure similar event structure and minimize distortions due to different acceptances, events that are selected for mixing have similar selection cuts [8]. The selection of events used for mixed-events must follow the following criteria:

- The selected mixed-event must have a V_z within ± 1 cm of the initial event
- The selected mixed-event must have its multiplicity within 5 % of the multiplicity percentile of the initial event

Having the two events be similar to each other in terms of their multiplicity percentile or position of V_z , allows for the distribution to less represent a random distribution of particles and more represent a true distribution of the particles in question that is only limited by the kinematics of the system. Furthermore, limiting the number of events used for mixing allows for a quicker analysis, as opposed to matching every event to every other event in the billions of collisions seen at the LHC. It is true that having more statistics is always encouraged when it comes to analyzing

something rare however, limiting the number of events mixed allows for a sufficient amount of statistics to be used without taking an impossibly large amount of time and memory.

Another possibility of how to create a background distribution of Λ and $K^\mp$ without it reconstructing to the $\Xi(1820)$ is to use data from the same event, but different charges. To create a “same-event” background that can be used to estimate the background distribution of the initial distribution, special attention to the decay channel must be given. As stated before, the $\Xi(1820)^\mp$ is reconstructed from the decay channel of $\Lambda+K^-$ in the normal particle case, or $\bar{\Lambda}+K^+$ in the anti-particle case. So it stands to reason that making an invariant mass plot using the decay channels of $\Lambda+K^+$ and $\bar{\Lambda}+K^-$ should not reconstruct to the $\Xi(1820)^\mp$, but instead simply create a distribution that corresponds to the actual distribution of $\Lambda+K^\mp$ without any $\Xi(1820)$. Furthermore, because these Λ and $K^\mp$ come from the same event, they should match better to the initial distribution since they have many of the same correlations.

The same-event background distribution used for this analysis is constructed using the geometric mean of the two distributions that should not reconstruct to the $\Xi(1820)^\mp$. Using the equation:

$$Background = 2 * \sqrt{(\Lambda K^+) * (\Lambda K^-)} \quad (38)$$

the same-event background is constructed and found to be very similar to the raw data’s distribution as seen in Figure 4.3.

This same-event background method is only useable for the $\Xi(1820)^\mp$ case, since $\Xi(1820)^0$ decays into two neutral particles, one of which is the K_s^0 which is its own anti-particle, thus making a similar calculation impossible.

4.2 Background Subtraction

Now that the signal has been observed and the combinatorial background has been estimated using either the same-event or mixed-event background, the process to extract the signal can begin. The main purpose of this entire process is to remove the contribution that the combinatorial background

has to the raw signal. This is primarily done so that the true yield results can be easier to extract, but also so that the raw signal can be better fit so as to determine the properties of the $\Xi(1820)$ resonance.

As stated in Section 4.1.6, the same-event and mixed-event backgrounds were constructed so that they would better match the shape of the original distribution. However, given how they are based on the previous description of the procedures, they will simply not match perfectly to the original distribution. Given the fact that the mixed-event distribution was created with about five times the statistics of the original distribution and the same-event distribution is simply not exactly the same of the original distribution, this means that both distributions need to be normalized to match the original distribution. This will allow for the distributions to better match the original distribution and further complement the process of removing the background from the calculations.

The distributions of the mixed-event and same-event background were normalized to match the integrals of the original distribution in mass regions of $1.70 < M_{\Lambda K} < 1.74 \text{ GeV}/c^2$ and $2.0 < M_{\Lambda K} < 2.1 \text{ GeV}/c^2$ respectively, i.e. just below and above the signal region, but sufficiently far away to not interfere with any kind of signal extraction. The main reason for the selection of these normalization ranges, as will be discussed in Section 6.2, is that these ranges corresponded to the lowest χ^2 for the fit performed at the end of the analysis.

The uncertainty of the selection for the normalization range was estimated by varying the normalization ranges as shown in the systematic uncertainty Section 6.2.

Figure 4.4 shows the mixed-event and same-event distribution when normalized to the region of $1.85 < M_{\Lambda K} < 1.95 \text{ GeV}/c^2$ from the raw signal. This was done so that both distributions can be seen clearly in the histogram and compared to each other. When comparing the same-event and mixed-event distribution of the $\Lambda(\bar{\Lambda})K^\mp$ invariant mass plot for Figure 4.4, it is observed that the same-event background distribution matches the initial distribution much better than the mixed-event distribution. This is because long range correlations between strange particles in the same

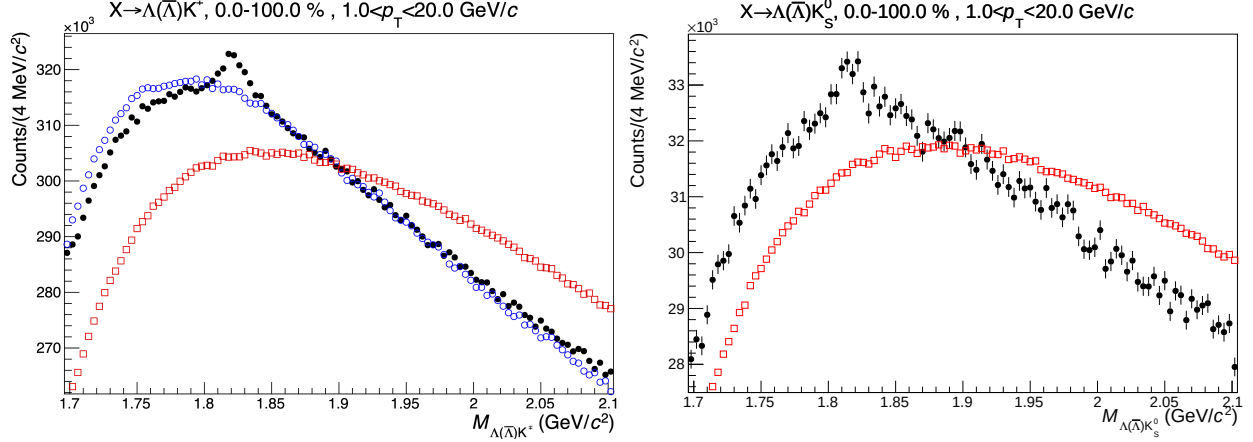


Figure 4.4: Left: Invariant mass plot for $\Lambda(\bar{\Lambda})K^\pm$ for p-p $\sqrt{s} = 13$ TeV minimum bias data. Right: Invariant mass plot for $\Lambda(\bar{\Lambda})K_s^0$ for p-p $\sqrt{s} = 13$ TeV minimum bias data. Blue data points correspond to same-event background. Red data points correspond to mixed-event background. Both distributions have been normalized to a region of $1.85 < M_{\Lambda K} < 1.95$ GeV/ c^2 .

event can not be captured using the mixed-event method. This can be seen in Eq. 39

$$M^2 = 2\cosh\Delta y \sqrt{p_{T,a}^2 + m_a^2} \sqrt{p_{T,b}^2 + m_b^2} - 2\cos\Delta\phi p_{T,a}p_{T,b} + m_a^2 + m_b^2 \quad (39)$$

where the difference between the azimuthal angles of the two particles used for reconstruction is represented in the second term as $\Delta\phi$ [49]. For the same-event distribution, the difference in angles between the two particles is enough that averaged over the large amount of data sets used, the contribution should approximate to $\langle \Delta\phi \rangle \sim 1$. For the mixed-event distribution, the difference between the angles can be anything due to the mixed nature of two collisions that are not dependent on each other. So the average of the contributions should approximate to $\langle \Delta\phi \rangle \sim 0$ over the large amount of collisions used. Because of this difference between same-event and mixed-event distributions, the mixed-event distribution is dismissed from the $\Lambda(\bar{\Lambda})K^\pm$ analysis for p-p data.

Other sources of correlations that can lead to the raw distribution and same-event distribution more closely matching each other come from the fact that some contributions are easier to notice when analyzing particles from one event as apposed to mixed-event. These other correlations are:

- Misidentified resonances: e.g. a Σ^* that decays to $\Lambda\pi^-$, but the pion is misidentified as a

kaon.

- Partially reconstructed decays: e.g. another particle decays to Λ , K^- , plus a third particle X , but we don't consider X in our analysis.
- Correlations due to elliptic flow.
- An $s\bar{s}$ pair is produced, then the s becomes part of a Λ and the $\bar{s}$ becomes part of a K^+ . The Λ and the K^+ may be correlated.

In addition, there are of course many random uncorrelated combinations. The mixed-event background has only uncorrelated combinations and completely misses the correlated part of the background. This may or may not be a problem, depending on the relative magnitudes of the correlated and uncorrelated parts. At low p_T , the uncorrelated part dominates and the mixed-event background gives a good description of the background. At higher invariant masses, the slopes of the mixed-event and same event backgrounds seem to become similar, which suggests that the uncorrelated part dominates there as well. However, there seems to be a large correlated contribution in the region under the $\Xi(1820)$ peak in the p_T range of interest, so the mixed-event background is not as useful in this case.

After the background is normalized to match the original distribution, the background is then subtracted from the original distribution. As can be seen in Figure 4.5, the $\Xi(1820)$ is much more noticeable after the normalized background is subtracted. However, with the subtraction of the background come negative counts for the invariant mass plot. These negative values are due to the nature of the subtracted same-event distribution to over-estimate the distribution when normalized to the initial distribution. Even though the same-event distribution matches the initial distribution very well, it is not perfect and this over-estimation causes the negative value when subtracted. The green line shown in the right figure of Figure 4.5 is the polynomial used to estimate the remaining background distribution that was not completely removed even after the background subtraction, referred to as the residual background. The red line corresponds to a Breit-Wigner fit added to the polynomial. This red fit is only used as a check and reference and is not used in any calculations

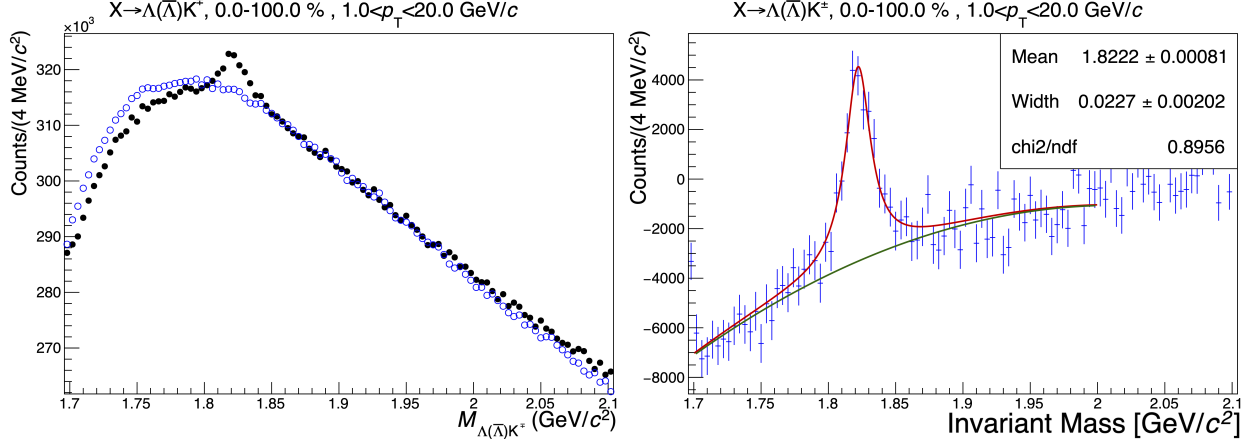


Figure 4.5: Left: Invariant mass plot for $\Lambda(\bar{\Lambda})K^\pm$ before background subtraction for p-p $\sqrt{s} = 13$ TeV minimum bias data. Right: Invariant mass plot of $\Lambda K^\pm$ after background subtraction for p-p $\sqrt{s} = 13$ TeV minimum bias data. Same-event background that has been normalized to a region of $2.0 < M_{\Lambda K} < 2.1 \text{ GeV}/c^2$ was subtracted from the raw data.

or analysis. This is because, as will be shown, there is still more to do in terms of extraction of the signal.

Despite the subtraction, the distribution is still not perfectly shaped like a Breit-Wigner function. This is mainly because of the contributions from residual correlated pairs, e.g. from decays of other resonances, misidentified decay daughters, and jets. Thus the distribution is fit with a 2nd degree polynomial over a range of $1.70 < M_{\Lambda K} < 2.00 \text{ GeV}/c^2$ that excludes the range $1.78 < M_{\Lambda K} < 1.86 \text{ GeV}/c^2$ ($1.75 < M_{\Lambda K} < 1.89 \text{ GeV}/c^2$ for Pb-Pb $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ data) which corresponds to the signal range. The excluded range is selected to be equivalent to the hypothetical range of the signal corresponding to approximately a four sigma deviation of the mass peak according to the PDG [12]. According to the PDG, with a mass of approximately $1.823 \text{ GeV}/c^2$ and a full-width half-maximum of $24 \text{ MeV}/c^2$, after converting the FWHM to sigma using Eq. 3 the range of a four sigma deviation from the mean would be $1.823 \pm 0.04 \text{ GeV}/c^2$ ($1.823 \pm 0.07 \text{ GeV}/c^2$ for Pb-Pb $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ data).

After the polynomial is selected and fit for the given range, the polynomial is subtracted from the aforementioned distribution resulting in a final distribution that contains only the signal with no polynomial or background effects. Figures 4.6 show the “polynomial subtraction” and how the

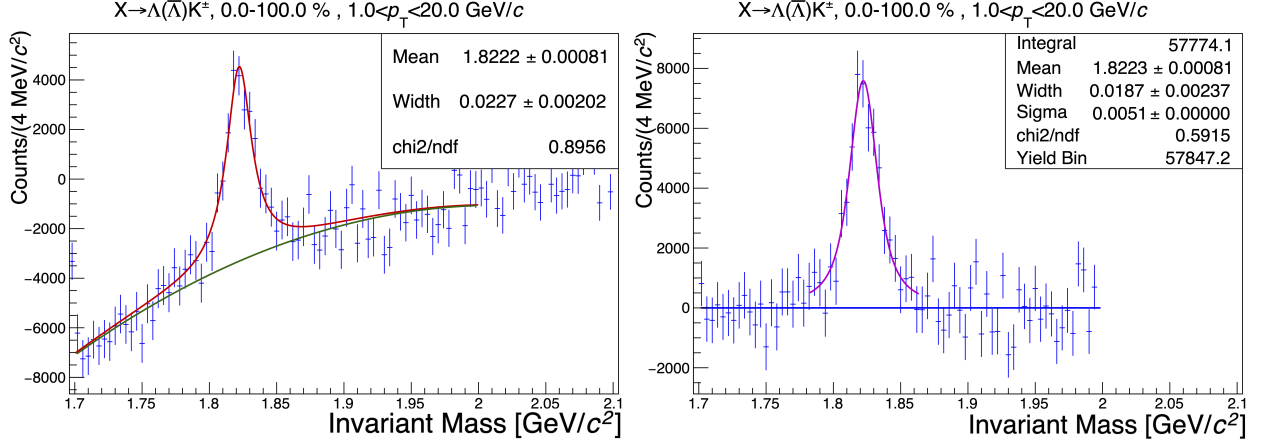


Figure 4.6: Left: Invariant mass plot for $\Lambda(\bar{\Lambda})K^\mp$ before polynomial subtraction. The green line shows the 2nd-degree polynomial used to estimate the residual background. Right: Invariant mass plot of $\Lambda K^\mp$ after polynomial subtraction. The purple line represents the Voigtian fit used to determine the mean, width, and yield of the fit.

peak becomes much more discernible from the process.

4.3 Fit Procedure

The final distribution is fit with a Voigtian function to be used for yield, mean, and width analysis.

The Voigtian function:

$$\frac{dN}{dm} = \frac{C\Gamma_0}{(2\pi)^{3/2}\sigma} \int_{-\infty}^{\infty} \exp\left[-\frac{(m-m')^2}{2\sigma^2}\right] \frac{1}{(m' - M_0)^2 + \Gamma_0^2/4} dm' \quad (40)$$

is a convolution of a Breit-Wigner, used to represent the resonance, and a Gaussian, used to account for the detector resolution [50].

The choice to use the Voigtian function instead of the previously mentioned Breit-Wigner function comes down to one factor. The width of the resonance may be the same order of magnitude as the resolution of the detectors [50]. While the $\Xi(1820)$ is documented as having a natural width of $24 \text{ MeV}/c^2$, there is still enough uncertainty in that value so that the actual width could be closer to the resolution of the detector [12]. To make sure the fit function used for this analysis accounts for the resolution of the detector, the Gaussian component is convoluted into the Breit-Wigner

function to create the Voigtian function.

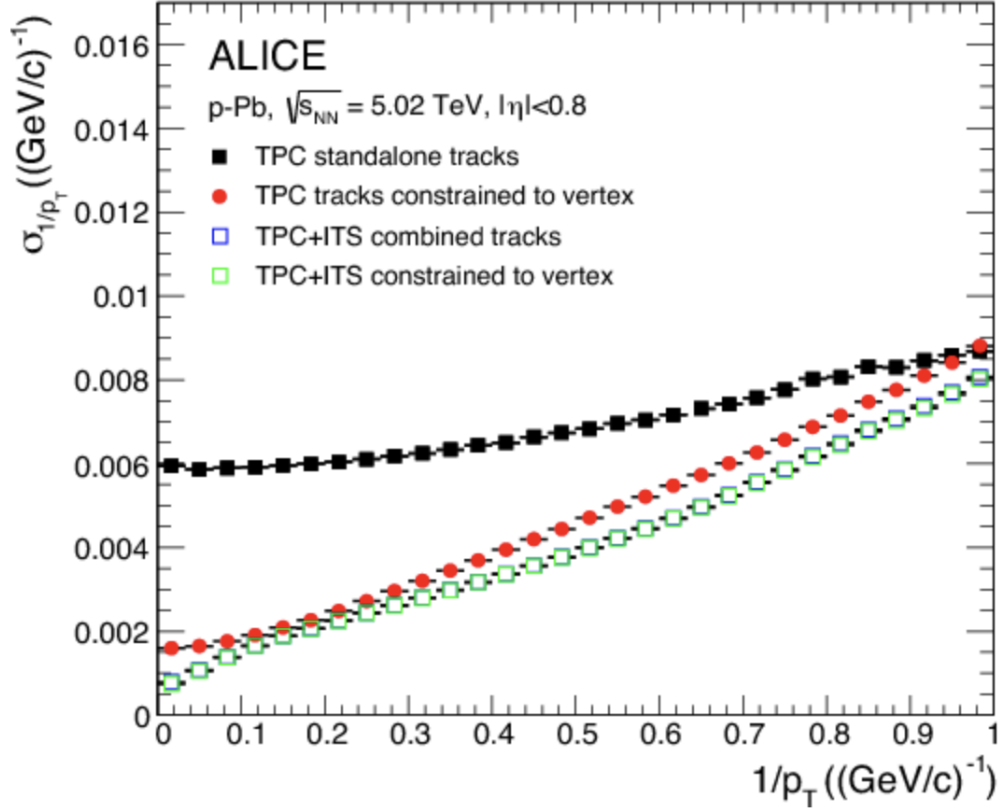


Figure 4.7: p_T resolution for p-Pb collisions using different combinations of tracking detectors [28].

As seen in Figure 4.7, the p_T resolution for a particle is dependent on the detectors used and the p_T of the particle. Particles with lower p_T correspond to lower relative uncertainty. In the range of inverse p_T ($\frac{1}{p_T}$) shown in Figure 4.7, the relative resolution of the p_T is about 0.9 % to 0.2 % between 1 to 10 GeV/ c respectively. While the $K^\pm$ tracks require at least one hit in the ITS, proton and π tracks from the Λ decay do not necessarily have ITS points and must be assumed to use only the TPC-standalone resolution.

The σ of the Gaussian is appropriate for determining the broadening of the signal due to detector resolution. This σ parameter of the Voigtian function is controlled by the momentum resolutions of the $\Xi(1820)$ decay products, although not necessarily in a simple way. Due to the fact that both σ of the Gaussian and Γ of the Breit-Wigner describe the width of their corresponding functions,

the σ is fixed to a value derived from Monte Carlo simulated collisions. This prevents the σ from interfering with the calculations of the free Γ whose value may change.

In the Monte Carlo data, as seen in Section 5.1.1, a simulation of $\Xi(1820)$ signals was performed so that it would be possible to estimate several key parameters needed in the calculation of yields. However, among these simulations of the $\Xi(1820)$ signal, there was also a calculation of the mass difference between $\Xi(1820)$ signals that were reconstructed and $\Xi(1820)$ that were generated. This mass difference plot, seen in Figure 4.8, is then fit with a Gaussian function to approximate the shape of the data points. The σ of the Gaussian is then related to the resolution of the detector, allowing for an estimate of the effect of momentum resolution on the invariant mass plot. Because this mass plot is the difference between the generated “true” $\Xi(1820)$ and the reconstructed “calculated” $\Xi(1820)$, it would make sense that the “resolution” of the plot is related to the “resolution” of the detectors used in the simulation.

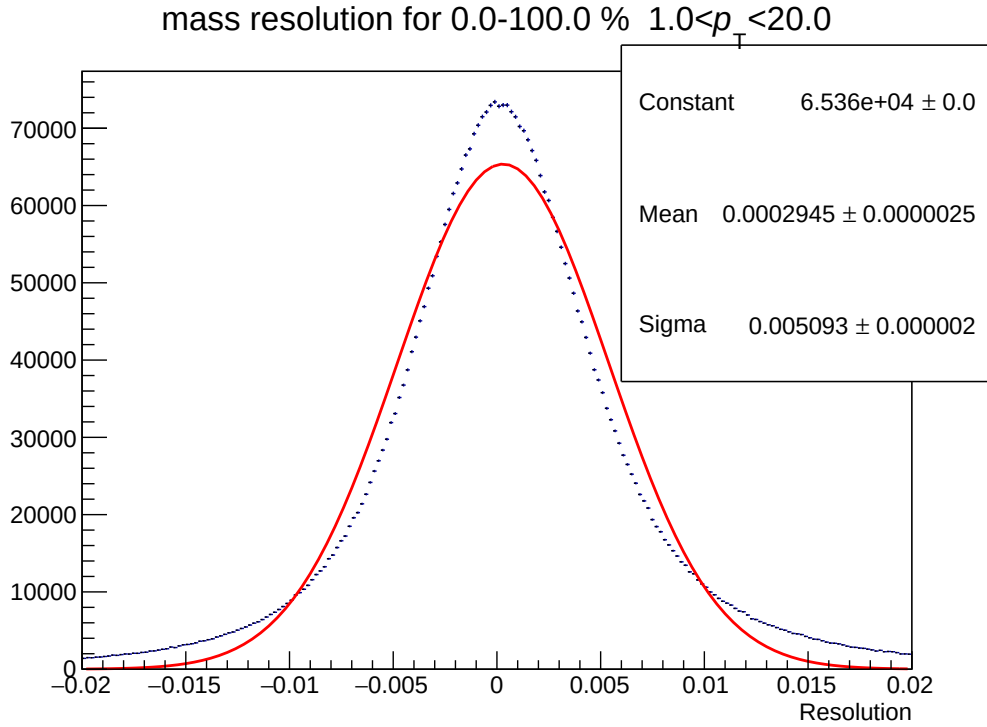


Figure 4.8: Monte Carlo simulation of mass difference between generated and reconstructed $\Xi(1820)^-$. Gaussian fit shown in red.

4.4 Spectrum Extraction

To determine the transverse momentum spectrum of the $\Xi(1820)$, several p_T intervals were constructed. These p_T intervals were initially set up so that each of the individual ranges would contain approximately the same raw yield. This would allow for easier comparison of different p_T intervals and fits since the statistics for the $\Xi(1820)$ are not large enough for a very finely binned analysis.

As seen in Figure 4.6, the raw yield of $\Xi(1820)$ is approximately 55-60000 counts for p-p collisions at $\sqrt{s} = 13$ TeV minimum bias from $1.0 < p_T < 20.0$ GeV/ c in 1.97 billion events. Thus five p_T intervals were constructed so that each had approximately 10000 counts for minimum bias p-p collisions at $\sqrt{s} = 13$ TeV as seen in Figure 4.9.

The process to extract the yield in each p_T interval will be discussed more in Section 6.1.

4.5 p-Pb and Pb-Pb Differences

As will be discussed in Section 7, there are a few key differences between the analysis method previously used for p-p $\sqrt{s} = 13$ TeV data and the procedures for p-Pb and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data.

While most of the procedure to extract the $\Xi(1820)$ signal from p-Pb data is similar to p-p data, there is one very important difference. Instead of using the same-event background to estimate the combinational background of the original distribution, the mixed-event background is used. This will be discussed in more detail in Section 7, but the main reason to use mixed-event instead of same-event background is due to the fact that the polynomial subtracted plots are very similar to each other. So similar in fact, that the only differences between the two plots of p-Pb $\sqrt{s_{NN}} = 5.02$ TeV data for 0-100% multiplicity from Figure 4.10 comes from the visual match of the background distributions to the initial distribution, χ^2/ndf , and the statistical uncertainties. Figure 4.11 for 0-30 % multiplicity percentile shows similar fits, but the statistical uncertainties are noticeably lower for the mixed-event background as opposed to the same-event background. The choice to look at p-Pb $\sqrt{s_{NN}} = 5.02$ data for 0-100 % and 0-30 % multiplicities will be discussed in Section 7.

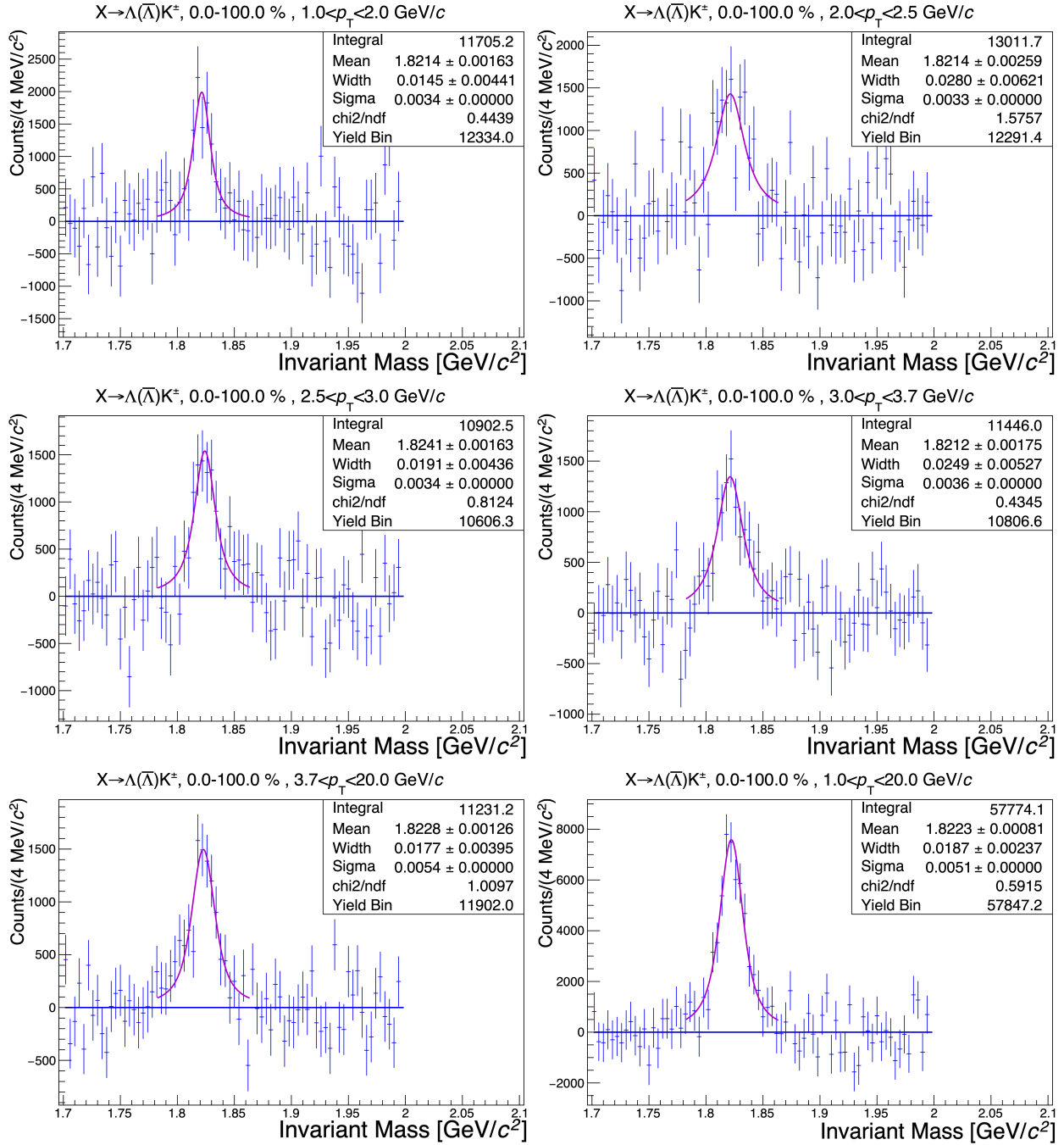


Figure 4.9: Invariant mass plots of $\Lambda(\bar{\Lambda})K^{\pm}$ for various p_T intervals for p-p collisions at $\sqrt{s} = 13$ TeV minimum bias from $1.0 < p_T < 20.0$ GeV/c.

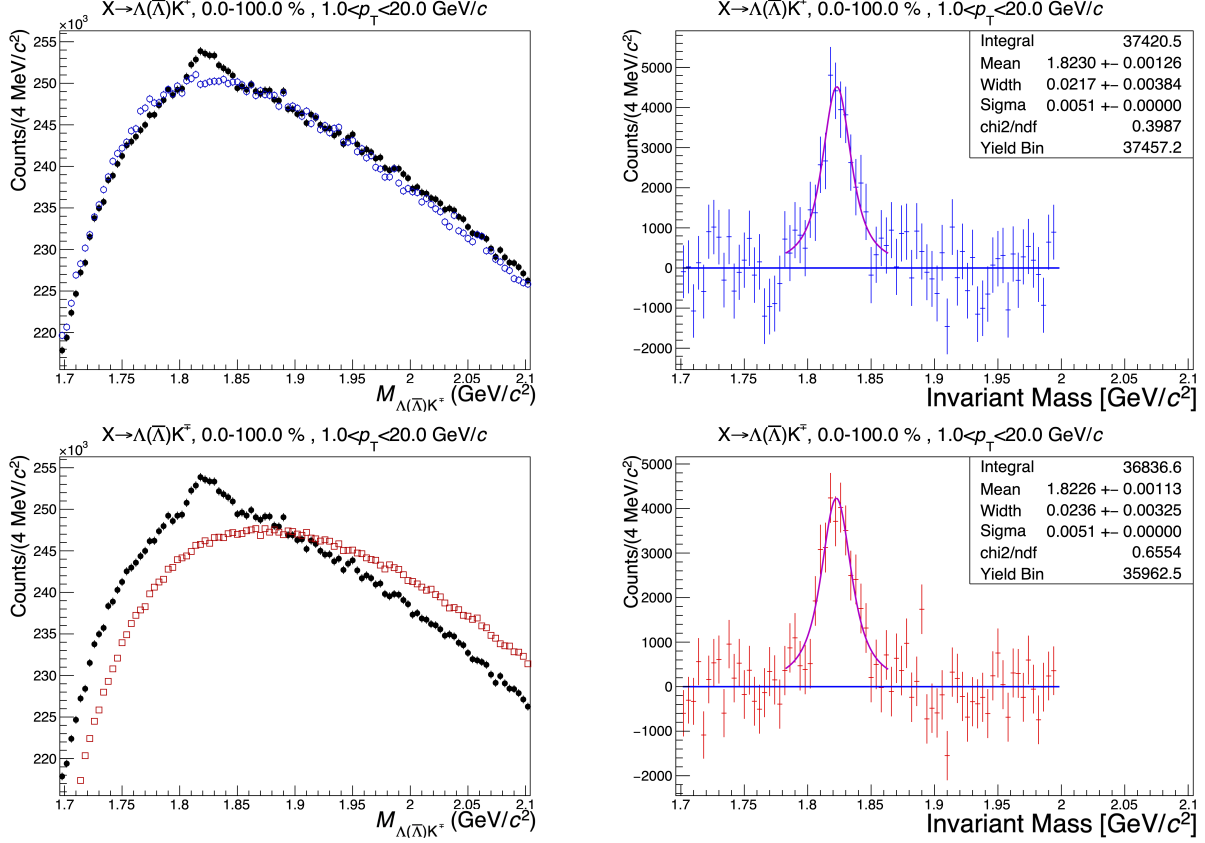


Figure 4.10: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^{\mp}$ for 0-100% for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Normalization ranges used for left plots are not those used for calculations and subtraction for right plots. Top left: Raw data and same-event background. Top right: Invariant mass plot with background subtraction of same-event background and residual polynomial background implemented. Bottom left: Raw data and mixed-event background. Bottom right: Invariant mass plot with background subtraction of mixed-event background and residual polynomial background implemented.

Do note that the left plots of Figures 4.10 and 4.11 have been normalized using a normalization region of $1.85 < \Lambda K^{\mp} < 1.95$ GeV/ c^2 so that each background can be seen clearly in the plots. The “default” normalization region for these same-event and mixed-event background distributions is determined to be $1.70 < \Lambda K^{\mp} < 1.74$ GeV/ c^2 , however the use of this normalization range would show the background distributions much higher than the raw data and possibly not visible on the same plot, especially for the mixed-event background. The use of normalization region of $1.85 < \Lambda K^{\mp} < 1.95$ GeV/ c^2 for the left plots is simply to have the backgrounds visually match the raw data.

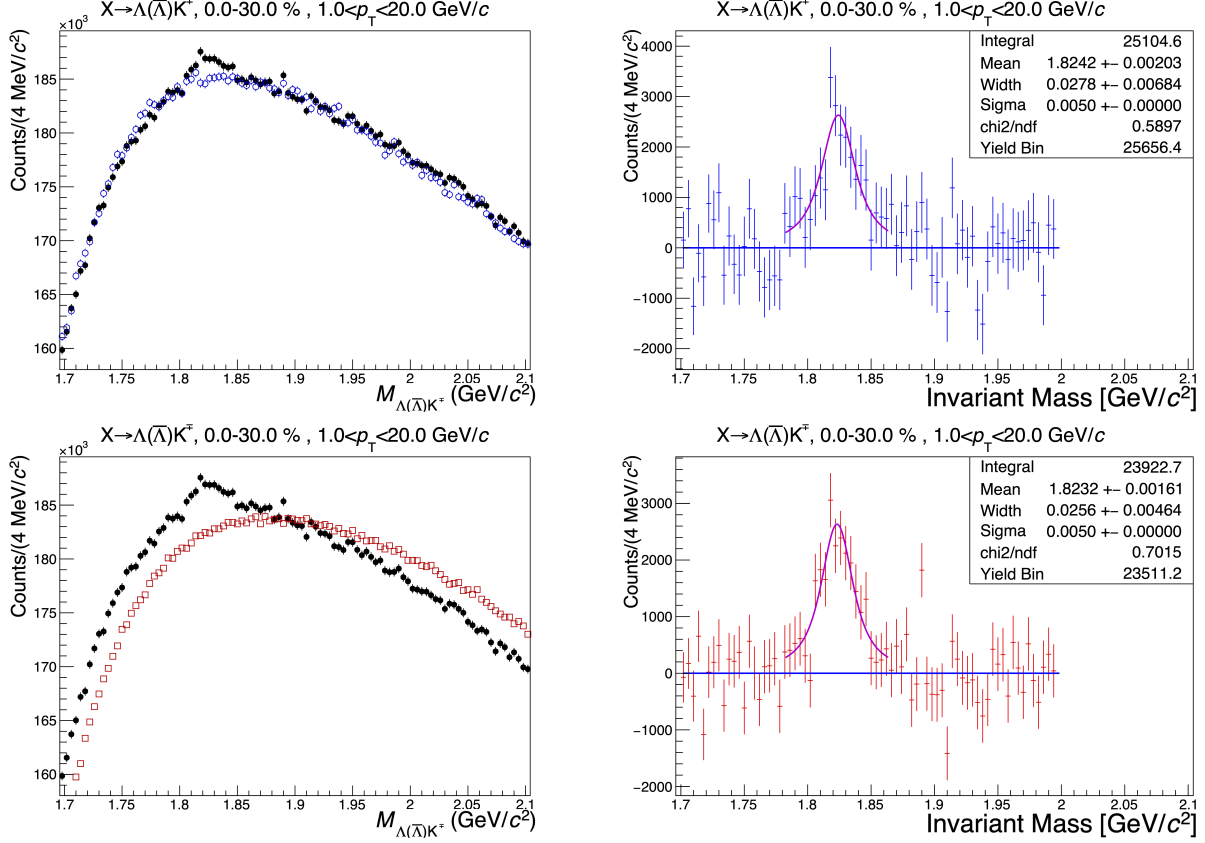


Figure 4.11: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^{\mp}$ for 0-30% for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Normalization ranges used for left plots are not those used for calculations and subtraction for right plots. Top left: Raw data and same-event background. Top right: Invariant mass plot with background subtraction of same-event background and residual polynomial background implemented. Bottom left: Raw data and mixed-event background. Bottom right: Invariant mass plot with background subtraction of mixed-event background and residual polynomial background implemented.

While it is true that the same-event background distributions for Figures 4.10 and 4.11 do have a lower χ^2/ndf when compared to the mixed-event distributions, and they seem to visually match the invariant mass plots better, it is the statistical uncertainty that caused the final decision to use the mixed-event background distribution as the “default” background. For both of the mixed-event background plots for Figures 4.10 and 4.11, the statistical uncertainty of mean and width of the Voigtian fits is lower when compared to the same-event plots. Furthermore, the mixed-event background matches the p-Pb data much more than mixed-event data for p-p data as discussed in Section 7.

It is expected that the mixed-event backgrounds will match the raw data distributions better with increasing multiplicity. With that in mind, and given that the errors for individual data points for mixed-event background will be about 0.45 times that of the same-event background, the choice to use the mixed-event background becomes more reasonable.

The procedure to extract $\Xi(1820)$ signal from Pb-Pb data contains two major differences from the procedure for p-p data.

First, as will be shown shortly, the polynomial fit for the background subtracted invariant mass plots has a different range of exclusion that corresponds to the signal range. Due to the wider signal for the $\Xi(1820)$ in these plots, the range that corresponds to the signal range is changed from a four sigma deviation from the mean to a seven sigma deviation from the mean. The previous four sigma deviation does not cover the entire signal, and as such can not be considered a fit for the $\Xi(1820)$ signal in question.

Second, the mixed-event background is used to estimate the combinational background as opposed to the same-event background. The reasons for this decision will be discussed in Section 4.5.1.

4.5.1 False peaks in Pb-Pb data

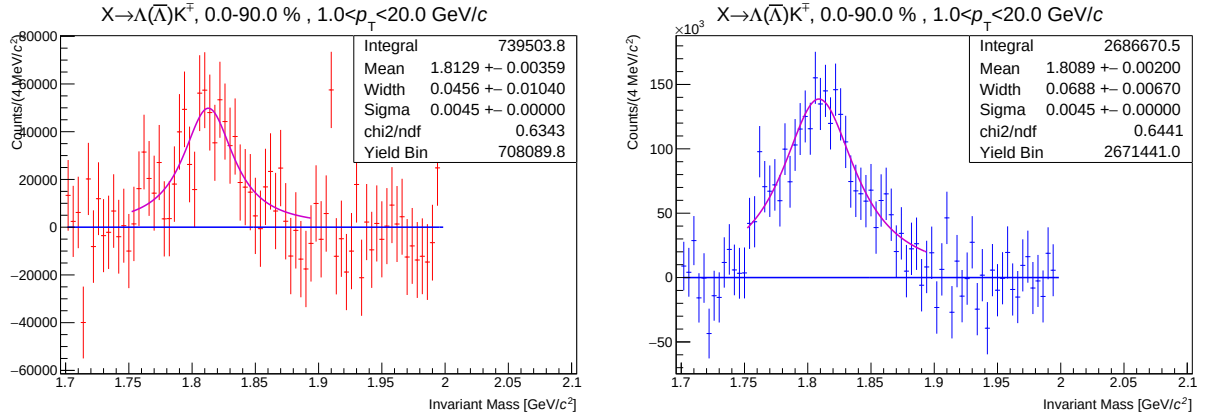


Figure 4.12: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^{\mp}$ for kINT7 triggered data (0-90%) for Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Invariant mass plot with background subtraction of mixed-event or same-event background and residual polynomial background implemented. Left: Mixed-event background. Right: Same-event background.

Upon first inspection of the invariant mass plots for Pb-Pb data, as seen in Figure 4.12, it would seem logical to assume that the same-event background plot is superior to the mixed-event background plot. The same-event background plot has the more noticeable signal, smaller uncertainties, and less statistical variation when compared to the mixed-event plot. In fact, taking the same-event background plot as is shown would seem to indicate several features of chiral symmetry restoration, such as the lower mean and higher width. However, the drastic size of the peak caused some concern as to the validity of the signal. With that, the Pb-Pb data was investigated more thoroughly. Several notable features were noticed that cast doubt on the validity of the same-event background plots.

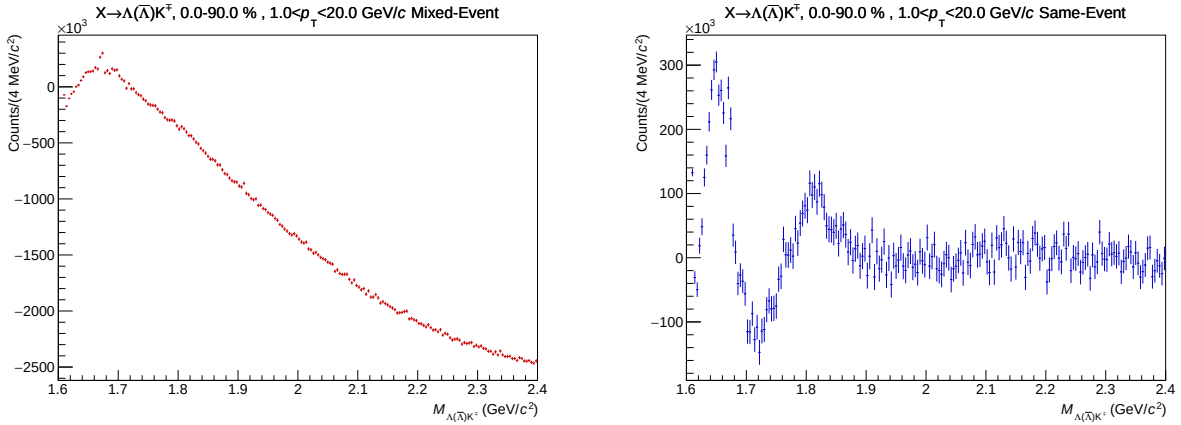


Figure 4.13: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^\pm$ for kINT7 triggered data (0-90%) for Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Invariant mass plot with background subtraction of mixed-event or same-event background. Left: Mixed-event background. Right: Same-event background.

First, the full range of the invariant mass plots with mixed-event and same-event background subtraction are investigated to make sure that the large signal from the same-event background is visible. As seen in Figure 4.13, there appears to be two “peaks” for the same-event background subtracted plots, one at around $1.650 \text{ GeV}/c^2$ and another around $1.820 \text{ GeV}/c^2$. These two peaks would seem to correspond to the Ω particle and the $\Xi(1820)$, but there are a few key points that call into question the validity of these peaks. Most notably, the large peak at $1.650 \text{ GeV}/c^2$ for the same-event background seems too large to correspond to the Ω particle. In fact, closer observation

of the same-event background seems to show the Ω peak as 3 data points just to the right of the peak at $1.650 \text{ GeV}/c^2$. This is similar to the much smaller peak seen in the mixed-event background around $1.672 \text{ GeV}/c^2$. Thus the “false” peak at $1.650 \text{ GeV}/c^2$ is called into question and also raises concern about the peak at $1.820 \text{ GeV}/c^2$.

Observation of the Pb-Pb $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ data for various p_T bins, shown in Figure 4.14, shows a very striking feature of the “false” peaks. They seem to move with changing p_T ranges. Lower p_T ranges show multiple “false” peaks, and as the p_T ranges increase, the peaks seem to shift as well. This shows the “false” peaks as not corresponding to particles, but instead to some kinematic feature or correlation.

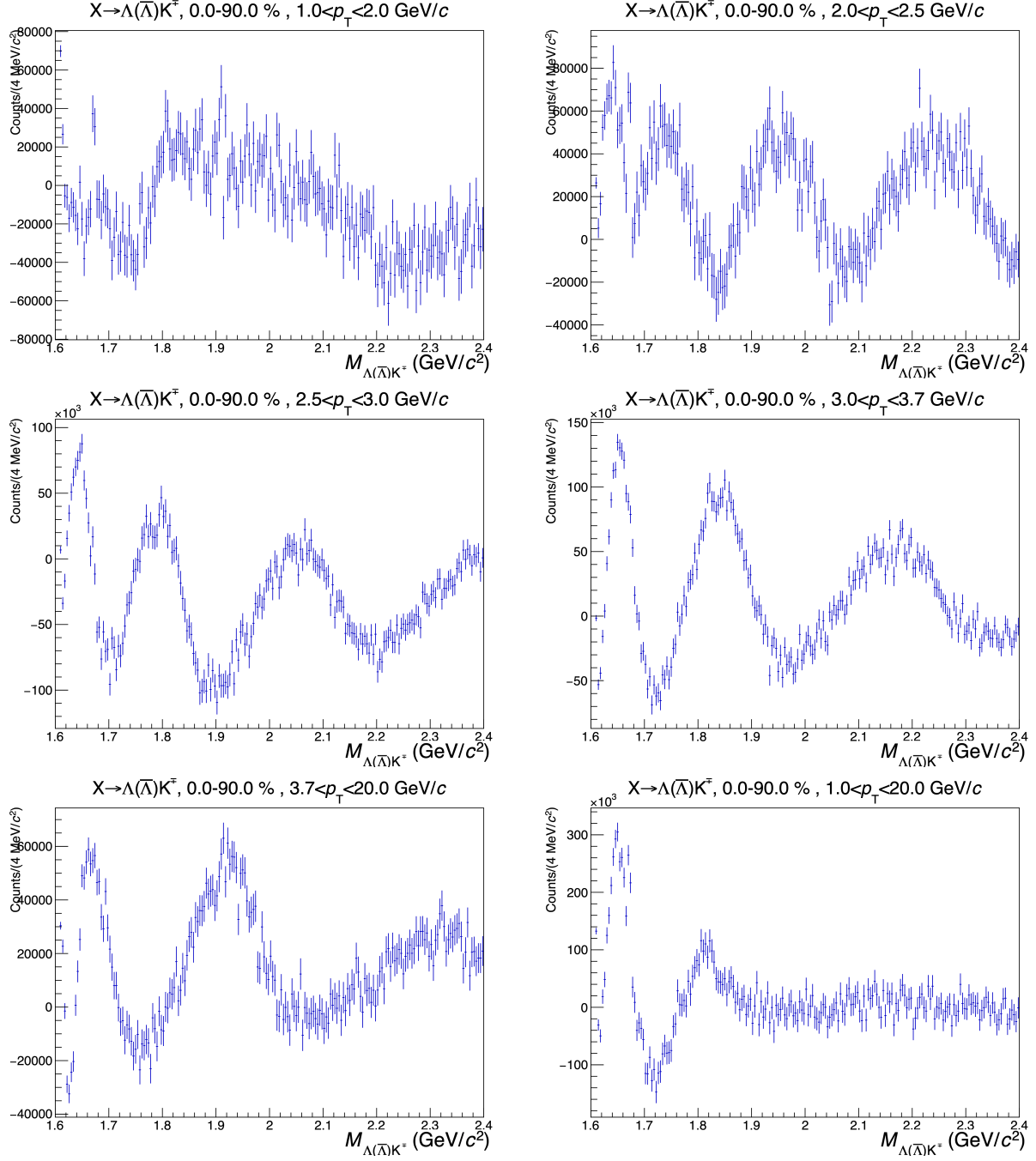


Figure 4.14: Invariant mass same-event background subtracted plots of $\Lambda(\bar{\Lambda})K^\pm$ for various p_T intervals for Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV kINT7 triggered minimum bias from $1.0 < p_T < 20.0$ GeV/ c .

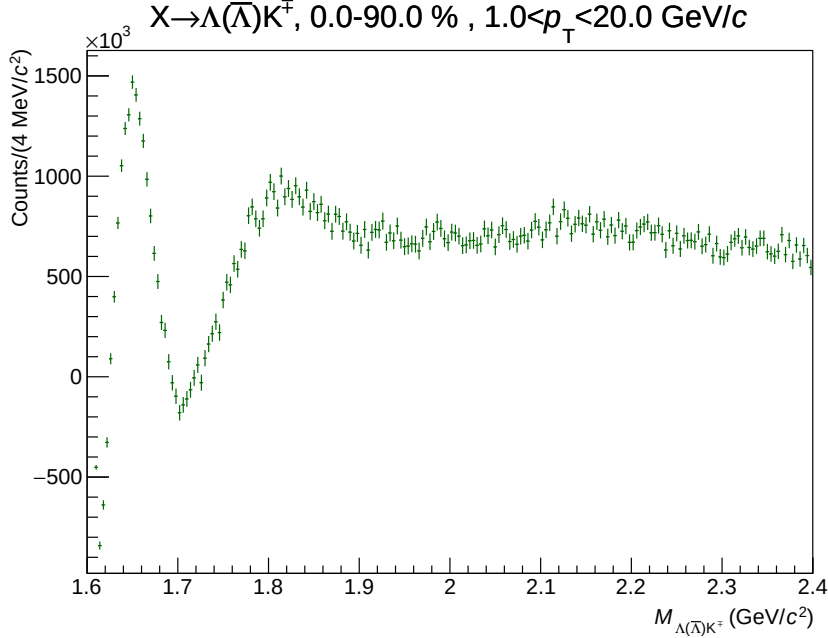


Figure 4.15: Invariant mass plots of mixed-event $\Lambda(\bar{\Lambda})K^\mp$ subtracted by mixed-event $\Lambda(\bar{\Lambda})K^\pm$ distribution Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV kINT7 triggered minimum bias from $1.0 < p_T < 20.0$ GeV/c.

The biggest evidence against the use of same-event background is from the invariant mass plots of mixed-event ΛK^- combined with mixed-event $\bar{\Lambda} K^+$ after being subtracted by mixed-event ΛK^+ combined with mixed-event $\bar{\Lambda} K^-$ distribution (mixed-event $\Lambda(\bar{\Lambda})K^\mp$ minus mixed-event $\Lambda(\bar{\Lambda})K^\pm$) seen in Figure 4.15. A mixed-event distribution should have no reconstructed peaks, let alone when subtracted by another mixed-event distribution. This is confirmed when there appears to be no 3 data points corresponding to the Ω peak just to the right of the $1.650 \text{ GeV}/c^2$ peak. The fact that these “false” peaks still appear in plots where there should be no peaks, means that plots with these “false” peaks can be rejected.

Similar to Figures 4.10 and 4.11 for p-Pb data, Figures 4.16, 4.17, and 4.18 show the mixed and same-event background distributions for Pb-Pb data. However, unlike the previous comparison of mixed-event and same-event backgrounds, the Pb-Pb data shows some unique features.

The most notable feature of the right plots for Figures 4.16 and 4.17 is the difference between

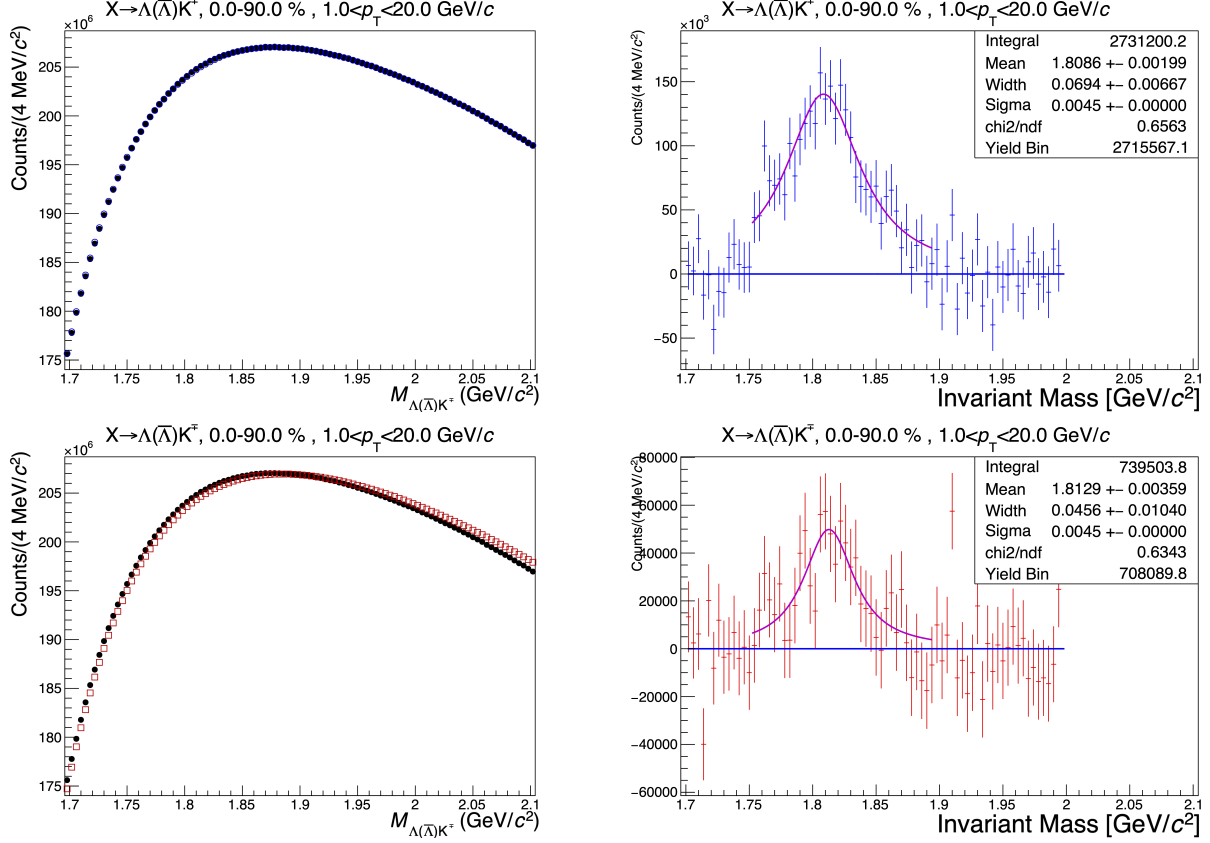


Figure 4.16: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^\pm$ for kINT7 (0-90 %) triggered Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Normalization ranges used for left plots are not those used for calculations and subtraction for right plots. Top left: Raw data and same-event background. Top right: Invariant mass plot with background subtraction of same-event background and residual polynomial background implemented. Bottom left: Raw data and mixed-event background. Bottom right: Invariant mass plot with background subtraction of mixed-event background and residual polynomial background implemented.

the same-event and mixed-event background subtracted plots. As discussed previously, the same-event background subtracted plots appear to contain some type of “false” peak that is masking any “true” $\Xi(1820)$ signal. It is for this reason, as well as the other reasons listed previously, that the same-event background distributions are rejected from the Pb-Pb analysis. Surprisingly, Figure 4.18 for kSemiCent triggered data does not contain the “false” peaks seen in the other triggered data. In fact, both distributions look incredibly similar, even having similar fits and parameters. Still, so as to not change the background selection procedure only for this triggered data, the mixed-event background plot that shows the lower statistical errors is selected.

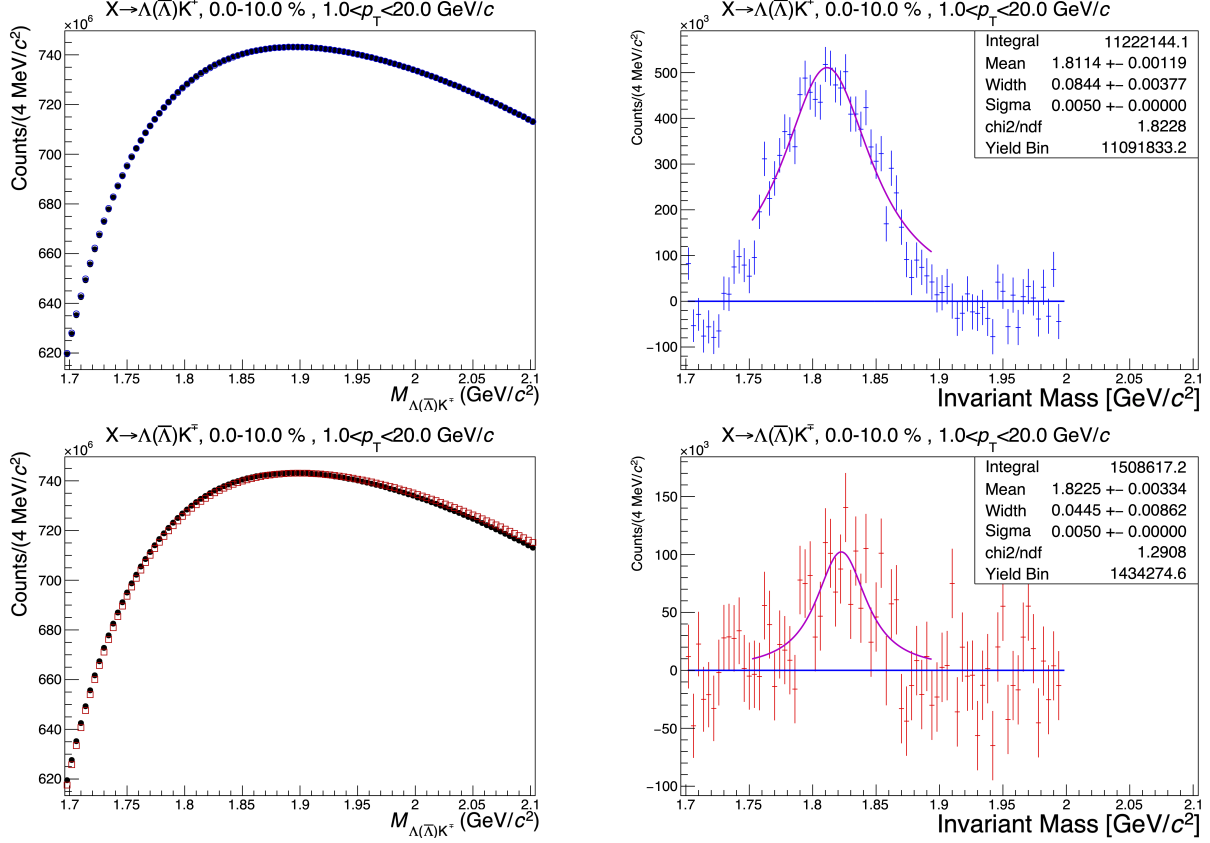


Figure 4.17: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^{\mp}$ for kCent (0-10 %) triggered Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Normalization ranges used for left plots are not those used for calculations and subtraction for right plots. Top left: Raw data and same-event background. Top right: Invariant mass plot with background subtraction of same-event background and residual polynomial background implemented. Bottom left: Raw data and mixed-event background. Bottom right: Invariant mass plot with background subtraction of mixed-event background and residual polynomial background implemented.

Another notable feature of the left plots for Figures 4.16 and 4.17 is that there is almost no difference between the same-event and mixed-event background distributions when compared to the raw distribution. As stated before in the p-Pb data comparison, as multiplicity increases, mixed-event data better matches the raw distribution. In fact, the $\Xi(1820)$ signal can not be seen in the left plots. However, given the statistics of the background subtracted right plots of Figures 4.16 and 4.17 having data points almost 1 % of the yield from the initial left plots, it would be difficult to notice a peak that is 1 % higher than the background for the left plots.

Furthermore, given that each of the triggered datasets for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV contains

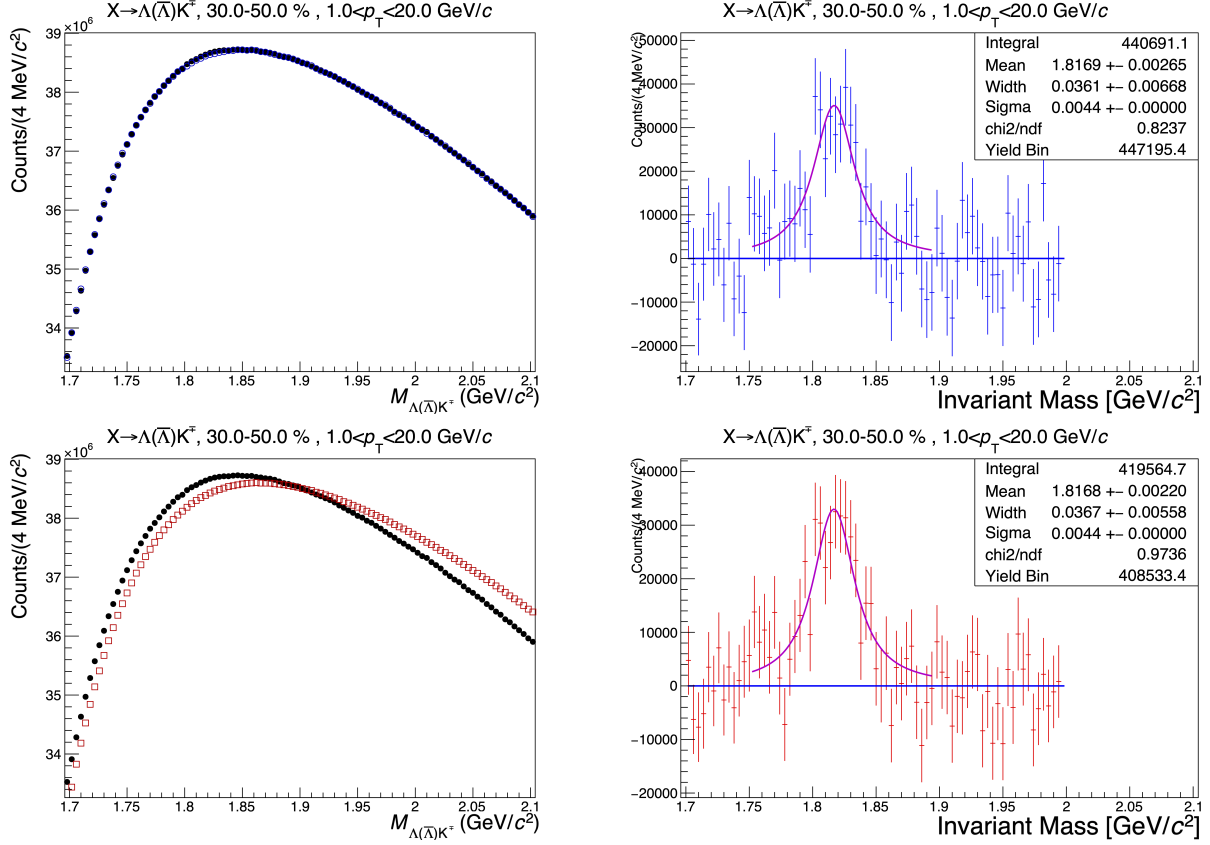


Figure 4.18: Comparison of mixed- and same-event backgrounds for invariant mass plots of $\Lambda(\bar{\Lambda})K^\pm$ for kSemiCent (30-50 %) triggered Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Normalization ranges used for left plots are not those used for calculations and subtraction for right plots. Top left: Raw data and same-event background. Top right: Invariant mass plot with background subtraction of same-event background and residual polynomial background implemented. Bottom left: Raw data and mixed-event background. Bottom right: Invariant mass plot with background subtraction of mixed-event background and residual polynomial background implemented.

similar statistics ($1.34 \cdot 10^8$ for kINT7, $9.6 \cdot 10^7$ for kCent, and $8.4 \cdot 10^7$ for kSemiCent), the choice to use or not use a specific triggered data set comes down to if the final invariant mass plots actually contain a usable $\Xi(1820)$ signal that can be analyzed.

5 Data Corrections

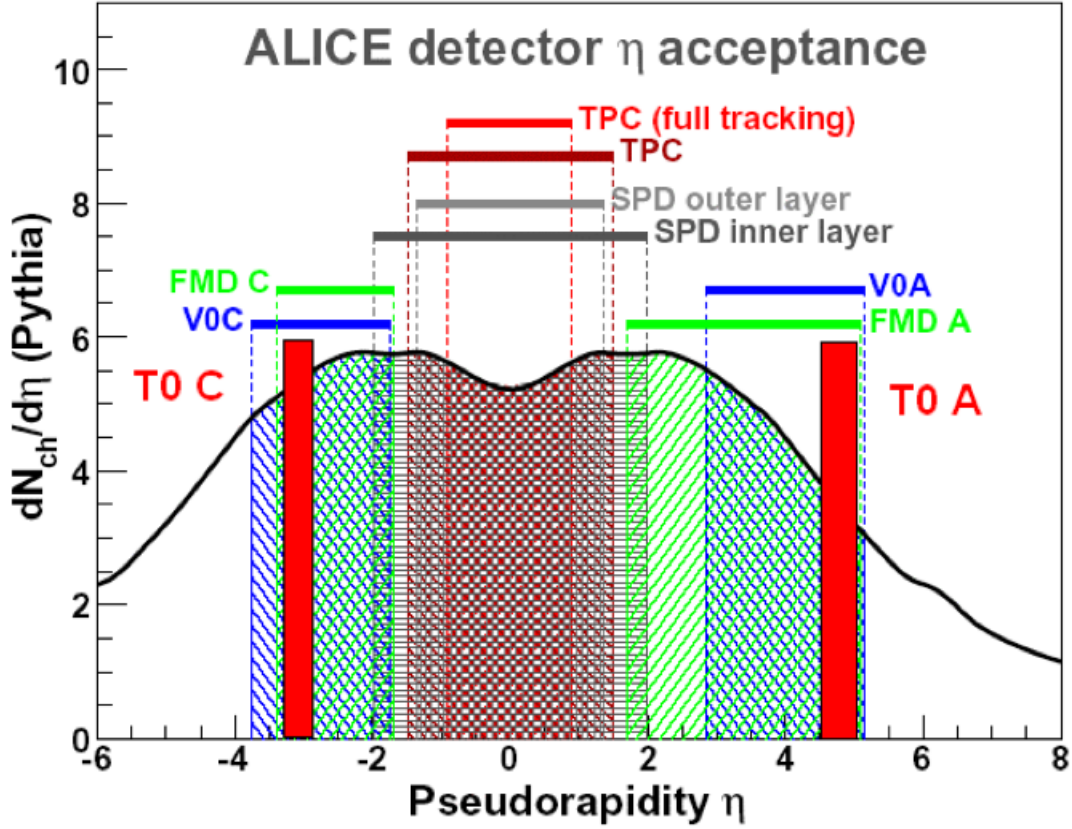


Figure 5.1: Pythia event generator prediction for p-p collisions at $\sqrt{s} \sim 14$ TeV. Superimposed is the coverage of various ALICE detectors for charged particle measurements vs pseudo-rapidity η [51].

Along with the calculations that have already been performed, there is one important factor that must be addressed before any result can be presented. As seen in Figure 5.1, ALICE does not encompass the entire span of conceivable particle tracks. Similarly, no one detector can cover the range of p_T used for the experiment as seen in Figure 2.5.

Due to the geometric limitations of the detectors, there will need to be a correction to any total yield calculation to account for particles that could not be detected simply because they were outside the physical range of the detector. This correction factor is referred to as the “acceptance” and can be calculated using Monte Carlo simulations. This “acceptance”, simply called a , also takes

into account the rapidity cuts and secondary decays of particles used to reconstruct the resonances.

Due to the non-ideal reality of detectors, there is an inherent possibility that not all particles can be tracked, identified, or reconstructed using the ALICE detectors. Thus, even if a certain number of particles are produced, not all will be correctly identified using the detectors or PID identification methods. A correction factor, known as the “efficiency” correction, is necessary to account for this possibility and is calculated using Monte Carlo simulations. This “efficiency”, simply called e or ε_{rec} , also accounts for the tracking efficiency and the cuts used to select good quality tracks coming from the primary vertex.

Despite the complexities, the “acceptance” and “efficiency” are both calculated at the same time using the procedure discussed in Section 5.1.2. However, instead of separately calculating each, the product of “acceptance” and “efficiency” is determined and used for this analysis.

5.1 Efficiency Corrections

5.1.1 Monte Carlo sets

Monte Carlo simulations are computer algorithms that obtain numerical results through the use of repeated random sampling.

Two event generators were used by the ALICE collaboration to simulate high energy collisions and the theoretical system dynamics of these collisions. PYTHIA8 is an event generator used to simulate hadron-hadron or lepton-lepton collisions [52]. Despite the limitation to only be used for simulating p-p collisions, PYTHIA8 is able to simulate several processes such as QCD processes, electroweak processes, Higgs processes, Supersymmetric (SUSY) processes, new gauge boson processes, left-right symmetric processes, resonances, and several more [52]. While PYTHIA8 is compared to physical experiments such as the LHC, a set of reference parameters and variables needs to be provided to better describe existing models and datasets [53]. Some of these lists of constraints, such as Monash 2013, are used to better “tune” the simulations to match the results and expectations of the experiments to which they are compared to [53]. Heavy Ion Jet Interaction

Generator (HIJING) is another event generator used to simulate A+A collisions [54]. While primarily used to simulate Pb-Pb collisions, HIJING does allow for simulations and calculations that simply are not possible in PYTHIA8, such as impact parameter calculations [54].

Another important code needed for simulations is the DPGSIM event generator. The DPGSIM is a steering script from the ALICE software package that takes information about the generator, run, energy, system, detector configuration, and other parameters to set up a complete Monte Carlo simulation.

The Monte Carlo simulation also included an “injected resonance” meaning that the number of resonances that were simulated and added on top of the collision is increased by about 1000 times the estimated amount of resonances normally produced. This “injected resonance” is meant to increase the number of resonances studied in the simulation since some resonances, like the $\Xi(1820)$, are such rare particles that even with a large number of simulated events, some of the properties that are being investigated may not be studied simply because there is not enough statistics. The injected resonances correspond to a flat p_T distribution, so as to increase statistics at high p_T .

To help further simulate collisions and the response of the ALICE detectors to the generated tracks, a computer package known as GEANT3 (Geometry And Tracking) is used to model the ALICE experiment with realistic descriptions of the detector responses.

The Monte Carlo simulation for p-p $\sqrt{s} = 13$ TeV data was constructed using the DPGSIM event generator with GEANT3 package and the Pythia8 Monash2013 package [55]. A sample of approximately 13 % of the number of minimum bias events was simulated to compute the corrections for the corresponding collisions.

The Monte Carlo simulation for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data was constructed using the HIJING event generator and the GEANT3 package [56]. A sample of approximately 2.8×10^5 0-10 % centrality, 8.0×10^5 10-50 % centrality, and 2.8×10^6 50-90 % centrality events were simulated, with the number of injected $\Xi(1820)$ (both positive and negative charged states) resonances for each event to be approximately 40, 10, and 5 respectively for each chosen centrality class.

Within these Monte Carlo simulations are a collection of simulated collisions with varying parameters such as multiplicity and impact parameter. These simulated collisions create “generated” particles that are then simulated to physically react with the modeled detectors of ALICE as they evolve in the system. For example, decaying particles are simulated to decay using the probability of their exponential decay laws, then the decay products are allowed to interact with the detectors used to identify them. The decay products are then reconstructed using the signals produced from the modeled detectors and the same cuts are used to identify such particles in real data. These “reconstructed” particles may not have the same physical properties of their “generated” particles due to the need to use the modeled detectors to reconstruct them. Identifying the difference between the “reconstructed” particles and the “generated” particles is fundamental to determining the effect that detectors have on the particle’s signal.

The main purpose of these Monte Carlo simulations is to determine the detector resolution (as discussed in Section 4.4) and the acceptance and efficiency for yields of the $\Xi(1820)$ (as will be discussed in Section 5.1.2). The detector resolution is determined by calculating the difference in mass between particles that have been “reconstructed” using the modeled detectors and their corresponding “generated” values. The acceptance and efficiency are calculated using the ratio of reconstructed over generated particles in the simulations to determine how many particles were reconstructed as opposed to how many were actually generated. While these values may seem trivial, they are important for any analysis to be compared to other experiments.

5.1.2 Correction Procedure

The product of acceptance and efficiency is calculated as the fraction of generated $\Xi(1820)$ that are reconstructed after passing through the Monte Carlo detector simulation. These $\Xi(1820)$ are subject to the same track quality, PID, topological, and pair rapidity cuts used for the real data analysis.

$$A \times \varepsilon_{rec} = \frac{reconstructed}{generated} \quad (41)$$

The acceptance * efficiency plots (called $a^*_{\epsilon_{rec}}$ for simplicity) are then fit with a linear polynomial after being projected onto their mass axis. Thus, the linear polynomial is a function of mass. As will be discussed later, the $a^*_{\epsilon_{rec}}$ plots can also be projected onto the p_T axis. As seen in Figure 5.2, the linear polynomial fit does seem to match the general trend of the $a^*_{\epsilon_{rec}}$ plot. The constant of the polynomial is interpreted as the average value of $a^*_{\epsilon_{rec}}$, while the slope can be considered a representation of the mass dependence of efficiency correction. The statistical uncertainty of the $a^*_{\epsilon_{rec}}$ plots is determined using standard error propagation for division of a subset.

The Voigtian fit is then modified to account for the mass dependent efficiency correction by including the following addition to the function:

$$Voigt \times [1 + 0.15 * (a * \epsilon_{rec})_{slope} * Erf((x - M)/(0.15 * 1.1283416))] \quad (42)$$

where Erf is the error function of a Gaussian. The use of Erf is to account for the linear variation of the $a^*_{\epsilon_{rec}}$ over the width of the peak, while not introducing unphysical behavior such as negative values for masses far away from the peak. The function was constructed so that the portion in the square brackets has a value 1 and a slope exactly equal to $a^*_{\epsilon_{rec}}$ at $x=M$. The factor 1.1283416 is necessary to achieve this. The 0.15 factor determines the range over which the Erf behaves similar to a straight line, before flattening out to constant values. Choosing 0.15 means the Erf behaves basically like a straight line across the whole $\Xi(1820)$ peak.

While Figure 5.2 shows the $a^*_{\epsilon_{rec}}$ vs. generated mass for one p_T range, it is equally important to understand how $a^*_{\epsilon_{rec}}$ changes across all of p_T . As seen in Figure 5.3, the $a^*_{\epsilon_{rec}}$ increases steadily until reaching a peak of 0.25 around 6 GeV/c. The $a^*_{\epsilon_{rec}}$ then slowly decreases until it reaches the point where no more $\Xi(1820)$ were generated for the Monte Carlo simulation around 15 GeV/c. Since the p_T range for the p-p $\sqrt{s} = 13$ TeV was selected to be 1-20 GeV/c, the value of $a^*_{\epsilon_{rec}}$ determined from the fit of Figure 5.2 of 0.20 seems reasonable as an average for the total $a^*_{\epsilon_{rec}}$.

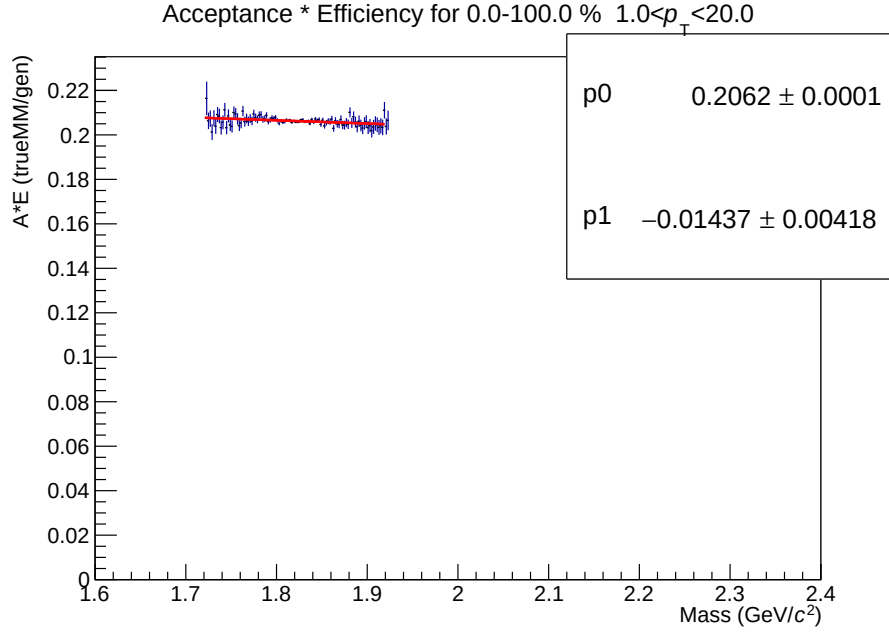


Figure 5.2: (Acceptance $\times$ Efficiency) $a^* \varepsilon_{rec}$ vs. mass plot for $\Xi(1820)^-$ for p-p $\sqrt{s} = 13$ TeV using Monte Carlo simulations. Fit with linear polynomial shown in red.

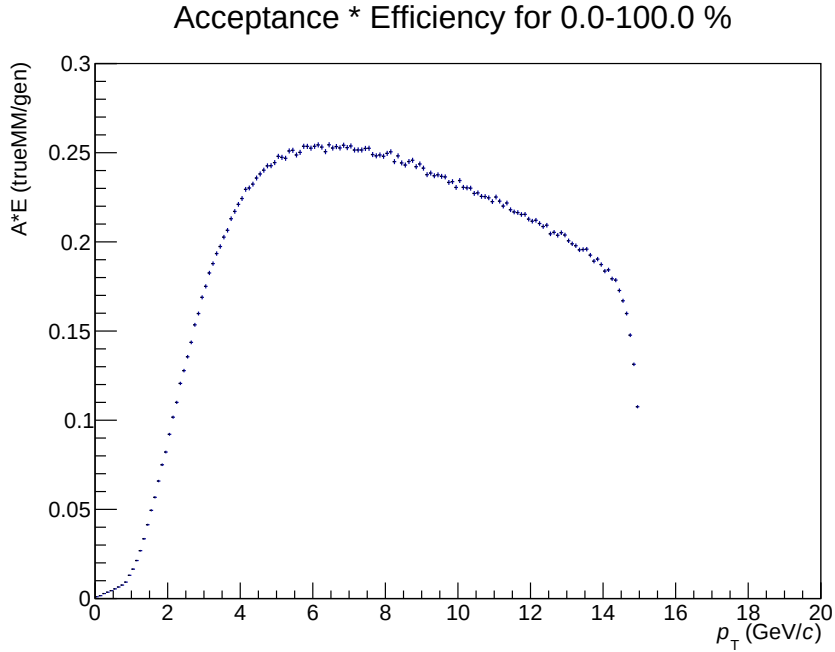


Figure 5.3: (Acceptance $\times$ Efficiency) $a^* \varepsilon_{rec}$ vs. p_T plot for $\Xi(1820)^-$ for p-p $\sqrt{s} = 13$ TeV using Monte Carlo simulations.

5.2 Acceptance Procedure

The p_T spectrum used for this analysis is not enough to determine the entire yield of the $\Xi(1820)$ p-p $\sqrt{s} = 13$ TeV. The spectrum needs to be fit with a function that not only describes the spectrum, but also can be used to calculate the yield from p_T ranges where the detector has no acceptance such as $0 < p_T < 1$ GeV/ c . The standard function that is used to determine the shape of the p_T -spectra for the corrected yields of the $\Xi(1820)$ for p-p $\sqrt{s} = 13$ TeV data is the Levy-Tsallis function. The Levy-Tsallis function is seen in Eq. 43:

$$f(mass, p_T, n, C, norm) = p_T * norm * \frac{(n-1)(n-2)}{nC(nC * mass(n-2))} \frac{1}{(1 + \frac{(m_T - mass)}{nC})^n} \quad (43)$$

where m_T is called the transverse mass ($m_T = \sqrt{p_T^2 + mass^2}$), n and C are fitting parameters, and $norm$ is the normalization constant [57].

To determine the systematic uncertainty of using the Levy-Tsallis function to describe the p_T -spectra, other fit functions are used to determine the difference between each fit as the uncertainty. Two functions in particular are used for this analysis, the Boltzmann function and the Boltzmann-Gibbs Blast-Wave function. The Boltzmann function is shown in Eq. 44

$$\frac{d^2N}{dydp_T} = p_T \frac{dN}{dy} m_T e^{\frac{-m_T}{T}} \quad (44)$$

where T corresponds to the temperature of the “blackbody” used to produce the observed spectra. The Boltzmann-Gibbs Blast-Wave function is shown in Eq. 45

$$f(r, m_T, p_T, \beta_{max}, T, n) = r * m_T * I_0\left(\frac{p_T \sinh(\text{arctanh}(\beta_{max} * r^n))}{T}\right) * K_1\left(\frac{m_T \cosh(\text{arctan}(\beta_{max} * r^n))}{T}\right) \quad (45)$$

where r is the radius, β_{max} is the magnitude of the common velocity, T is the freeze out temperature, and n is the shape of the velocity profile. I_0 and K_1 correspond to modified Bessel functions of the first and second kind respectively. The Boltzmann-Gibbs Blast-Wave function attempts to mimic

radial flow by assuming that all particles move radially outward with a common velocity profile.

When looking back at Figure 5.3, the $a^*_{\epsilon_{rec}}$ varies significantly over the individual p_T bins. Due to this factor, it becomes necessary to reweigh the generated and reconstructed $\Xi(1820)$ spectra in the Monte Carlo simulation [8]. It should be noted that the yield calculations have been quoted after accounting for the mid-rapidity yield with one unit of rapidity ($dy = (0.5) - (-0.5)$).

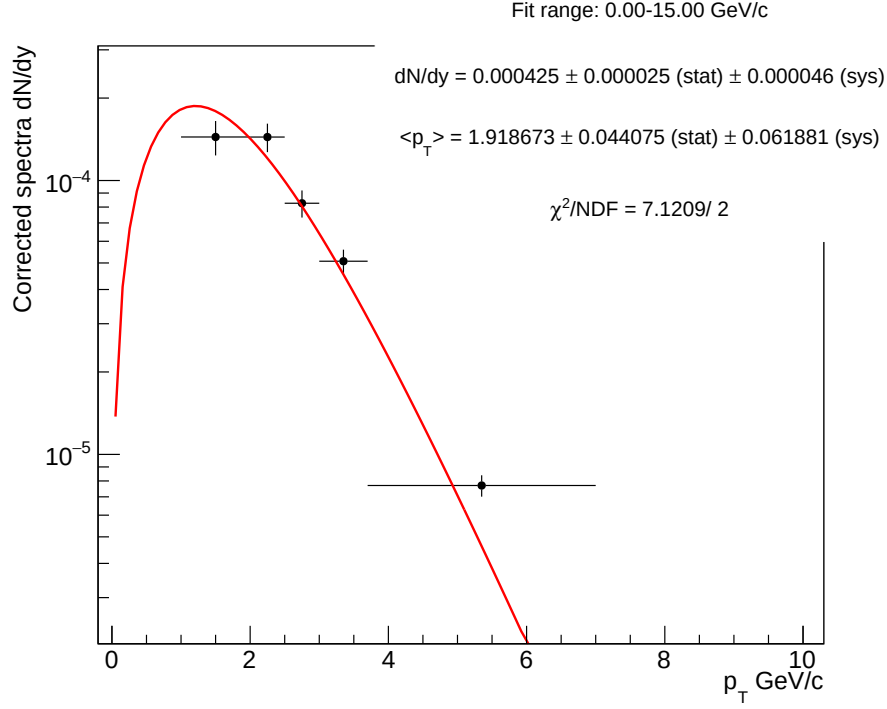


Figure 5.4: Real corrected $\Xi(1820)$ spectrum fit with Levy-Tsallis function for the minimum-bias events. No re-weighting is applied.

The first Levy-Tsallis fit of the measured $\Xi(1820)$ spectrum (without weighing) reaches its peak around $p_t \sim 1.9 \text{ GeV}/c$ as seen in Figure 5.4. It should be noted that the fit function used to determine the yield is fit from $0 < p_T < 15 \text{ GeV}/c$ due to the limit from the Monte Carlo simulation which does not generate any particles after $15 \text{ GeV}/c$. The correction factor ϵ , also known as the unweighted efficiency of $a^*_{\epsilon_{rec}}$, is observed to change rapidly over this p_T range when observed in Figure 5.3 [8]. Therefore, a re-weighting procedure on $a^*_{\epsilon_{rec}}$ is performed in order to

make sure that the generated and reconstructed $\Xi(1820)$ p_T spectrum have the same general shape as the measured $\Xi(1820)$ spectrum [8]. The process to determine the re-weighted efficiency uses an iterative procedure as described below [8].

1. The correction factor, ϵ , is calculated.
2. This ϵ is used to correct the measured $\Xi(1820)$ spectrum.
3. The corrected $\Xi(1820)$ spectrum is fitted using one of the previously mentioned functions used to determine the shape of the p_T spectra.
4. The fit is used to weight the simulated $\Xi(1820)$ spectrum. A p_T -dependent weighting is applied to the generated $\Xi(1820)$ spectrum so that it follows the fit. The same weight is applied to the reconstructed $\Xi(1820)$ spectrum.
5. The weighted ϵ is calculated and used to re-correct the $\Xi(1820)$ spectrum.
6. Steps 2-5 are repeated (with the weighted ϵ from step 5 used as the input for step 2) until the ϵ values are observed to change by $<0.1\%$ (relative) between iterations.

After all of the re-weighting is complete, the value of the integrated yield obtained from the fit functions finally account for most sources of error and can be considered complete. It is observed that three iterations are sufficient for the procedure to converge for the Levy-Tsallis, Boltzmann, and Boltzmann-Gibbs Blast-Wave functions. The final results of the reweighed functions are shown in Section 7.4.

6 Error Analysis

The error analysis of the procedure is divided into two sections. Statistical uncertainty corresponds to the procedure to estimate the uncertainty of the calculations due to statistical limits and fit procedures. Systematic uncertainty corresponds to the procedure to estimate the uncertainty of the calculations due the specific procedure used to calculate any values or fits.

6.1 Statistical Uncertainty

As expected in any rare particle search and analysis, the statistical uncertainties dominate in all results presented here. As has been mentioned several times in this discussion, the $\Xi(1820)$ is a very rare particle to investigate due to its low statistics and difficulty in reconstruction. It also has a low $a^*\varepsilon_{rec}$ at low p_T , about 0.09 at 2 GeV/ c according to Figure 5.3, where most of the $\Xi(1820)$ are produced. Thus billions of events are necessary to find the few tens of thousands of particles that are investigated.

Any parameter used in the construction of a fit, such as the mean or width of a Voigtian, has an associated statistical uncertainty that depends on the χ^2 of the fit. The applied procedure determines the uncertainty of each parameter as the amount of increase or decrease in the parameter that will increase χ^2 by a value of 1. Thus, the more the change in the parameter affects the χ^2 of the fit, the higher the statistical uncertainty for that parameter. This is different than standard methods of determining the statistical uncertainty, such as the “subsample” method, mainly because of the low statistics of the $\Xi(1820)$. While methods like the “subsample” method require large statistics and the ability to subdivide their samples into smaller subsamples to determine the effect that statistics play on the signal, the $\Xi(1820)$ sample can not be subdivided and still has a reasonable signal to be analyzed. Since the main focus of this discussion is on the mean and width of the $\Xi(1820)$, those are the only parameters that have their statistical uncertainty estimated using this procedure. The yield will be discussed shortly.

After determining which combination of variations to the fits leads to the lowest χ^2 calculation,

the corresponding fit is selected as the “default” fit, as is discussed in Section 6.2. After the “default” fit is selected, the corresponding mean and width are selected as the final mean and width from the modified Voigtian fit calculations. The corresponding statistical uncertainties are also determined for the “default” fit using the previously described procedure.

Using the same “default” fit function as is used for the mean and width calculations, the total yield for each multiplicity percentile and p_T range is calculated and the average between two different methods is used to determine the yield. This is to allow for a more accurate calculation of the yield given two methods, as well as determine the uncertainty associated with using one method over another. The two yield methods are functional integral and bin counting. Do note that the difference between these two methods is accounted for in the Section 6.2 systematic uncertainty.

The first method used for determination of the yield is the use of the integral function. The integral function calculates the integral of the Voigtian function used to fit the final invariant mass plot. While this function is only used to fit the signal in a range of $1.823 \pm 0.04 \text{ GeV}/c^2$ ($1.823 \pm 0.07 \text{ GeV}/c^2$ for Pb-Pb), the function itself can still be extrapolated to much further ranges. The integral of the Voigtian is determined from a range of the minimum mass threshold of the decay channel of interest to $2.4 \text{ GeV}/c^2$, to account for all of the potential yield that would not be seen in the plots but still contribute. This yield must be divided by the size of the bins of the invariant mass plot so that it can match the values seen in the invariant mass plot.

The process to determine the uncertainty of the integral function is thankfully very simple. After calculating the yield of the integral of the function, this value is then multiplied by the error of the scaling constant used to scale the Voigtian function to match the data points, divided by the value of the scaling constant. Thus, the error of the integral is related to the scaling constant of the Voigtian function used to determine the integral. This uncertainty is then divided by the size of the bins of the invariant mass plot so that it can match the yield.

The other method known as bin counting is, as the name implies, counting of the individual values in a bin. As stated before in Section 4.2, one of the main reasons that the invariant mass plots have their raw signal subtracted from same or mixed-event distributions and then further

subtracted by the polynomial representation of the residual background, is to create an invariant mass plot with no further influence from the background. Thus, the final invariant mass plots should only contain the signal, along with any small fluctuations in the signal that can persist throughout the process. By summing the values of the bins in a range of the signal, as discussed in Section 4.2, it is possible to determine the amount of counts (yield) for the signal.

However, because the functional integral is calculated from the mass threshold set to $2.4 \text{ GeV}/c^2$, while the bin counting method is only calculated from the chosen range of $1.823 \pm 0.04 \text{ GeV}/c^2$, there is still a certain amount of “yield” that has not been accounted for. Using the Voigtian function that was used to determine the “total” yield, two more yields are calculated using the functional integral. These contributions, which are called the “high yield” and “low yield”, correspond to the ranges of the Voigtian fit function that occupy higher and lower parts of the invariant mass plot and function respectively. The high yield is calculated in the range $1.823 + 4 \times \sigma$ ($0.04 \text{ GeV}/c^2$) to $2.4 \text{ GeV}/c^2$. The low yield is calculated in the range of the minimum mass of the decay channel for the $\Xi(1820)$, namely the mass of the $\Lambda + K$ of about $1.610 \text{ GeV}/c^2$, to the lower range of the signal $1.823 - 4 \times \sigma$ ($0.04 \text{ GeV}/c^2$).

The final yield calculation using the bin counting method is simply the sum of the actual bin counting, plus the high and low yield contributions, after being divided by the size of the invariant mass bins, $4 \text{ MeV}/c^2$, to account for the area that each bin covers.

Using these procedures, the only real difference between the two methods is the yield calculations between the range of $1.823 \pm 0.04 \text{ GeV}/c^2$. The functional integral method uses the function, while the bin counting uses the actual data points.

Determination of the statistical uncertainty for the bin counting method is similar to the method for determining the yield from bin counting. However, instead of counting the value of each bin, the uncertainty of each bin is counted.

Initially for the invariant mass plots, the uncertainty of each bin can be estimated as the square root of the value for each bin. For example, as seen in the left plot of Figure 4.4, the initial counts of the raw data point at $1.820 \text{ GeV}/c^2$ are approximately 3.2×10^5 , which leads to an error for that

data point of 565 counts. The same-event background will have approximately the same count and error for its data point at the same position since it has the same number of events used in the analysis as the raw data, hence same-event. The mixed-event background however was constructed to have 5 times the events used than those used for the raw data or same-event background, thus the counts for the data point at $1.820 \text{ GeV}/c^2$ will be about 5 times higher than the data point from the raw data. The corresponding error for the mixed-event data point would then be about $\sqrt{5}$ times greater than the raw data point if each point has the same error.

After the background distributions have been normalized to their corresponding normalization region, the errors for each data point are also normalized by the same factor. While the same-event background will have normalization around 1, with a similar effect on the error, because the mixed-event distribution has 5 times the events, the normalization will be around $1/5$. With this normalization factor, the errors for each data point are also normalized by $1/5$. With these calculations and normalization factor, the mixed-event background will normally have the lowest errors bars of any distribution, $\sqrt{5}/5 \simeq 0.45$ times the error for the raw or same-event distribution.

After subtracting the background distribution, seen in Figure 4.5, the errors from each data point of the initial and background plots are used in the following equation:

$$\sigma_{\text{IndividualBinError}} = \sqrt{\sigma_{\text{Signal}}^2 + \sigma_{\text{Background}}^2} \quad (46)$$

to determine the final error for each data point of the background subtracted plot. However, the error of the polynomial used to estimate the residual background still need to be accounted for in the total statistical error calculations for the bin counting method. Because the polynomial does not take into account the signal and only fits the residual background, the integral error function, another standard ROOT function, is used to determine the error of the polynomial function's yield. While the yield due to the residual background is removed from the calculations due to the polynomial subtraction, the error still must be noted and used for the statistical error calculations. The integral error function is similar to the integral function discussed before, however instead of

calculating the integral (yield) for a function, it calculated the error of the integral (error of yield) for a function. This error contribution is divided by the size of the invariant mass bins to match the yield calculations that have already been performed.

Following this, the uncertainty from the bin counting method is determined using the following equation:

$$\sigma_{\text{BinCountingError}} = \sqrt{\sum \sigma_{\text{IndividualBinError}}^2 + \sigma_{\text{ResidualBackgroundError}}^2 + \sigma_{\text{HighYieldError}}^2 + \sigma_{\text{LowYieldError}}^2}. \quad (47)$$

Finally, just like with the mean and width calculations, the total yield calculation is performed for the fit variation and cuts that have the lowest χ^2 . Do note that “total” yield and error in this sense still needs one more step, that will be discussed next, before is it called the “final” yield and error. The yield calculations from the functional integral and bin counting method are averaged to produce the total yield for the specified fit. The statistical error from the functional integral and bin counting method are similarly averaged to calculate the total yield statistical error.

This total statistical uncertainty for the yield, along with the value of the yield, is complete when the $a * \varepsilon_{rec}$ corrections from the Monte Carlo simulations are incorporated into the calculations. The “final” statistical uncertainty calculation of the yield is calculated using the equation:

$$\sigma_{\text{FinalYieldStat}} = \frac{\sigma_{\text{TotalYieldStat}}}{(a * \varepsilon_{rec}) \times \text{NumberOfEvents} \times dy \times \Delta p_T \times B.R.} \quad (48)$$

where dy corresponds to the range of the pseudo-rapidity used for the analysis (set to 1.0), Δp_T is the range of the p_T used for the calculation, and $B.R.$ is the Branching ratio of the decay channel of interest (set to 100 %).

6.2 Systematic Uncertainty

To calculate the systematic uncertainty, the same procedure which was done to get the results with the default cuts is performed several times over the same data by varying the possible permutations of the analysis [8]. These permutations, which will be described later, include but are not limited to

variations in the topological cuts, different methods for signal extraction, and so on. The procedure for calculating the systematic uncertainty is as follows [8]:

1. Select one set of parameters and cuts, with the lowest χ^2 , for the analysis in question as a default cut
2. Calculate the deviation of the yield, mean, or width when one parameter or cut is changed
3. A Barlow check is preformed to determine if the variation observed is due to a systematic effect instead of statistical influence [59] (See Section 6.2.1)
4. The total systematic uncertainty, considering all of the different sources, except the default cut, that pass the Barlow check, is the sum in quadrature of each source

Under the assumption that the modified Breit-Wigner function still captures all of the underlying physics processes, we can determine a systematic uncertainty using our methods combined with the Barlow checks. The total systematic uncertainty is computed using a process very similar to the normal process used to determine the standard deviation of any set of data points. Namely, each variation or permutation of the cuts, fits, and so on is considered a “value” of the observable in question. These “values” are then used to determine the standard deviation of the distribution of each observable according to the variations in the procedures.

The various permutations of the topological cuts and fits are as follows:

- To determine the effect that the DCA from the primary vertex of the reconstructed primary particle has on the analysis, the values of the DCA were changed from values that were considered “loose” to values that were considered “tight” in comparison. This is done since the initial cuts were designed to be loose and allow for more reconstructed particles to be analyzed for this discussion. Tightening these cuts then allows for fewer reconstructed particles to be analyzed, but also means that fewer misidentified particles are used for the analysis. The Λ DCA cut is tightened from a value of 0.4 cm to 0.2 cm. The K_s^0 DCA is tightened from a value of 0.3 cm to 0.15 cm.

- To determine the effect that the DCA of the daughter particles to their reconstructed mother particle has on the analysis, the cuts of the DCA were changed from values that were considered “loose” to values that were considered “tight” in comparison. Similar to the variations done for the DCA to primary vertex, the loose cuts allow for more particles to be used in the analysis with a higher chance of misidentified particles being introduced, while the tighter cuts allow for less particles to be used in the analysis with lower chance of misidentified particles be introduced. The Λ daughter DCA cut is tightened from a value of 1.0 cm to 0.3 cm. The K_s^0 daughter DCA is tightened from a value of 1.0 cm to 0.3 cm.
- The pointing angle is defined as the angle between the reconstructed particle’s direction of momentum and its displacement vector from the primary vertex. If a particle is reconstructed, but the direction the particle was coming from is not pointing toward the primary vertex, then it can be inferred that either the reconstructed particle did not come from the initial collision, or that the reconstruction of two daughter particles did not correspond to a “true” mother particle. Regardless, the smaller the angle, the closer the particle is to coming from the primary vertex. Instead of using a value of the angle, the cosine of the angle is calculated. The smaller the angle, the closer cosine is to 1, the more likely that a particle came from the primary vertex. The Cosine of the Pointing angle for the Λ is tightened from 0.99 to 0.995. The Cosine of the Pointing angle for the K_s^0 is tightened from 0.97 to 0.995.
- To determine the effect that the DCA of V^0 daughters particles to the xy plane of the primary vertex has on the analysis, the values of the DCA were tightened. However, instead of decreasing the value of the DCA, as seen in the previous DCA cuts, the DCA is increased. This is because the DCA of V^0 daughter particles in the xy plane is set to a minimum distance away from the primary vertex. By increasing the DCA, candidates will need to be further away from the xy axis of the primary vertex. Thus, increasing the DCA from 0.06 cm to 0.07 cm will further limit the number of candidates that can be identified as V^0 daughter particles.
- As stated before in section 2.1, the TPC is a very good detector used for identifying charged

particles. While the decay products of the Λ and K_s^0 have a constant and very loose 5σ cut for the proton and charged pion, the value of σ for the particles that are detected directly, namely the $K^\pm$, are a bit more flexible. Because of the fact that the $K^\pm$ can be detected directly from the TPC and TOF, the n_σ values can be defined more rigorously. Namely, the value of the n_σ for the TPC can be set to match those of the standard cuts used for $K^\pm$ analysis. Also, instead of simply tightening the values of the cuts, as was the case for most of the previous cuts, the value can also be loosened further. This allows for analysis to be performed to determine how the default, looser, and tighter cuts all compare to each other. In order to determine the effect of the TPC dE/dx selection for the charged kaon, the n_σ is varied from its initial value of 2.0 to values of 1.5 and 2.5.

- Similar to the TPC, the TOF can directly detect the $K^\pm$, without the need of an overly loose initial cut. Instead, this cut can match the cuts that are normally used for other $K^\pm$ analysis. In order to determine the effect of the TOF selection for the charged kaon, the n_σ is varied from its initial value of 3.0 to values of 2.5 and 3.5.
- The background distribution data selected for p-p, p-Pb, and Pb-Pb was selected due to the lowest value of χ^2 seen across each selected distribution. For this analysis, same-event background for p-p $\sqrt{s} = 13$ TeV data was selected, while mixed-event background for p-Pb and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data was selected. Due to the noticeable difference between both background distributions for each system, it was determined that the background distribution would not be incorporated into the systematic error analysis. For p-p data, the mixed-event background is excluded because the same-event background clearly provides a better definition of the true combinatorial background. For Pb-Pb data, the same-event background is excluded for introducing “false” peaks that do not correspond to actual particles. This is mainly seen in the fact that the differences between the two signals and fits are so large that they seem to have a difference greater than the uncertainties calculated for each distribution. Thus, each distribution is considered an “independent” distribution and is not counted as a source of

systematic error for the other distributions.

- In order to determine the effect of the normalization range on the distribution, several ranges were selected to estimate the impact of the final values calculated. Four normalization ranges are used. 1.70-1.74 GeV/ c^2 corresponds to the segment of the distribution that is toward the left of (lower than) the signal. 1.85-1.95 GeV/ c^2 corresponds to the segment of the distribution that is slightly to the right of (higher than) the signal. 2.0-2.1 GeV/ c^2 corresponds to the segment of the distribution that is further to the right of (higher than) the signal. 2.3-2.4 GeV/ c^2 corresponds to the segment of the distribution that is much farther to the right of (higher than) the signal. The normalization range that contributes to the lowest χ^2 for the corresponding system is denoted as the default. In the case of $\Xi(1820)^\mp$ in p-p collisions at $\sqrt{s} = 13$ TeV, the same-event like-sign distribution has its lowest χ^2 for the 3rd distribution (2.0-2.1 GeV/ c^2) while the mixed-event distribution has its lowest χ^2 for the 1st distribution (1.70-1.74 GeV/ c^2).
- In order to determine the effect that the polynomial's degree has on the distribution, several polynomials were selected to estimate the impact of the final value calculated. Three polynomials were used: 2nd degree (quadratic), 3rd degree (cubic), and 4th degree. The polynomial that contributed to the lowest χ^2 for the corresponding system was denoted as the default. The other polynomials are used to determine the systematic uncertainty of using a polynomial of n-th degree to estimate the residual background distribution. Do note that given the nature of polynomial fits to data, in almost all cases a higher degree polynomial would fit better to a larger amount of data points. A 1st degree polynomial perfectly matches 2 data points, a 2nd degree polynomial perfectly matches 3 data points, an n-th degree polynomial perfectly matches n+1 data points. So the higher degree polynomials are only selected as the default if the change in χ^2 is significant enough (χ^2 is lower than the lower degree polynomials by 1 or more), the higher degree polynomial matches the distribution visually better than the lower degree polynomials, or the errors for the parameters are greatly reduced.

- The range of the polynomial used to determine the residual background distribution is varied in order to determine the effect in the final values calculated. Three ranges were used: 1.7-2.0 GeV/ c^2 , 1.74-1.96 GeV/ c^2 , and 1.74-2.0 GeV/ c^2 . Due to the increased range of the signal for the Pb-Pb data, the ranges used were changed to 1.7-2.0 GeV/ c^2 , 1.73-1.93 GeV/ c^2 , and 1.73-2.0 GeV/ c^2 so as to allow for enough data points to be used for the polynomial estimation of the background subtracted plots, without interfering with the actual range of the signal and Voigtian fit. It must be noted that, while the polynomial is fit to the residual background in the range of the signal, the fit is performed so as to ignore the potential effects of signal in the range of 1.782-1.864 GeV/ c^2 (1.752-1.895 GeV/ c^2 for Pb-Pb data). The range of the polynomial that corresponds to the lowest χ^2 for the corresponding system was denoted as the default. The other ranges are used to determine the systematic uncertainty of using a polynomial with the corresponding range to estimate the residual background.
- In order to determine the effect that the σ of the Voigtian fit has on the final value calculated, the fit is performed where the σ is free to move in a range of 3 to 6 MeV/ c^2 while initially being set to a value of 4 (5 for Pb-Pb data) MeV/ c^2 . Because the σ determined from the Monte Carlo simulations is set as the default, the effect that the “free” σ has on the parameters is estimated as the systematic uncertainty. Do note that the σ will have its largest effect on the calculations of width since they are technically used for the same principle.

The sources of uncertainty that fail the Barlow checks, seen in Section 6.2.1, are finally used to determine the systematic uncertainty using the standard deviation of these sources.

6.2.1 Barlow Checks

Calculations of systematic uncertainty for yield, mean, and width normally consist of calculating the same values for different analysis parameters such as cuts, background estimation, fit range, etc..., and determining the variance between these values for each different parameter. It is therefore necessary to ensure that the variations are not caused by statistical fluctuations. It is not necessary

to account for difference between two measurements in the systematic uncertainties if it is found that they are similar to each other within the quadrature difference of their statistical uncertainties [58]. This can be investigated using the “Barlow checks” [59]. The procedure is identical to the procedure from the K^* analysis note and is explained as follows [58].

Consider two values for width of the $\Xi(1820)^\mp$, one obtained from the default setting and a second obtained with an alternative parameter (for example: different normalization range). Let us denote the width and the statistical uncertainty as Γ_{def} and σ_{def} for the default case and as Γ_{al} and σ_{al} for the alternate parameter. For each alternate parameter, we can estimate Δ/σ_{cc} , where Δ is the difference between the default and the alternative measurement ($\Delta = \Gamma_{def} - \Gamma_{al}$) and σ_{cc} is the difference in the quadrature of the statistical uncertainties of the measurements ($\sigma_{cc} = \sqrt{|\sigma_{def}^2 - \sigma_{al}^2|}$) [58]. Assuming that the two measurements are consistent, it is expected that the distribution for all alternate cuts, fits, etc. involved between these two measurements of Δ/σ_{cc} would have a mean near 0, a standard deviation near 1, and that 68% (95%) of the entries would lie within $|\Delta/\sigma_{cc}| < 1$ ($|\Delta/\sigma_{cc}| < 2$) [58]. For this analysis, we do not consider a cut or variation to the fit as a systematic source of uncertainty if at least 3 of the following 4 criteria are satisfied:

1. $|\Delta/\sigma_{cc}| < 0.12$
2. Standard deviation < 1.3
3. Fraction of entries within ± 1 ($I_1 > 0.55$)
4. Fraction of entries within ± 2 ($I_2 > 0.75$)

The selection of these criteria comes from the “ideal” Gaussian distribution. Assume that an “ideal” Gaussian distribution was formed with a mean of 0 and a standard deviation of 1. If another Gaussian distribution was formed that had a standard deviation of 1.3, then I_1 for this distribution would be 80 % of the “ideal” value for I_1 determined from the “ideal” Gaussian. $I_1 = 0.55$ is 80 % of the “ideal” value for I_1 . $I_2 = 0.75$ is 80 % of the “ideal” value of I_2 .

This process allows for the analysis of the systematic uncertainties to exclude sources that will

not greatly contribute to the uncertainty and are consistent with statistical fluctuations. Caution must be used when using the Barlow check, as passing the Barlow check means that the source is not a source of systematic uncertainty, while failing the Barlow check means that the source is included in the systematic uncertainty.

6.3 Summary of Systematic Uncertainty Contributions

The following tables are a list of the sources of systematic uncertainty for the yield, mean, and width calculations. These values do include the Barlow checks described in Section 6.2.1. The systematic uncertainties listed in Tables 5, 6, and 7 are calculated using the fractional uncertainty, which is each specific uncertainty divided by the value of the parameter in question. So, if the width had a value of $24 \text{ MeV}/c^2$ with a systematic uncertainty of $\pm 3 \text{ MeV}/c^2$, the fractional uncertainty would be $3/24 = 12.5 \%$.

Do note that the data for Table 5 is for the systematic uncertainties for only the p-p collisions at $\sqrt{s} = 13 \text{ TeV}$ for $\Lambda+K^\mp$ for the 0-100 % multiplicity percentile and transverse momentum $1.0 < p_T < 20.0 \text{ GeV}/c$ for same-event background. The data for Table 6 is for the systematic uncertainties for only the p-Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ for $\Lambda+K^\mp$ for the 0-100 % multiplicity percentile and transverse momentum $1.0 < p_T < 20.0 \text{ GeV}/c$ for mixed-event background. The data for Table 7 is for the systematic uncertainties for only the Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ for $\Lambda+K^\mp$ for the “central” kCent triggered (0-10 %) multiplicity percentile and transverse momentum $1.0 < p_T < 20.0 \text{ GeV}/c$ for mixed-event background. Other systems, energies, multiplicities, momenta, and so on would have different values of systematic uncertainty and contributions. Furthermore, the segments that are labeled with X correspond to variations that pass the Barlow check and thus not considered sources of systematic uncertainty. Segments labeled with XX correspond to sources of systematic uncertainty that do not apply to that corresponding value’s systematic uncertainty. This is only used for two cases: calculation of the systematic uncertainty that arises from the use of using both the integral of the fit function vs. the bin counting method

for determining the yield, or inclusion of the error for $a^*\varepsilon_{rec}$ incorporation into the final yield calculations. Since this has no effect on the mean and width calculations, it is not included in their calculations. Segments labeled with XXX correspond to sources of systematic uncertainty that are so great that they can be concluded to come from insufficient fits, and must not be included into the systematic uncertainty. Also, it must be noted that the total systematic uncertainty is not the sum of each systematic uncertainty listed, rather it is the sum in quadrature of each source. The final line corresponds to the statistical uncertainty which is calculated from Section 6.1.

Source of uncertainty	Yield(%)	Mean(%)	Width(%)
Normalization Range	0.60	X	X
Free σ of Voigtian	0.04	0.001	3.16
Polynomial degree	0.33	X	0.44
Fit Range	2.13	0.003	3.00
Λ Daughter DCA to each other	6.08	0.005	5.23
Λ DCA to PV	2.12	0.020	X
Pointing Angle $\cos\theta$	0.26	X	0.79
DCA of V^0 daughters to PV xy axis	0.50	X	X
$K^\mp$ TPC n_σ	3.41	0.013	2.87
$K^\mp$ TOF n_σ	0.95	X	X
Integral vs. Bin counting	0.03	XX	XX
$a^*\varepsilon_{rec}$ uncertainty	1.04	XX	XX
Total systematic uncertainty	7.80	0.025	7.44
Total statistical uncertainty	6.75	0.044	12.68

Table 5: Systematic and statistical uncertainty of $\Lambda+K^\mp$ for p-p $\sqrt{s} = 13$ TeV with 0-100 % multiplicity percentile and transverse momentum $1.0 < p_T < 20.0$ GeV/ c with same-event background.

From Table 5, several features become apparent as each source of systematic uncertainty is investigated. For example, some of the largest sources of systematic uncertainty come from the topological cuts such as the DCA of the Λ daughter particles to each other. The fit range has a similar effect on the systematic uncertainty as the width for the free σ , even though the effect that the free σ has on the width is almost 1 to 1, in that if the sigma of the Voigtian fit increases by a value of 1 MeV/ c^2 , then the width will decrease by a value of 1 MeV/ c^2 .

However, caution should be used when using the tables such as Table 5 since some values may at first appear concerning. For example, the mean would appear to be very stable according to the

systematic analysis and the low uncertainty. Given the value of the mean however, which comes to about $1822.3 \text{ MeV}/c^2$ even a hypothetical $24 \text{ MeV}/c^2$ uncertainty (which would be of the same magnitude as the width of the Voigtian fit) would only have a $24/1822.3 = 1.32 \%$ systematic uncertainty. Thus while 0.025% is certainly low, it would correspond to a systematic uncertainty of about $\pm 0.4 \text{ MeV}/c^2$. Stable yes, but do not forget that most of the “variance” in the fits would normally be seen in the width of the fits, not the mean unless it is drastic enough.

Source of uncertainty	Yield(%)	Mean(%)	Width(%)
Normalization Range	2.33	X	2.1
Free σ of Voigtian	0.38	X	2.1
Polynomial degree	0.44	X	0.4
Fit Range	5.49	0.010	5.4
Λ Daughter DCA to each other	6.68	X	X
Λ DCA to PV	5.99	0.019	2.3
Pointing Angle $\cos \theta$	0.68	X	X
DCA of V^0 daughters to PV xy axis	0.004	X	X
$K^\pm$ TPC n_σ	3.59	0.015	0.9
$K^\pm$ TOF n_σ	X	0.008	X
Integral vs. Bin counting	0.56	XX	XX
$a^* \epsilon_{rec}$ uncertainty	2.09	XX	XX
Total systematic uncertainty	11.61	0.027	6.6
Total statistical uncertainty	8.33	0.062	13.74

Table 6: Systematic and statistical uncertainty of $\Lambda+K^\pm$ for p-Pb $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ with 0-100 % multiplicity percentile and transverse momentum $1.0 < p_T < 20.0 \text{ GeV}/c$ with mixed-event background.

From Table 6, several features become apparent as each source of systematic uncertainty is investigated. While the mean and width systematic uncertainties are similar to the p-p data from Table 5, the contributions are slightly different. Furthermore, the systematic uncertainty for the yield is higher than for Table 5. This is mostly due to the fact that the variations and procedures for the fits and topological cuts for p-p $\sqrt{s} = 13 \text{ TeV}$ data simply do not work as well for p-Pb $\sqrt{s_{NN}} = 5.02$ data.

One important point to be addressed is the use of the Voigtian’s σ for the p-Pb data. As discussed previously, Monte Carlo simulations were constructed for p-p and Pb-Pb data, but not

for p-Pb data. There was simply not enough time or resources to perform the simulation. However, given the fact that no corrected yield calculations were needed for p-Pb data, and the fact that the resolution (σ) of the detectors from the simulations of p-p and Pb-Pb data are very similar (5.1 to 5.0 MeV/ c^2), the p-p Monte Carlo data was used to determine the σ of resolution of the detectors. It should be enough to at least get a good estimate of the mean and width for p-Pb data. Furthermore, because σ and the width of the Voigtian have an almost 1 to 1 relation to each other, if the resolution from the Pb-Pb Monte Carlo simulation is used instead, it would only decrease the σ by 0.1 MeV/ c^2 , while increasing the width by 0.1 MeV/ c^2 .

Source of uncertainty	Yield(%)	Mean(%)	Width(%)
Normalization Range	XXX	XXX	XXX
Free σ of Voigtian	0.12	X	0.68
Polynomial degree	0.98	0.007	0.64
Fit Range	5.43	0.052	2.21
Λ Daughter DCA to each other	4.52	0.026	4.37
Λ DCA to PV	6.95	0.080	1.74
Pointing Angle $\cos \theta$	0.18	X	4.56
DCA of V^0 daughters to PV xy axis	2.35	0.020	6.84
$K^\mp$ TPC n_σ	6.25	0.064	7.42
$K^\mp$ TOF n_σ	3.67	X	8.91
Integral vs. Bin counting	1.24	XX	XX
$a^* \varepsilon_{rec}$ uncertainty	3.13	XX	XX
Total systematic uncertainty	12.54	0.120	15.16
Total statistical uncertainty	13.57	0.184	19.36

Table 7: Systematic and statistical uncertainty of $\Lambda+K^\mp$ for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV with kCent triggered data (0-10 % multiplicity percentile) and transverse momentum $1.0 < p_T < 20.0$ GeV/ c with mixed-event background.

From Table 7, several features become apparent as each source of systematic uncertainty is investigated. The most notable feature is the larger uncertainties of the yield, mean, and width when compared to Tables 5 and 6. This is most likely caused by the fact that the variations and procedures used to determine the systematic uncertainty that have been optimized for p-p collisions simply do not perform as well when applied to p-Pb or Pb-Pb collisions. To reach similar systematic uncertainties across all system sizes, different variations and procedures may need to be

implemented for each system size.

Another important feature of Table 7 is that the normalization range contains fits that are bad enough to not be considered in the systematic uncertainties. These insufficient fits mostly come from the inability of other normalization ranges, aside from the default normalization range, to accurately describe the combinatorial distribution. When non-default normalization ranges are used for Pb-Pb data, the polynomial fits appear to be greatly distorted, to the point that they can not be used as a true estimation of the residual background. These insufficient fits from other normalization ranges are not included into the systematic uncertainties.

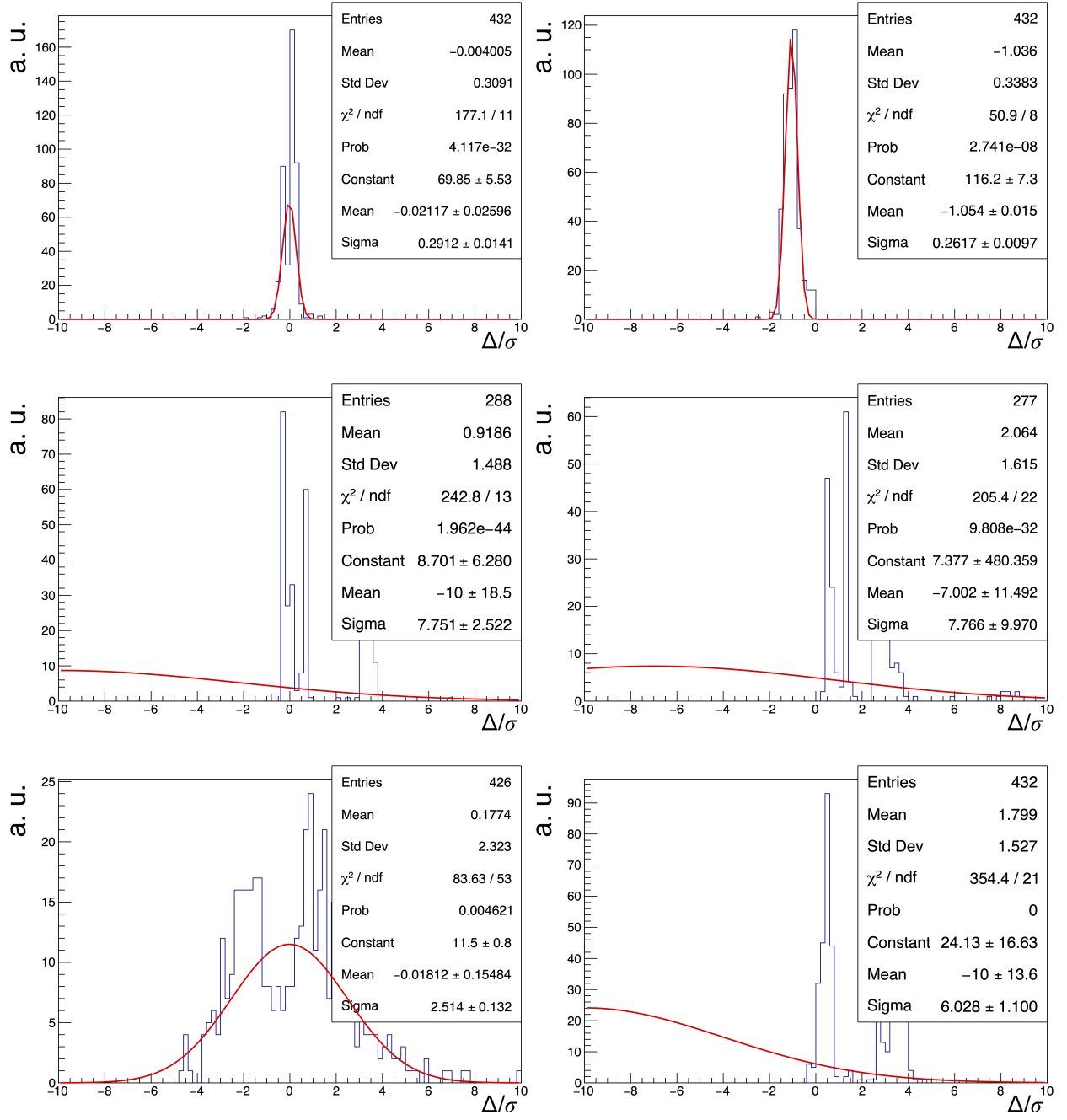


Figure 6.1: Barlow check plots for the mean and width of the Voigtian fits for p-p collisions at $\sqrt{s} = 13$ TeV with same-event background used. $\Lambda + K^\mp$, 0-100 % multiplicity percentile, $1.0 < p_T < 20.0$ GeV/c. Left plots: Barlow check for mean. Right plots: Barlow check for width. Top plots: Barlow checks for changes in the polynomial degree. Middle plots: Barlow checks for changes to the TPC PID cuts. Bottom plots: Barlow checks for changes to the range of the fit.

Examination of Barlow plots such as Figure 6.1 show several Barlow check plots for several variations and permutations. All plots were fitted with a Gaussian function to determine the mean and σ of the plots and if they correspond to the Barlow checks listed before. The integrals I_1 and I_2 are not shown here, but are still calculated.

The left plots of Figure 6.1 shows the Barlow check plots for the mean of the Voigtian fits for p-p collisions at $\sqrt{s} = 13$ TeV with same-event background subtraction used on $\Lambda K^\mp$ reconstructed data from 0-100 % multiplicity percentile and transverse momentum $1.0 < p_T < 20.0$ GeV/c. The right plots of Figure 6.1 also show the same Barlow checks for the same data set as previously discussed however, instead of the Barlow checks being performed on the mean of the Voigtian fits, these checks are performed on the width of the Voigtian fits.

The top plots of Figure 6.1 shows the Barlow check when the polynomial degree changed. As can be seen from the mean and σ of the left plot, both of these values match the Barlow check criteria listed before. Furthermore, just by observing the plot, almost all entries are within ± 1 of the center of the distribution. Thus, it can be concluded based off the Barlow checks that changing the polynomial degree has almost no effect on the systematic uncertainty of the mean, and as such can be ignored as a source of systematic uncertainty. The right plot shares several similarities with the left plot. However, while the peak from the left plot is centered close to 0 as can be seen from the mean of the gaussian fit, the peak from the right plot is toward the left, centered around a value of -1.0. Even though the σ is small enough to pass the Barlow check, the mean and fraction of entries within the ranges ± 1 and ± 2 do not satisfy the Barlow check criteria. Thus, for the width, changing the polynomial degree does have enough of an effect on the systematic uncertainty that it must be included as a source of systematic error.

The middle plots of Figure 6.1 shows the Barlow checks when the TPC PID cuts are changed. The Gaussian fit for the left and right plots just can not match the distribution very well, which is not surprising given the three peaks that would need to be fit by a function that only has one peak. Even if the Gaussian could fit the distribution to better describe its shape, it is clear that there are enough data points past the range of ± 1 and ± 2 that the fractions of entries would not pass the

3rd or 4th Barlow check. Thus, it can be concluded by the Barlow checks that changing the TPC PID cuts does have a noticeable effect on the systematic uncertainty of the mean and width, and as such must be included as a source of systematic uncertainty.

The bottom plots of Figure 6.1 show the Barlow checks when the range of the fit function was changed. These distributions, when compared to the top plots which look like a singular peak or the middle plots that have several peaks of various heights and widths, is much closer to what would be considered a Gaussian distribution. However, even though the bottom left distribution is very close to a Gaussian, the fit of a Gaussian to the distribution shows that the σ are slightly higher than the criteria for the Barlow checks listed before. The right plot shows a much more unreliable fit, and still the fraction of entries above ± 1 and ± 2 would not pass the Barlow checks. Thus it can be concluded, based on the Barlow checks, that changing the fit range has an effect on the systematic uncertainty of the mean and width, and as such must be included as a source of systematic uncertainty.

These Barlow check plots are created for the yield, mean, and width for every variation of the fits and procedure. Even though one source may not be considered a source of systematic uncertainty for one value, that same source may be a large source of systematic uncertainty for another value.

7 Results

The mass, width, and yield of the $\Xi(1820)$ have been measured in p-p collisions at $\sqrt{s} = 13$ TeV, p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. From these measurements, the mean and width of the $\Xi(1820)$ for different system sizes have been observed and analyzed. Integrated particle yield-ratios for p-p $\sqrt{s} = 13$ TeV have also been obtained. However, integrated particle yield-ratios for p-Pb and Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV were not calculated due to insufficient statistics.

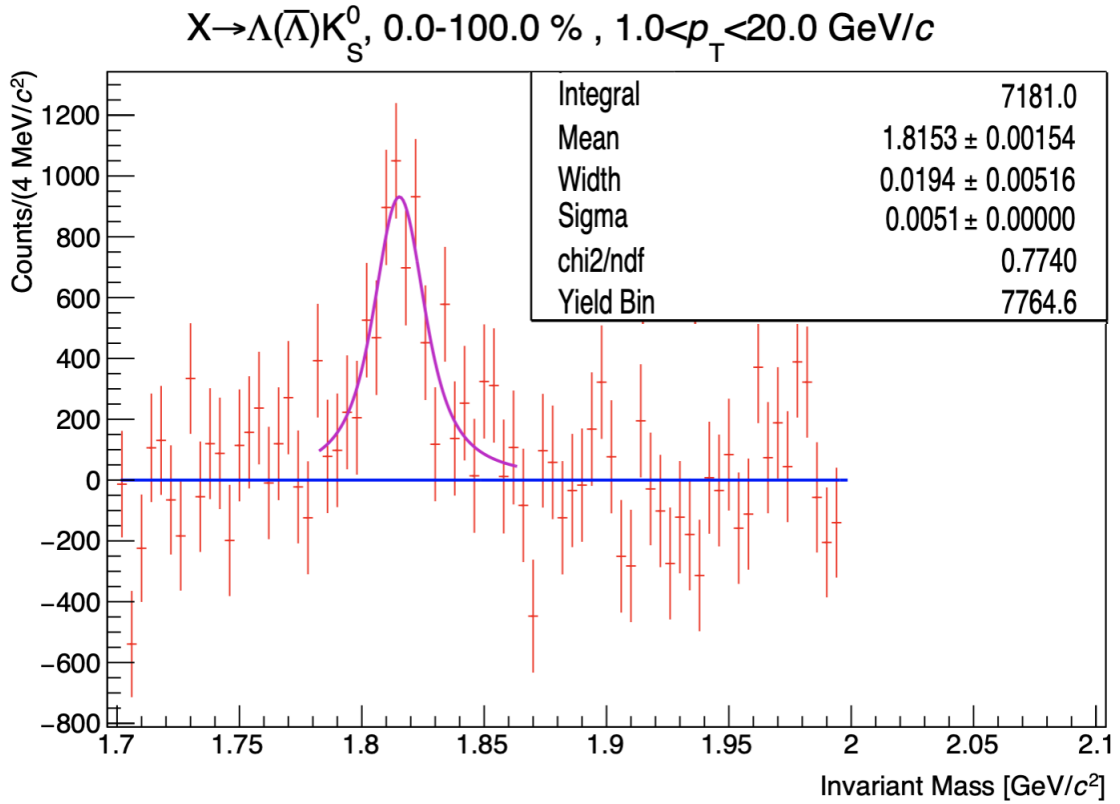


Figure 7.1: $\Xi(1820)^0$ for p-p collisions at $\sqrt{s} = 13$ TeV for 0-100 % multiplicity percentile and $1 < p_T < 20$ GeV/c.

Do note that Figure 7.1 is the only figure of the $\Xi(1820)^0$ state. Due to the low statistics of this “best case” minimum bias plot, further subdivision of the $\Xi(1820)^0$ into different multiplicity classes or p_T bins was not possible. Nevertheless, the $\Xi(1820)^0$ is seen and is fit very well. This

one measurement may not be useful in determining if chiral symmetry is restored, but it is useful in determining the differences between the $\Xi(1820)^0$ and $\Xi(1820)^\mp$ charged states that have been considered the same until now.

The following figures are the final invariant mass plots for p-p $\sqrt{s} = 13$ TeV, p-Pb $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data for the charged $\Xi(1820)$. Both charge states, i.e. particle and anti-particle, have been added. All of the final mean, width and statistical uncertainty values can be observed in these plots. These plots are not efficiency corrected.

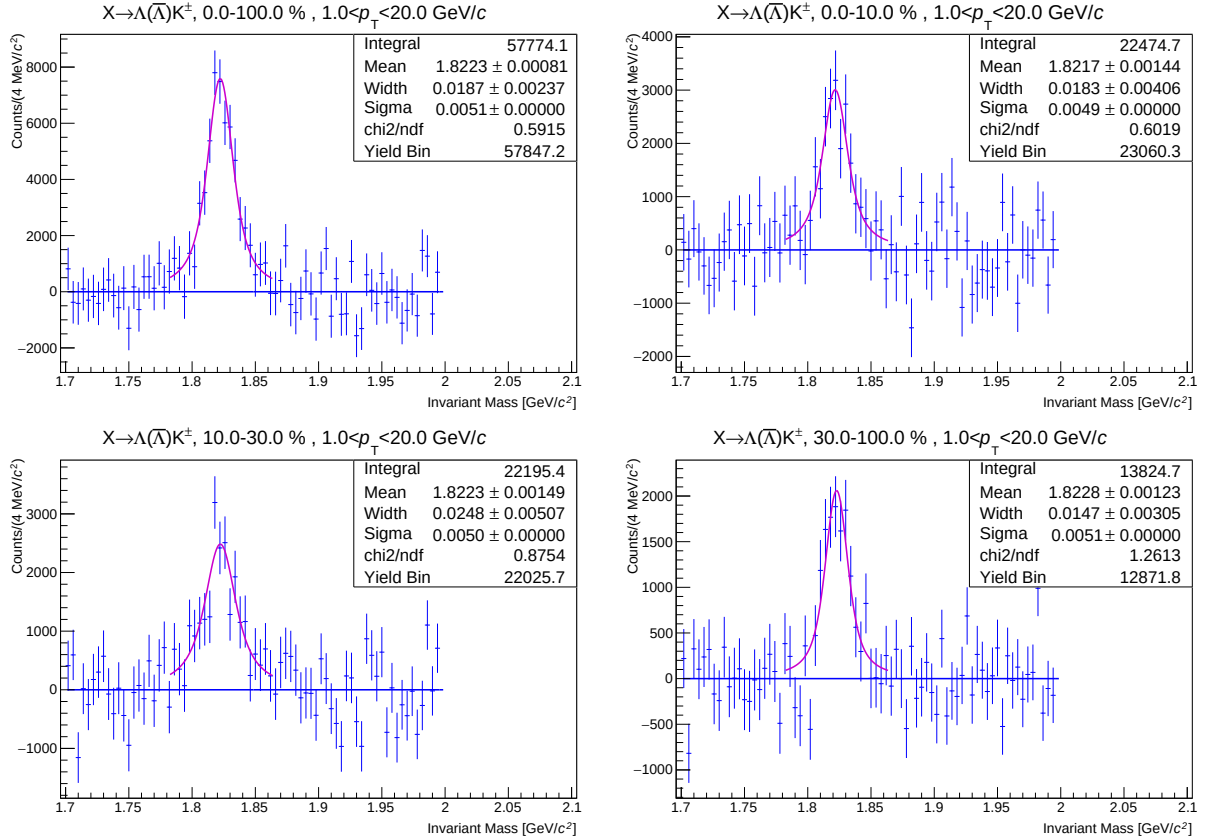


Figure 7.2: Invariant mass plots of $\Lambda(\bar{\Lambda})K^\pm$ for various multiplicity percentiles using p-p collisions at $\sqrt{s} = 13$ TeV, $1.0 < p_T < 20.0$ GeV/c. Background subtraction of same-event background and residual polynomial background implemented.

The plots of Figure 7.2 show the final invariant mass plots for p-p collisions at $\sqrt{s} = 13$ TeV, $1.0 < p_T < 20.0$ GeV/c at various multiplicity percentiles. Each plot shows a clear and noticeable signal that is fit with a modified Voigtian function to determine the mean, width, and functional

yield calculations.

The top left plot of Figure 7.2, for 0-100 % multiplicity percentile, shows the $\Xi(1820)$ signal that will be the basis of comparison to other invariant mass plots.

For example, when comparing the top left plot of Figure 7.2 for $\Xi(1820)^\mp \rightarrow \Lambda + K^\mp$ to the plot in Figure 7.1 for $\Xi(1820)^0 \rightarrow \Lambda + K_S^0$, the most striking feature is the difference in statistics between the two minimum bias plots. Figure 7.1 only has a raw yield of about 7-8000 counts, with the highest data point around 1000 counts with an error of about 200 counts. The top left plot of Figure 7.2 has a raw yield of about 55-60000 counts, with the highest data point around 8000 counts with an error of about 800 counts. The $\Xi(1820)^\mp \rightarrow \Lambda + K^\mp$ plot has about 8 times the statistics of the $\Xi(1820)^0 \rightarrow \Lambda + K_S^0$ plot. While at first this may appear confusing, under the assumption that an equal number of $\Xi(1820)^-$, $\bar{\Xi}(1820)^+$, and $\Xi(1820)^0$ should be produced in the initial collision, after accounting for the fact that the $\Xi(1820)^0$ requires four particles to be reconstructed into two V^0 particles, and the efficiency of such a reconstruction, it is no wonder that $\Xi(1820)^0$ can not reach the same statistics as the $\Xi(1820)^\mp$. As opposed to having the $\Xi(1820)^0$ with such low statistics that no signal can be seen, having a $\Xi(1820)^0$ signal at all is truly remarkable.

Another striking feature is the difference in masses between the two decay channels. The $\Xi(1820)^0 \rightarrow \Lambda + K_S^0$ has a mass of $1815.3 \pm 1.5 \text{ MeV}/c^2$, while the $\Xi(1820)^\mp \rightarrow \Lambda + K^\mp$ has a mass of $1822.3 \pm 0.8 \text{ MeV}/c^2$. As stated in Section 3.3, the $\Xi(1820)^0$ and $\Xi(1820)^\mp$ should have different masses and widths. With a difference in mass of $7.0 \text{ MeV}/c^2$ and only considering for now the statistical uncertainties of $1.7 \text{ MeV}/c^2$, this is strong evidence for mass and width difference between the neutral and charged states of the $\Xi(1820)$. The only limitation for any conclusive statement is the low statistics of the $\Xi(1820)^0$. Higher statistics for the $\Xi(1820)^0$ would lower the statistical uncertainty so that confidence in a conclusion about the difference in mass would be boosted.

The 0-10 %, 10-30 %, and 30-100% plots from Figure 7.2 each show roughly equivalent values for yield, mean, width, and statistical uncertainty. While some plots have some of their values noticeably different from other plots, such as the 10-30 % width of $24.8 \text{ MeV}/c^2$, they still show a

$\Xi(1820)$ signal that is fit and contains higher statistics than Figure 7.1. Interestingly, despite the 0-10 % and 10-30 % having very similar yields, they have slightly different widths and uncertainties, though still within error of each other. This may be due to the 10-30 % distribution not looking “Voigtian” enough, when compared to the 0-100 % plot, but may simply be due to statistical fluctuations.

One interesting thing to note is that previous analysis into some subranges of the total p-p data set, such as using only p-p $\sqrt{s} = 13$ TeV data from ALICE collected from 2016-2017 while 2018 data was still being collected, showed a similar trend of the 10-30 % distribution having a larger width than other multiplicity distributions for p-p data. This was concluded as simply due to low statistics at the time, as when 2018 data from ALICE was collected and added to the total dataset, the statistics and yield of the 10-30 % almost matched the yield of the 0-100 % distribution for 2016-2017 data.

Regardless, each distribution shown in Figure 7.2 is used to determine if the mass or width of the $\Xi(1820)$ does change with multiplicity density.

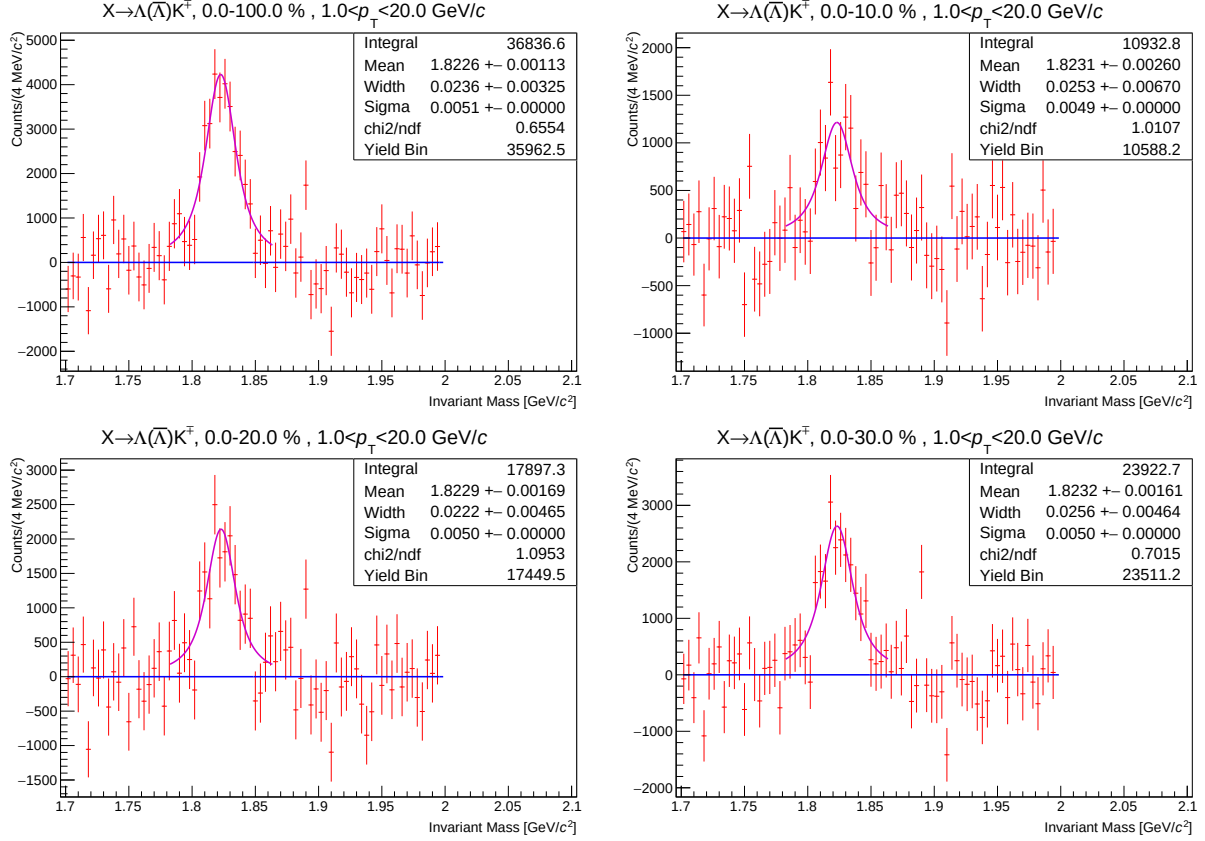


Figure 7.3: Invariant mass plots of $\Lambda(\bar{\Lambda})K^\pm$ for various multiplicity percentiles for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV data. Background subtraction of mixed-event background and residual polynomial background implemented.

The plots of Figure 7.3 show the final invariant mass plots for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, $1.0 < p_T < 20.0$ GeV/c at various multiplicity percentiles.

The top left plot of Figure 7.3, for 0-100 % multiplicity percentile, shows a clear and noticeable signal that is fit with a modified Voigtian function to determine the mean, width, and functional yield calculations. The other multiplicity classes however need to be discussed further.

The 0-10 % and 0-20 % multiplicity percentile plots, top right and bottom left plots of Figure 7.3 respectively, are not considered for the p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV analysis mostly due to the low statistics of the signals. The 0-10 % signal is simply not large enough when compared to the background statistical variations to conclusively determine that a signal will be captured by the Voigtian fit. The 0-20 % signal does not seem to have a clearly curved structure that would be

represented by the Voigtian fit.

The 0-30 % multiplicity percentile, bottom right plot of Figure 7.3, does seem to have notable signal and good representation with a Voigtian fit. The signal is seen above the background, the data seems to be in a curved shape that could be represented with a Voigtian fit, and the fit has a χ^2/ndf below 1. This plot seems to satisfy enough conditions and criteria to be used for this analysis.

The use of the p-Pb data in comparison to the p-p data would at first seem to indicate that chiral symmetry is partially restored. The width of the 0-100 % p-Pb plot from Figure 7.3 is higher than the width of the 0-100 % p-p plot from Figure, 7.2, and the 0-30 % p-Pb plot is higher still. However, given the statistical uncertainties of these plots and the fact that chiral symmetry restoration should be more pronounced in Pb-Pb collisions, no conclusive statement can be issued based only off these plots.

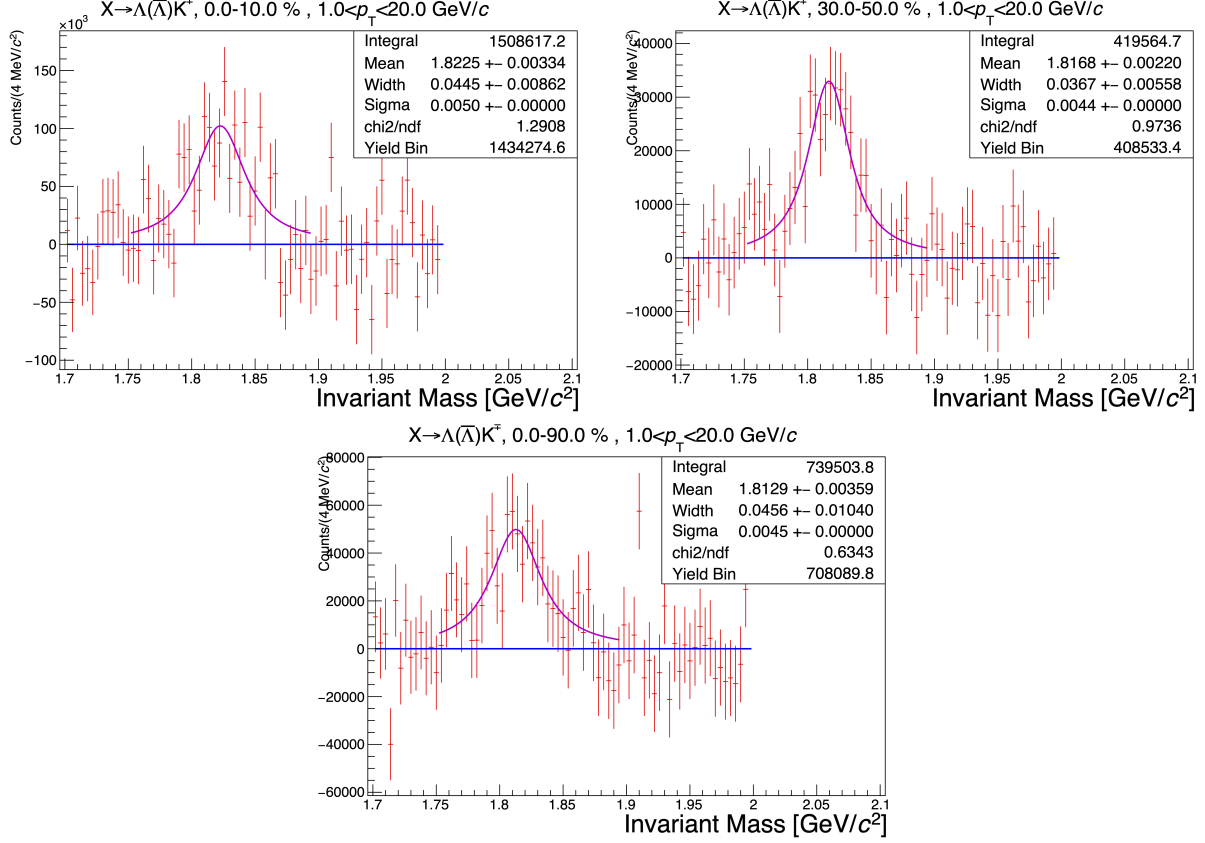


Figure 7.4: Invariant mass plots of $\Lambda(\bar{\Lambda})K^\pm$ for various triggers for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data. Background subtraction of mixed-event background and residual polynomial background implemented. Top left: “Central” kCent trigger (0-10 %). Top right: “Semi-Cent” kSemiCent trigger (30-50 %). Bottom: “Minimum Bias” kINT7 (0-90 %).

Figure 7.4 shows the final invariant mass plots for Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, $1.0 < p_T < 20.0$ GeV/c at various centrality triggers.

The bottom plot, for “minimum-bias” kINT7 (0-90 %) triggered data, seems to indicate that there is a signal, however the statistical fluctuations and large uncertainties limits the accuracy of the modified Voigtian fit. In terms of χ^2/ndf , the fit looks very good. However, the same fit has the highest statistical uncertainties of the Pb-Pb data show in Figure 7.4. While a signal is seen and can be considered, the large uncertainties most likely come from a larger background contribution.

The “central” kCent (0-10 %) and “semi-central” kSemiCent (30-50 %) triggered data, top left and top right plots of Figure 7.4 respectively, also shows noticeable signals for the $\Xi(1820)$.

While the kSemiCent triggered data appears to have a notable signal, the kCent triggered data is a bit different. The kCent triggered data contains much more statistical fluctuation than the kSemiCent or kINT7 triggered data. While the kCent plot does show a signal, the higher statistical fluctuations correspond to higher statistical uncertainties and a higher χ^2/ndf . Regardless, all three Pb-Pb triggered plots contain a signal that can be used to determine the properties of the $\Xi(1820)$ and see if chiral symmetry is indeed restored.

One very important feature of the Pb-Pb plots is that they seem to hint at two of the signatures of chiral symmetry restoration. The kINT7 and kSemiCent triggered data seem to indicate that a mean lower than 1820 MeV/ c^2 is observed, but that is not the case in kCent triggered data, where the largest chiral symmetry effect is expected to be seen. While the systematic uncertainty still needs to be calculated to be certain and make a definitive statement, it should be noted that none of the previously shown plots for p-p or p-Pb data had a width for the Voigtian fits above 30 MeV/ c^2 , while the Pb-Pb plots have three widths above 35 MeV/ c^2 . The statistical uncertainty will limit any conclusive statement, but this is very encouraging for confirming chiral symmetry restoration.

7.1 Mass

The mass of the $\Xi(1820)$ for p-p, p-Pb, and Pb-Pb collisions is shown in Figure 7.5. The statistical and systematic uncertainties are shown as lines and capped boxes error bars respectively on the plot. The first, third, and fourth blue data points (30-100 %, 10-30 %, and 0-10 % respectively) for Figure 7.5 are subsets of the second data point (0-100 %) for p-p $\sqrt{s} = 13$ TeV data, thus the minimum bias (0-100 %) p-p data point is chosen to make the final comparison to the Pb-Pb data. The mass from the kINT7 triggered Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data is not included in Figure 7.5 due to the larger statistical uncertainties when compared to other Pb-Pb triggered data sets.

Individual data points for the mass, statistical uncertainty, and systematic uncertainty for $\Xi(1820)^\mp$ are shown in Table 8. $\Xi(1820)^0$ for p-p $\sqrt{s} = 13$ TeV at 0-100 % is included at the bottom of the table only so that it can be compared to $\Xi(1820)^\mp$ for p-p $\sqrt{s} = 13$ TeV at 0-100 %.

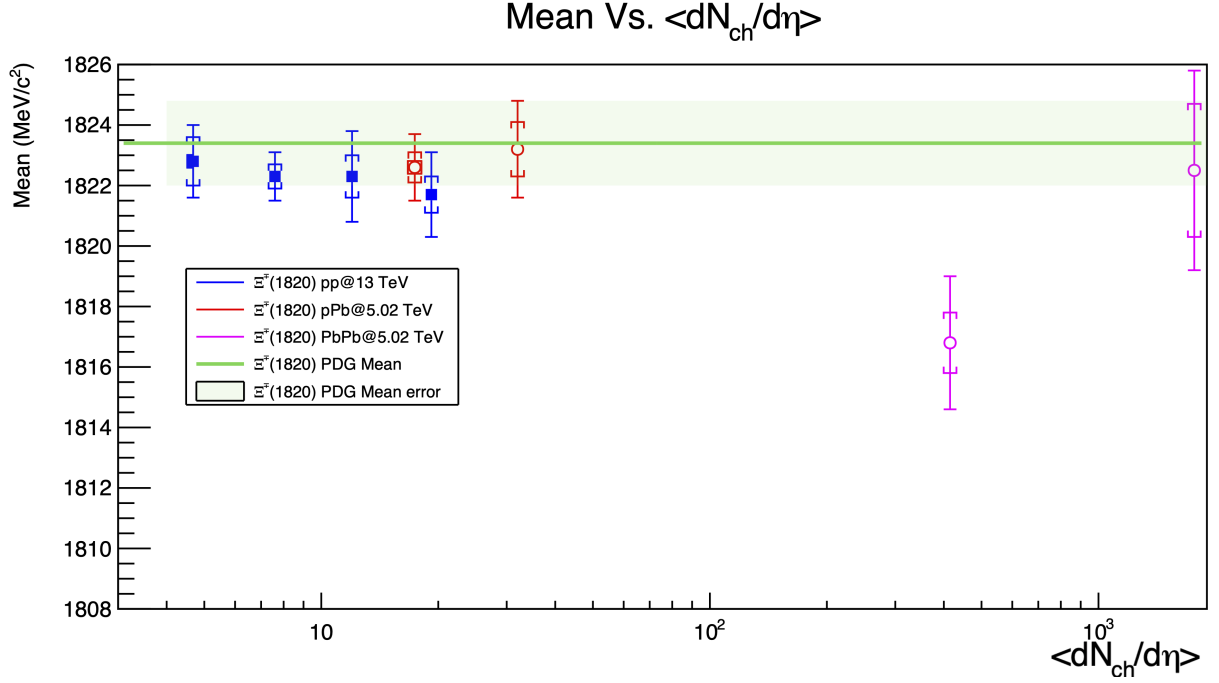


Figure 7.5: Mean of the Voigtian fits of $\Xi(1820)^\mp$ for different system sizes and multiplicity percentile. Statistical uncertainties are shown as lines and systematic uncertainties are shown as capex boxes.

The momentum range of p-p $\sqrt{s} = 13$ TeV, p-Pb at $\sqrt{s_{NN}} = 5.02$, and Pb-Pb at $\sqrt{s_{NN}} = 5.02$ data corresponds to $1 < p_T < 20$ GeV/c. The distribution used to estimate the background of the p-p $\sqrt{s} = 13$ TeV data corresponds to same-event distribution, while the distribution used to estimate the background of the p-Pb at $\sqrt{s_{NN}} = 5.02$ and Pb-Pb at $\sqrt{s_{NN}} = 5.02$ TeV data corresponds to mixed-event distribution.

Two major choices that need to be discussed are the p_T binning and the multiplicity classes for each system.

The choice to not include the range of $0.0 < p_T < 1.0$ GeV/c comes from the fact that there was no $\Xi(1820)$ signal for this p_T range due to the very low $a^* \varepsilon_{rec}$. Invariant mass plots need to show a reasonable peak that can be analyzed, so the fact that $0.0 < p_T < 1.0$ GeV/c did not show a signal while other p_T bins did, see Section 7.3, means that $0.0 < p_T < 1.0$ GeV/c can be ignored and only $1.0 < p_T < 20.0$ GeV/c need to be considered.

Collision Sys, Multi. Percentile (%)	$\langle dN_{ch}/d\eta \rangle \pm \text{Sys}$	Mean $\pm \text{Stat} \pm \text{Sys}$ (MeV/ c^2)
p-p, 30-100 %	4.68 ± 0.16	$1822.8 \pm 1.2 \pm 0.8$
p-p, 0-100 %	7.60 ± 0.14	$1822.3 \pm 0.8 \pm 0.4$
p-p, 10-30 %	12.0 ± 0.3	$1822.3 \pm 1.5 \pm 0.7$
p-Pb, 0-100 %	17.4 ± 0.7	$1822.6 \pm 1.1 \pm 0.5$
p-p, 0-10 %	19.2 ± 0.5	$1821.7 \pm 1.4 \pm 0.6$
p-Pb, 0-30 %	32.0 ± 0.8	$1823.2 \pm 1.6 \pm 0.9$
Pb-Pb, 30-50 %	415 ± 10	$1816.8 \pm 2.2 \pm 1.0$
Pb-Pb, 0-90 %	545 ± 6	$1812.9 \pm 3.6 \pm 7.6$
Pb-Pb, 0-10 %	1765 ± 36	$1822.5 \pm 3.3 \pm 2.2$
p-p, 0-100 % $\Xi(1820)^0$	7.60 ± 0.14	$1815.3 \pm 1.5 \pm 1.1$

Table 8: All $\Xi(1820)^\mp$ mean calculations for p-p $\sqrt{s} = 13$ TeV, p-Pb $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data. Statistical and systematic uncertainties are included. $\Xi(1820)^0$ mean for p-p $\sqrt{s} = 13$ TeV at 0-100 % is included [40] [41] [42] [43].

The choice to use the multiplicity percentile classes of 0-100 % and 0-30 % for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV was again limited by the statistics of the data sets. As discussed in the previous section, while the 0-100 % classes did contain enough statistics for a $\Xi(1820)$ signal to be analyzed, the same could not be said for the other multiplicity classes. 0-10 % did not contain enough statistics or a sizable $\Xi(1820)$ signal to be analyzed with a high enough degree of certainty that could be compared to other peaks. 0-20 % did show an increased amount of statistics and a better $\Xi(1820)$ signal, but χ^2 analysis showed that fits for the invariant mass plots were not up to standards that would be recommended for analysis of such a limited signal such as chiral symmetry restoration. 0-30 % contained enough statistics and a well fit peak that was used in the analysis. 30-100 % was considered at one point to be included in the analysis, however due to the fact that 0-30 % barely contained enough statistics and a good signal, and the fact that it contained about 2/3 of the total yield for the system, it was concluded that the 30-100 % for p-Pb $\sqrt{s_{NN}} = 5.02$ TeV will not be analyzed.

Figure 7.5 shows that, as the charged particle multiplicity increases with system size, the mean of the $\Xi(1820)^\mp$ shows a slight decrease, except for the kCent triggered Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data. Comparison of the Pb-Pb kINT7 triggered (minimum bias) data point (not shown) to the p-p minimum bias 0-100 % multiplicity percentile data point shows a notable lower value, outside

the error of the two values. The difference between these two data points corresponds to a 1.11σ difference when all error bars are accounted for. To further complement this trend is the kSemiCent triggered data for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV, which has a 2.13σ difference. However, the kCent triggered data point, that should have the highest decrease in mass if this trend was to continue, has a mass incredibly similar to the p-p minimum bias 0-100 % multiplicity percentile data point.

A very notable feature of Figure 7.5 is the size of the statistical and systematic uncertainties for the Pb-Pb data points in comparison to p-p and p-Pb data points. It is normally assumed that systematic uncertainty should be similar across all data points if the procedure is the same for all of the analysis, which is mostly the case for this analysis procedure. However, when comparing the systematic uncertainties of Pb-Pb data from Table 7 to p-p data or p-Pb data from Tables 5 or 6 respectively, it is obvious that the sources of systematic uncertainties and total systematic uncertainties for Pb-Pb are very different than for p-p or p-Pb. Furthermore, the statistical uncertainties increase in a similar style as the systematic uncertainties.

The most questionable feature of Table 8 is the 0-90% (kINT7 triggered) data for Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV. It contains the lowest mean, the highest statistical uncertainty, and the highest systematic uncertainty for any data point in Table 8. While it might be concerning that the Pb-Pb triggered data set with the highest statistics also has the highest statistical uncertainty, it also contains the highest amount of background. This, along with the high systematic uncertainty, is the main reason that the Pb-Pb kINT7 triggered data point is not included in Figure 7.5.

While there does appear to be a slight decrease in the mean of the $\Xi(1820)^\mp$ with increased charged particle multiplicity, which could be interpreted as a signature of chiral symmetry restoration, a conclusive statement is not possible given the kCent triggered Pb-Pb data point. However, this trend of decreasing mean with the increasing system size does give encouragement that further analysis with potentially higher statistics and greater systematic analysis could provide better evidence of a chiral symmetry restoration signal seen in the mean of $\Xi(1820)^\mp$. Based on the calculations shown in Section 3.5.1, there is expected to be more of an effect in the width measurements than in the mean measurements.

As mentioned previously, the difference between the $\Xi(1820)^0$ and $\Xi(1820)^\mp$ resonances is informative so that a distinction can be made between the neutral and charged state that have been considered the same in the PDG [12]. With the $\Xi(1820)^0$ having a $7.0 \text{ MeV}/c^2$ lower mass when compared to the $\Xi(1820)^\mp$, this would show a similar trend seen for other Ξ states. Furthermore, this difference corresponds to a 3.4σ difference when all uncertainties are accounted for.

7.2 Width

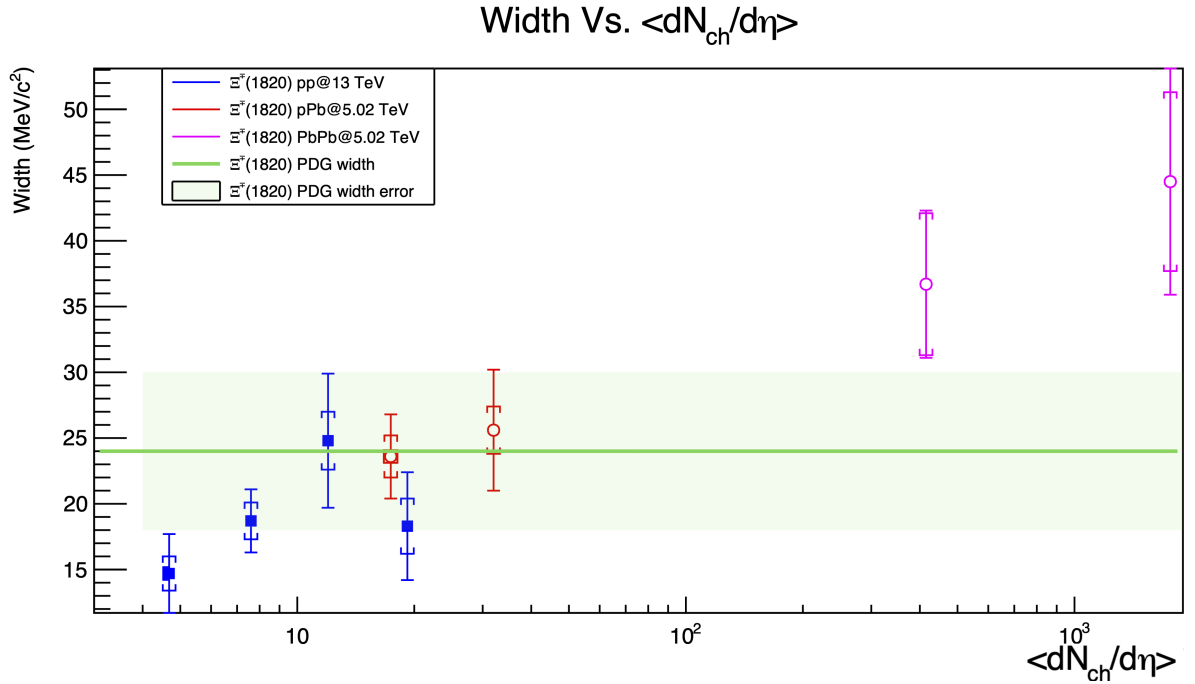


Figure 7.6: Width of the Voigtian fits of $\Xi(1820)^\mp$ for different system sizes and multiplicity percentile. Statistical uncertainties are shown as lines and systematic uncertainties are shown as capped boxes.

The width of the $\Xi(1820)^\mp$ for p-p, p-Pb, and Pb-Pb collisions is shown for Figure 7.6. The first, third, and fourth blue data points (30-100 %, 10-30 %, and 0-10 % respectively) for Figure 7.6 are subsets of the second data point (0-100 %) for p-p $\sqrt{s} = 13 \text{ TeV}$ data, thus the minimum bias (0-100 %) p-p data point is chosen to make the final comparison to the Pb-Pb data. The width from the kINT7 triggered Pb-Pb $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ data is not included in Figure 7.6 due to the

larger statistical uncertainties when compared to other Pb-Pb triggered data sets. Individual data points for the width, statistical uncertainty, and systematic uncertainty for $\Xi(1820)^\mp$ are shown in Table 9. $\Xi(1820)^0$ for p-p $\sqrt{s} = 13$ TeV at 0-100 % is included at the bottom of the table only so that it can be compared to $\Xi(1820)^\mp$ for p-p $\sqrt{s} = 13$ TeV at 0-100 %.

Collision Sys, Multi. Percentile (%)	$\langle dN_{ch}/d\eta \rangle \pm \text{Sys}$	Width $\pm \text{Stat} \pm \text{Sys}$ (MeV/ c^2)
p-p, 30-100 %	4.68 ± 0.16	$14.7 \pm 3.0 \pm 1.3$
p-p, 0-100 %	7.60 ± 0.14	$18.7 \pm 2.4 \pm 1.4$
p-p, 10-30 %	12.0 ± 0.3	$24.8 \pm 5.1 \pm 2.2$
p-Pb, 0-100 %	17.4 ± 0.7	$23.6 \pm 3.2 \pm 1.6$
p-p, 0-10 %	19.2 ± 0.5	$18.3 \pm 4.1 \pm 2.1$
p-Pb, 0-30 %	32.0 ± 0.8	$25.6 \pm 4.6 \pm 1.8$
Pb-Pb, 30-50 %	415 ± 10	$36.7 \pm 5.6 \pm 5.4$
Pb-Pb, 0-90 %	545 ± 6	$45.6 \pm 10.4 \pm 10.6$
Pb-Pb, 0-10 %	1765 ± 36	$44.5 \pm 8.6 \pm 6.8$
p-p, 0-100 % $\Xi(1820)^0$	7.60 ± 0.14	$19.4 \pm 5.2 \pm 4.3$

Table 9: All $\Xi(1820)^\mp$ width calculations for p-p $\sqrt{s} = 13$ TeV, p-Pb $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data. Statistical and systematic uncertainties are included. $\Xi(1820)^0$ width for p-p $\sqrt{s} = 13$ TeV at 0-100 % is included [40] [41] [42] [43].

Figure 7.6 shows that as the charged particle multiplicity increases with system size, the width of the $\Xi(1820)^\mp$ signal increases. Comparison of the Pb-Pb “central” kCent triggered (0-10 %) data point to the p-p minimum bias 0-100 % multiplicity percentile data point shows a higher value that corresponds to a 2.38σ difference.

The increase in the width of the $\Xi(1820)^\mp$ with increased charge particle multiplicity would appear to be a hint of the signature of chiral symmetry restoration. While normally a 3σ difference between two data points or a model is required for a result to be considered “evidence”, a 2σ difference would still be very notable and encouraging. The 2.38σ difference seen in this data is more than encouraging, but a less than conclusive (3σ) result. Furthermore, a clear increasing trend of width vs. $dN_{ch}/d\eta$ can be seen from the minimum bias cases for the p-p, p-Pb, and Pb-Pb systems. The main issue with any conclusive statement is, as is often the case, the uncertainty of the largest multiplicity point. Further statistics for Pb-Pb $\sqrt{s_{NN}} = 5.02$ would prove instrumental in reducing these errors and increasing the confidence in this signal.

Another important feature of Figure 7.6 is the comparison of the p-p $\sqrt{s} = 13$ TeV minimum bias data point to the PDG value, that shows a 0.80σ difference [12]. While the width of the p-p $\sqrt{s} = 13$ TeV minimum bias data point does appear to be on the lower end of the PDG value and error, it should be noted that this Voigtian function does have a $5 \text{ MeV}/c^2 \sigma$. If the same invariant mass plots were fit with a normal Breit-Wigner function, the value of the width would increase by about $5 \text{ MeV}/c^2$, reaching much closer to the value listed in the PDG. While having the minimum bias data point closer to the value listed in the PDG would seem to further complement our analysis when compared to other studies, it would not change the results of the comparison between the Pb-Pb central 0-10 % multiplicity percentile data point to the p-p minimum bias 0-100 % multiplicity percentile data point. If both were fit with the normal Breit-Wigner function, both would have an increase in width of about $5 \text{ MeV}/c^2$, the difference between them would not change.

7.3 Yield

Figure 7.7 shows the completed yield calculations and p_T spectra for p-p collisions at $\sqrt{s} = 13$ TeV for the 0-100 % multiplicity percentile. As stated before, each p_T bin was selected to have approximately the same raw yield in each, so that yield comparisons would be easier for analysis into how the yield of the $\Xi(1820)$ might change. Each fit was reweighted using the procedure discussed in Section 5.2.

The individual data points used for the p_T spectra are listed in Table 10. These yield calculations have not been reweighted using the procedure discussed in Section 5.2.

p_T range (GeV/c)	Final Yield ($\frac{\text{TotalYield}}{(a*\epsilon_{rec}) \times \text{NumberOfEvents} \times dy \times \Delta p_T \times B.R.}) \pm \text{Stat} \pm \text{Sys}$)
$1.0 < p_T < 2.0$	$(1.44 \pm 0.21 \pm 0.11) * 10^{-4}$
$2.0 < p_T < 2.5$	$(1.44 \pm 0.17 \pm 0.15) * 10^{-4}$
$2.5 < p_T < 3.0$	$(8.3 \pm 0.9 \pm 1.2) * 10^{-5}$
$3.0 < p_T < 3.7$	$(5.1 \pm 0.5 \pm 0.5) * 10^{-5}$
$3.7 < p_T < 7.0$	$(7.7 \pm 0.7 \pm 0.7) * 10^{-6}$

Table 10: All $\Xi(1820)^\pm$ yield calculations for p-p $\sqrt{s} = 13$ TeV 0-100 % data. The 1st and 2nd uncertainties are the statistical and systematic uncertainties respectively. No reweighted procedure applied.

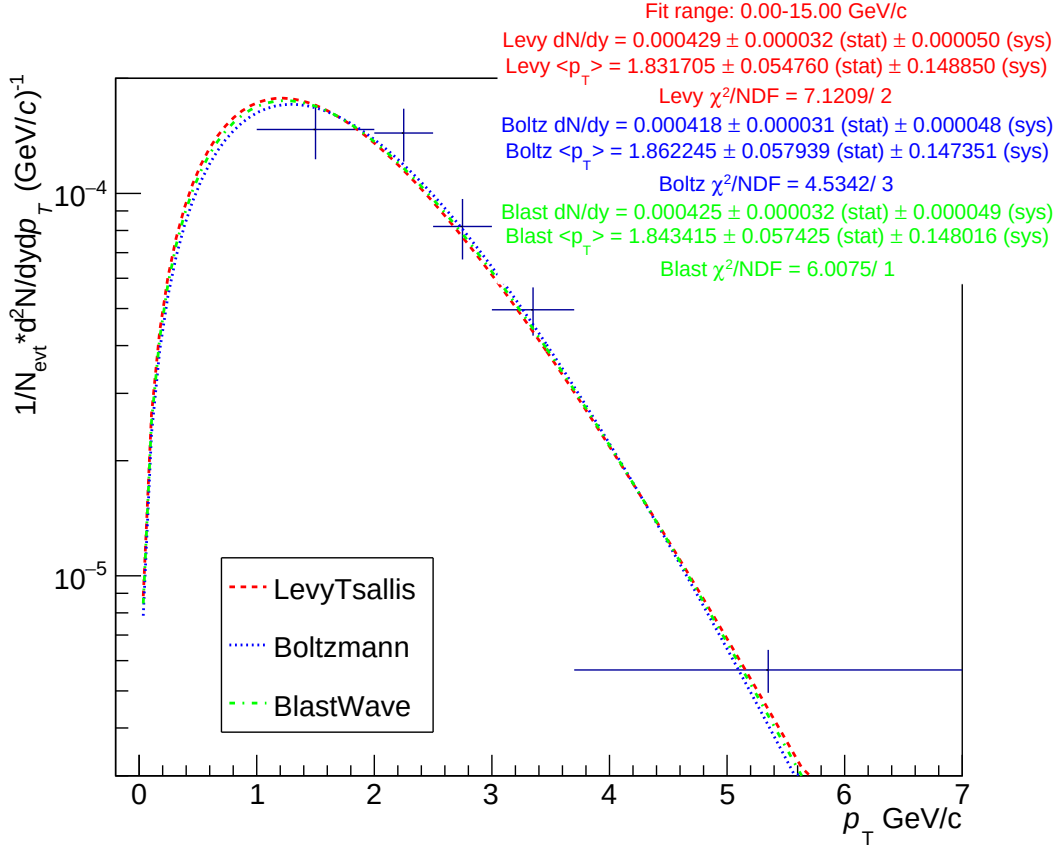


Figure 7.7: Reweighed corrected p_T spectra for $\Lambda+K^\mp$ from p-p collisions at $\sqrt{s} = 13$ TeV with 0-100 % multiplicity percentile. Levy-Tsallis (Red), Boltzmann (Blue), and Boltzmann-Gibbs Blast-Wave (Green) functions are shown after all iterations of the reweighing procedure have been completed. Both statistical and systematic uncertainty have been combined into one error bar.

It should be noted that the p_T bins used to calculate yield and spectra are slightly different than what was described previously. While the mean and width of the previous plots were calculated for the p_T ranges of $1.0 < p_T < 20.0$ GeV/c, the spectra was calculated using subranges of the total p_T range. Furthermore, while the previously used invariant mass plots from Section 4.4 had the final p_T bin correspond to $3.7 < p_T < 20.0$ GeV/c, the final p_T bin for the spectra corresponds to $3.7 < p_T < 7.0$ GeV/c. The reduction of the last p_T bin for the spectra analysis is due to the need for the fit functions of the p_T spectra to have reasonable values and uncertainties. For example, if the final bin is changed to $3.7 < p_T < 20.0$ GeV/c and fit using the fit procedure discussed, then

the $\langle p_T \rangle$ would increase by 30 %, a completely impossible increase that would occur from a contribution that should have a minimum influence on the $\langle p_T \rangle$.

Another important note is the yield fit function range used to determine the total yield. Once the Levy-Tsallis function is fit for any range of data points, it can be extrapolated to ranges that are not listed in the data points, such as the $0 < p_T < 1.0$ GeV/ c bin. However, the main concern is the high end of the fit function. As mentioned previously, Monte Carlo data only simulates $\Xi(1820)$ up to 15 GeV/ c . Thus, the fit function to determine the total yield, which requires the Monte Carlo corrections, is fit only from $0 < p_T < 15.0$ GeV/ c .

As has been discussed in previous sections, the main reason for the functions used to fit the p_T spectra is to determine the total yield of the $\Xi(1820)$. Given the limited statistics of the $\Xi(1820)$ and the general shape of the fits, it is clear that not all of the yield is encapsulated in the data points used to determine the p_T spectra. Analysis into the total yield from the fit functions shows that about 75 % of the total yield is accounted for in the selected p_T bins. About 25 % of the total yield would come from extrapolating the fit functions to higher and lower p_T ranges such as $0 < p_T < 1.0$ GeV/ c and $7.0 < p_T < 15.0$ GeV/ c .

The integrated yield calculation for the Levy-Tsallis function is selected as the “default” fit and used to determine the statistical uncertainty. The difference between the Levy-Tsallis function and the other 2 fit functions is incorporated into the systematic uncertainty. Thus, the final integrated yield for $\Xi(1820)$ from p-p collisions at $\sqrt{s} = 13$ TeV is $dN/dy = 429 \pm 32$ (stat.) ± 51 (sys.) $\times 10^{-6}$. This calculation is important for any comparison of the yield of the $\Xi(1820)$ to other experiments.

Another important value from these fits is the mean transverse momentum $\langle p_T \rangle$ which can be regarded as the average transverse momentum of the $\Xi(1820)$ for this system. Normally, the $\langle p_T \rangle$ has a value similar to the mass of the particle it is fit for. The final mean transverse momentum for $\Xi(1820)$ from p-p collisions at $\sqrt{s} = 13$ TeV is $\langle p_T \rangle = 1.83 \pm 0.06$ (stat.) ± 0.15 (sys.) GeV/ c , very similar to the expected mass of the $\Xi(1820)$. While the uncertainties for the $\langle p_T \rangle$ may seem large, $\langle p_T \rangle$ is not used for any calculations and serves more as a check that

the fit functions behave as expected.

One interesting feature to note is the relation between the systematic and statistical uncertainty when comparing the integrated yield and $\langle p_T \rangle$. For the $\langle p_T \rangle$, the systematic uncertainty is almost 3 times the statistical uncertainty, while the integrated yield has the systematic uncertainty about 50 % higher than the statistical uncertainty. This most likely relates to the property of these fit functions to be very sensitive to the p_T bins and ranges used for the calculations of the $\langle p_T \rangle$ as opposed to the integrated yield of the fit functions. This is also most likely why, when the final p_T bin is changed from $3.7 < p_T < 7.0$ GeV/ c to $3.7 < p_T < 20.0$ GeV/ c , it is the $\langle p_T \rangle$ that is affected the most, increasing by 30 %, while the integrated yield is far less affected.

7.4 Comparison to other particles

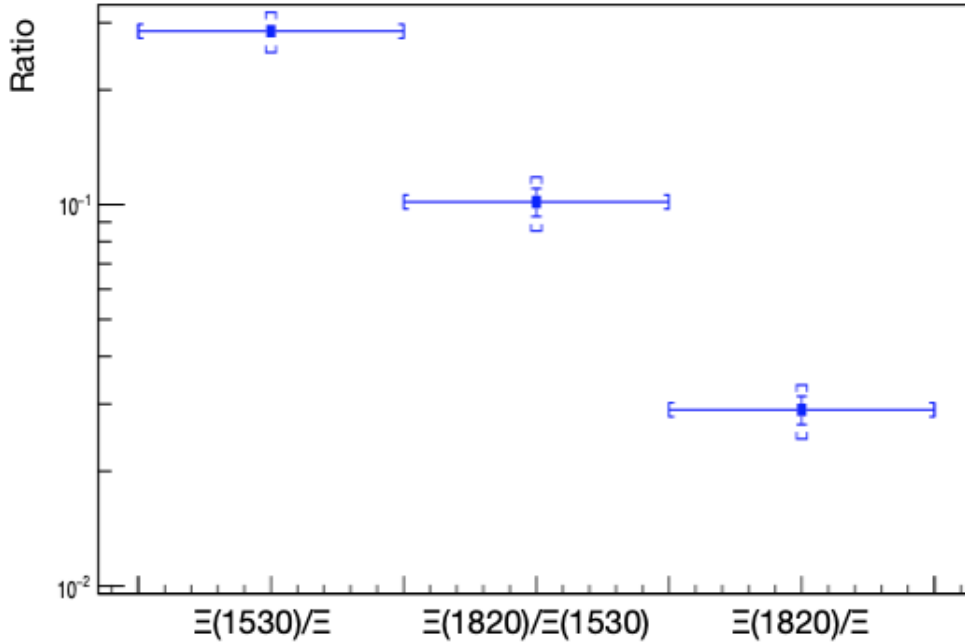


Figure 7.8: Ratio of corrected yields fits of $\Xi(1322)$, $\Xi(1530)$, and $\Xi(1820)$ for minimum bias p-p $\sqrt{s} = 13$ TeV 0-100 % data [60] [61]. Statistical uncertainties are shown as lines and systematic uncertainties are shown as capped boxes.

The yield ratios of $\Xi(1530)/\Xi$, $\Xi(1820)/\Xi(1530)$, and $\Xi(1820)/\Xi$ for minimum bias p-p $\sqrt{s} =$

13 TeV 0-100 % are shown in Figure 7.8 [60] [61]. Table 11 shows the values of the ratios from Figure 7.8. Do note that the $\Xi(1820)^\mp \rightarrow \Lambda + K^\mp$ branching ratio is not well known and has been set to 100 % as a high limit evaluation.

Particles used for ratio	Ratio $\pm$ Stat $\pm$ Sys
$\Xi(1530)/\Xi$	$0.285 \pm 0.005 \pm 0.034$
$\Xi(1820)/\Xi(1530)$	$0.102 \pm 0.009 \pm 0.016$
$\Xi(1820)/\Xi$	$0.029 \pm 0.002 \pm 0.005$

Table 11: All ratio calculations for $\Xi(1820)$, $\Xi(1530)$, and Ξ for p-p $\sqrt{s} = 13$ TeV 0-100 % data. Statistical and systematic uncertainties are included [60] [61].

In the event that chiral symmetry is restored and the mass of the particle has decreased, then the yield of $\Xi(1820)$ should decrease in comparison to chiral symmetry remaining broken. As stated before, the decay channel of interest for this discussion is the $\Lambda + K$ channel. In theory, the mass of the $\Xi(1820)$ could drop to such a point that it would cross the minimum mass threshold. In such an event, the $\Xi(1820)$ would not decay into the $\Lambda + K$, but into another decay channel such as $\Xi + \pi$. Regardless, the $\Xi(1820)$ would be lost to our decay channel and as such would not be reconstructed, thus having a lower yield than if the $\Xi(1820)$ kept its same mass. This, however, is a very extreme case and is expected to have a minor effect, if any, on the calculations.

The more likely case is that, due to the lower mean or higher width of the signal, as seen in Figure 3.1, some of the yield will simply not be counted in the calculations. Similar to the previous case, though not as extreme, the $\Xi(1820)$ may drop its mass to a low value, 1700 MeV/ c^2 for example, and this mass, while affecting the fit calculations to some small degree, will not be enough to be measured. Thus, even with the new fit, the yield will appear to be lower.

Because of the fact that the Pb-Pb collisions will by their very nature have about 50 times the number of particles being produced in a collision, regardless of the creation of the QGP or chiral symmetry restoration, a ratio needs to be created that can help observe the potential effects of the changing yields. According to the lattice QCD calculations from Figure 1.8, the $\Xi(1530)$ should have its mass independent of the temperature of the system. So, if a ratio is constructed between a particle that may have its yield change over a particle that should not have its yield change, such as

$\Xi(1820)/\Xi(1530)$, then if the ratio decreases with increased particle multiplicity, chiral symmetry restoration is indicated.

To determine if there is any signature of chiral symmetry restoration, a change in the ratio of $\Xi(1820)/\Xi(1530)$ should be observed. For such a ratio to be calculated, p_T spectra for a system with chiral symmetry still broken and a system with chiral symmetry restored needs to be constructed. While possible in the p-p $\sqrt{s} = 13$ TeV minimum bias case, there is just not enough statistics available in the Pb-Pb $\sqrt{s_{NN}} = 5.02$ TeV data. As stated before, there is barely enough data for the kINT7 triggered Pb-Pb data to create an invariant mass plot, let alone the need to further divide the invariant mass plot into separate p_T ranges and fit with good quality. As such, while the $\Xi(1820)$ ratios for p-p $\sqrt{s} = 13$ TeV are interesting, they are simply not enough to provide any further information on chiral symmetry restoration. Further investigation into Pb-Pb data sets with higher statistics or detailed chiral symmetry restoration model analysis would prove helpful in finding any possible signature.

8 Conclusion and Discussion

In conclusion, the investigation into the $\Xi(1820)$ resonance has produced several important results. While some of these results are minor in comparison to others, each one is an important feature that has not been measured to the same accuracy in past experiments [12].

First, there is a notable difference between the $\Xi(1820)^\mp$ and $\Xi(1820)^0$ resonance mass and width. While some sources treat these two resonances as the same, such as the PDG, this analysis does confirm some of the differences [12]. While the difference between the widths of both resonances is small enough to be within error of each other, the difference between the masses of both resonances is large enough to have a 3.4σ difference when statistical and systematic uncertainties are included. This difference is the largest obtained in this analysis, but is again limited by the statistics of the analysis and low yield of the $\Xi(1820)^0$ resonance. Larger statistics and more analysis would provide a way to conclusively divide these two resonances into separate values used in future analysis.

Second, the basic information of the mass and width for the $\Xi(1820)^\mp$ resonance from p-p collisions at $\sqrt{s} = 13$ TeV for the minimum bias 0-100 % multiplicity percentile yields a much smaller uncertainty than any result that is listed in the PDG [12]. In some cases, the uncertainty in the mass and width from this analysis are about 1/3 of the uncertainties from the PDG [12]. This analysis can use the results of mean and width from p-p collisions at 0-100 % minimum bias to better compare to Pb-Pb kCent triggered data with higher certainty than what can be found in the results from the PDG [12]. In addition, there is a suggestion to add these results to the PDG summary results for the $\Xi(1820)$, to further improve the standard properties of the $\Xi(1820)$.

Third, the investigation into chiral symmetry restoration using the yield, mass, and width of the $\Xi(1820)$ is very interesting. The yield ratio analysis for the $\Xi(1820)$ in p-p collisions at $\sqrt{s} = 13$ TeV, while interesting, provides no comment on the potential effects seen from parity doubling and chiral symmetry restoration. More statistics, especially for the Pb-Pb data, are required before any conclusive statement can be made about chiral symmetry restoration. The yield ratios analysis can not provide any comment on chiral symmetry restoration until sufficient statistics in

higher multiplicity densities are available and a realistic estimate of the branching ratio for the analyzed channels can be obtained. The mass analysis for the $\Xi(1820)$ is much more informative than the yield analysis. There does appear to be a slight trend of decreasing mass of the $\Xi(1820)$ with increasing $\langle dN_{ch}/d\eta \rangle$. In fact, comparison of kINT7 and kSemiCent triggered Pb-Pb data to p-p minimum bias data shows a 1.11 and 2.13 σ difference respectively. However, the inability of this trend to be seen in kCent triggered Pb-Pb data prevents any conclusive statement from being made about chiral symmetry restoration. The width analysis for the $\Xi(1820)$ is the most informative analysis for this discussion. There is a notable general trend of increasing width with increasing $\langle dN_{ch}/d\eta \rangle$. More notable than the mass analysis, this 2.38 σ difference is close to being considered evidence for a statement about chiral symmetry restoration to be made. These results, in particular the width analysis of the $\Xi(1820)^\mp$, would seem to indicate evidence of parity doubling, which is a signature of chiral symmetry restoration, but is currently not conclusive enough. Furthermore, the fact that this analysis was inspired by lattice QCD calculations brings more confidence in their predictive power and acceptance. More investigation into the $\Xi(1820)$ would prove to be very informative about chiral symmetry restoration analysis, but more statistics is required.

While the investigation into chiral symmetry restoration using the $\Xi(1820)$ has led to very interesting results, there is more than enough here to be very encouraged for future analyses using this same procedure. While currently using the highest statistics available at ALICE is still not enough, with the potential improvements and further analysis at ALICE and the LHC currently underway, it may only be a matter of time. In fact, the current upgrades implemented in ALICE for the future campaign starting in 2021 include increasing the event readout rate by a factor 30-50 up to 50 kHz for Pb-Pb collisions [62], thus greatly increasing the rate of collecting Pb-Pb data to a point that a full analysis into the $\Xi(1820)$ can be completed. Still, more statistics for the $\Xi(1820)$ and a potentially independent verification through another parity partner pair is needed. After all, the ρ analysis for the NA60 experiment and the potential chiral symmetry restoration evidence is still being debated, even 10 years after its publication [19]. Other chiral symmetry restoration

investigation efforts, like the lattice QCD calculations from the Swansea collaboration, seem to only look for the potential effects using simulations and not physical data [24]. This analysis, which is very encouraging, may be the foundation for a whole new wave of experiments and analysis into chiral symmetry restoration. This analysis illustrates one of the most unique approaches that can be used for science, the ability to investigate something that may seem so small and insignificant, but can lead to the complete redefinition of one of the fundamental properties of the universe.

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