

Sequential Strangeness Freeze-out via Statistical Hadronization in Relativistic Heavy Ion and Elementary Particle Collisions

by
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*for Ninis, Pepé, and Tony –
all of who, decades ago, first interrupted my dogmatic slumber.*

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ABSTRACT

Quantum Chromodynamics (QCD), the theory describing the interaction between quarks and gluons, suggests that the coupling strength of the strong interaction asymptotically decreases at large momentum scales. Hence, at large enough densities and/or temperatures – like the conditions present in the early Universe – the quarks and gluons composing hadrons can exist in a deconfined state of matter known as the Quark-Gluon Plasma (QGP).

A Large Ion Collider Experiment (ALICE) at the Large Hadron Collider (LHC) and the Solenoidal Tracker at RHIC (STAR) experiment at the Relativistic Heavy Ion Collider (RHIC) have gathered experimental evidence of QGP formation generated through relativistic heavy ion collisions. The QGP lifetime is accepted to be on the order of 10^{-21} seconds, with a temperature reaching 10^{12} K, before the system cools and undergoes a phase transition, back into confined ground-state hadrons and hadronic resonances. Whether this transition from quark to hadron degrees of freedom occurs at a uniform temperature for all quark flavors remains a question of interest.

Assuming a thermally equilibrated system, experimental particle yields from ALICE and STAR serve as anchors for the determination of common freeze-out parameters – namely, the chemical freeze-out temperature (T_{ch}) and the baryo-chemical potential (μ_{B}) – in the QCD phase diagram through thermal fits within the Statistical Hadronization Model (SHM) framework. In this dissertation, I present evidence supporting a flavor-dependent chemical freeze-out in the crossover region of the QCD phase diagram via SHM calculations of the freeze-out parameters at ALICE and STAR at collision energies ranging from $\sqrt{s_{\text{NN}}} = 11.5$ GeV to 5.02 TeV. I also show the system size dependence of T_{ch} at vanishing μ_{B} via the same approach in pp, pPb and PbPb collisions at LHC energies – supporting the applicability of the SHM in small collision systems and that flavor-dependent freeze-out parameters lead to a natural explanation of strangeness enhancement from small to large collision systems at ALICE.

At first there was only darkness wrapped in darkness; all that existed then was void and formless. That One which came to be, enclosed in nothing, arose at last, born of the power of heat.

– Rig Veda, *Mandala X; Hymn CXXIX: Creation* (c. 1700 – 1100 BC)

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1 INTRODUCTION

Cosmological observations suggest our Universe emerged from a rapidly expanding single point in spacetime approximately 13.8 billion years ago. While the notion of an expanding fireball generating the Cosmos dates back to antiquity, the contemporary model describing this phenomenon did not materialize until the early 20th century.

In 1922, following astronomical observations of neighboring galaxies receding from Earth by Vesto Slipher and Carl Wilhelm Wirtz [1–3], Alexander Friedmann derived a set of solutions to Albert Einstein’s field equations of General Relativity; one of which suggested a non-static, expanding Universe [4, 5]. Shortly thereafter, Georges Lemaître developed the theoretical framework of an expanding Universe whose genesis could be traced back in finite time to a point-like *primeval atom* in space by observing a direct relationship between the recessional velocities of galaxies and their apparent distances from Earth [6, 7]; this was later confirmed by Edwin Hubble in 1929 [8]. By 1931, Lemaître’s highly regarded primeval atom hypothesis, subsequently known as the Big Bang Theory, became the reigning paradigm describing the evolution of the observable Universe [9]. Through this paradigm, it is possible to chronicle the evolution of the Universe after the Big Bang. A terse, non-inclusive Universal timeline leading to the creation of hadronic matter is given below.

Immediately after the Big Bang, the Universe’s average temperature and energy density were high enough for the gravitational, electromagnetic, weak and strong interactions of nature to be indistinguishable; the Universe continued to rapidly expand and began to cool. Approximately 10^{-43} seconds following this initial expansion, the Universe reached a temperature around 10^{33} K [10]. During this epoch, the gravitational force decoupled from the other three fundamental forces. As the temperature dropped to 10^{28} K, a phase transition occurred. During this transition, about 10^{-36} seconds after the Big Bang, the strong nuclear force decoupled from the electromagnetic and weak nuclear forces. At this

temperature scale, the latter two forces remain unified as a single interaction known as the *electroweak force* [11–18]. The Universe continued to rapidly expand and supercooled to a temperature of 10^{22} K, 10^{-32} seconds after the big bang, upwardly scaling its size by a factor on the order of 10^{26} . Throughout this period of cosmic inflation [19], the temperature of the Universe was much too high for composite subatomic particles, known as *hadrons*, to coalesce. Instead, a plasma of deconfined quarks and gluons is present as state of matter known as the *Quark-Gluon Plasma* [20, 21]. As the age of the Universe approached 10^{-12} seconds and reached a temperature of 10^{15} K, the electroweak force underwent a phase transition and decoupled into the weak nuclear force and the electromagnetic force [22], rendering the four distinct fundamental interactions prevailing the present day Universe. However, at this temperature scale, the Quark-Gluon Plasma (QGP) persisted. Figure 1.1 shows the temperature evolution of the early Universe as a function of time after the electroweak phase transition took place. Around 10^{-5} seconds in the Cosmic timeline, the temperature dropped below 10^{12} K and induced a third phase transition; at this temperature the QGP transitioned into a hadronic gas as quarks were confined into hadrons [23, 24]. This transition is accessible through modern experimental methods [25] and can be understood from a thermodynamic perspective [26]. Whether this transition from quark to hadron degrees of freedom occurred at a uniform temperature for all quark species is the primary focus of this dissertation.

Relativistic heavy ion collisions serve as an excellent experimental tool for QGP production, where the generated temperature and energy densities reach scales similar to those present microseconds after the Big Bang; providing a laboratory setting where the QGP transition into hadronic matter can be thoroughly investigated. In the late 1970s, experimentation with relativistic heavy ions began at the Joint Institute for Nuclear Research (JINR) and Lawrence Berkeley National Laboratory (LBNL) [27–29]. Shortly thereafter, these efforts were followed by various fixed target experiments at both the Alternating Gradient

Synchrotron (AGS) in Brookhaven National Lab (BNL) and the Super Proton Synchrotron hosted by the European Organization for Nuclear Research (CERN) [30,31]. By the turn of the 21st century, the Relativistic Heavy Ion Collider (RHIC) at BNL and the Large Hadron Collider (LHC) at CERN confirmed signatures of QGP formation via relativistic heavy ion collisions [32].

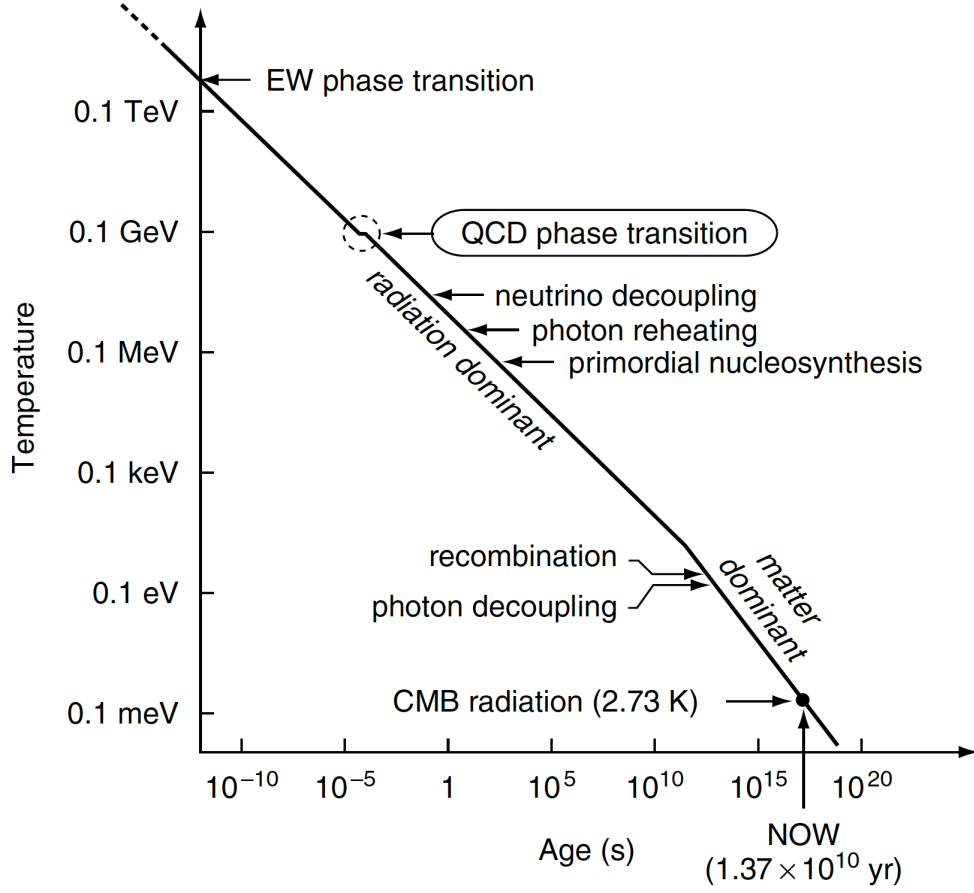


Figure 1.1: Temperature evolution of the early Universe as a function of time in seconds following the electroweak phase transition. For reference, $1 \text{ eV} \approx 1.16 \times 10^4 \text{ K}$. Figure from [33].

While there exist various observables that may signal the transition from hadronic matter into a state of deconfined quarks and gluons, the present dissertation places its focus on the production of final state strange hadrons [34,35] measured in relativistic heavy ion

collisions at the Solenoidal Tracker at RHIC (STAR) in BNL and A Large Ion Collider Experiment (ALICE) at the LHC. In more recent studies, the ALICE Collaboration showed that strangeness production in high-multiplicity proton-proton (pp) collisions reaches values similar to those achieved by heavy ion collisions – namely, lead-lead (PbPb) – where QGP formation occurs, and agree with results from proton-lead (pPb) collision results. [36].

Final state particle yields can be successfully reproduced by Statistical Hadronization Models (SHMs) to nine orders of magnitude over a wide energy range in high energy collisions of heavy ions [37, 38]. Assuming thermal equilibrium, SHMs can determine common freeze-out parameters in the cross-over region of the Quantum Chromodynamics (QCD) phase diagram, in particular, the baryo-chemical potential (μ_B) and the chemical freeze-out temperature (T_{ch}), from experimental particle yields measured by ALICE and STAR. These calculations are usually performed with a single value for T_{ch} , assuming all particle species “freeze-out” uniformly. Nevertheless, a point of contention arises when comparing the extracted freeze-out parameters obtained using different sets of particle species in the thermal fit; namely, the difference in T_{ch} between the light hadrons and the strange hadrons. These parameter calculations can be compared to measurements of higher order fluctuations of net-particle distributions and/or susceptibility calculations of conserved quantum numbers from lattice QCD, where a difference in the determined freeze-out temperatures between flavors has previously been shown [39].

This dissertation presents calculations of the freeze-out parameters from STAR and ALICE at center-of-mass collision energies ranging from $\sqrt{s_{NN}} = 11.5$ GeV to 5.02 TeV within the SHM framework and provides evidence for a flavor-dependent chemical freeze-out in the crossover region of the QCD phase diagram. Additionally, thermal model calculations of experimental yield ratios for several multiplicity classes in pp, pPb and PbPb collisions measured by the ALICE detector are performed employing a flavor-dependent chemical freeze-out

framework and are shown to be in excellent agreement to experimental data.

The general outline of the dissertation is as follows: The present chapter introduces the Standard Model of Particle Physics, Quantum Chromodynamics, characteristics of the QGP, as well as an overview of heavy-ion collisions. Chapter 2 provides a brief history of strangeness as well as a description of strangeness production as a signature of QGP formation; the notion of sequential strangeness freeze-out is then introduced. The ALICE detector is outlined in Chapter 3 and a comparison is made to the STAR experiment. The Statistical Hadronization Model is described in Chapter 4, as well as the analysis methodology used to acquire the main results of the dissertation. The results acquired from the analysis described in Chapter 4 are shown and discussed in Chapter 5. Lastly, Chapter 6 serves as the conclusion.

1.1 The Standard Model of Particle Physics

Throughout the latter half of the present century, experimental developments in high energy physics have established a descriptive model of elementary particles and their interactions. This generalized framework, known as the Standard Model of Particle Physics, describes three of the four fundamental forces in the Universe (the electromagnetic, weak, and strong interactions) through a local $SU(3) \times SU(2) \times U(1)$ gauge symmetry [40, 41].

The Standard Model is composed of two integral parts: fermions and bosons. The former refer to the matter-generating elementary particles which make up a portion of the observable universe, while the latter correspond to the force carriers which mediate the interactions between fermions [42]. Figure 1.2 depicts a schematic of the elementary particles of the Standard Model including their respective mass, charge, and spin.

Fermions are categorized into three generations of quarks and leptons, which possess half-integer spin values and are subject to the Pauli Exclusion Principle [43]. Each of these twelve particles possesses an associated antiparticle with all features, such as spin and mass, being

identical to their particle counterparts and differ only by an inversion of their respective charges. The fermions in the Standard Model are shown on the left-hand side of Figure 1.2, stratified into the three generations of quarks (white) and leptons (gray).

Quarks consist of six flavors: up (u), down (d), strange (s), charm (c), bottom (b), and top (t). The second generation (s and c) and third generation (b and t) quarks are unstable and weakly decay into quarks from previous generations – in this formulation, the only generation of stable quarks is the first (u and d).

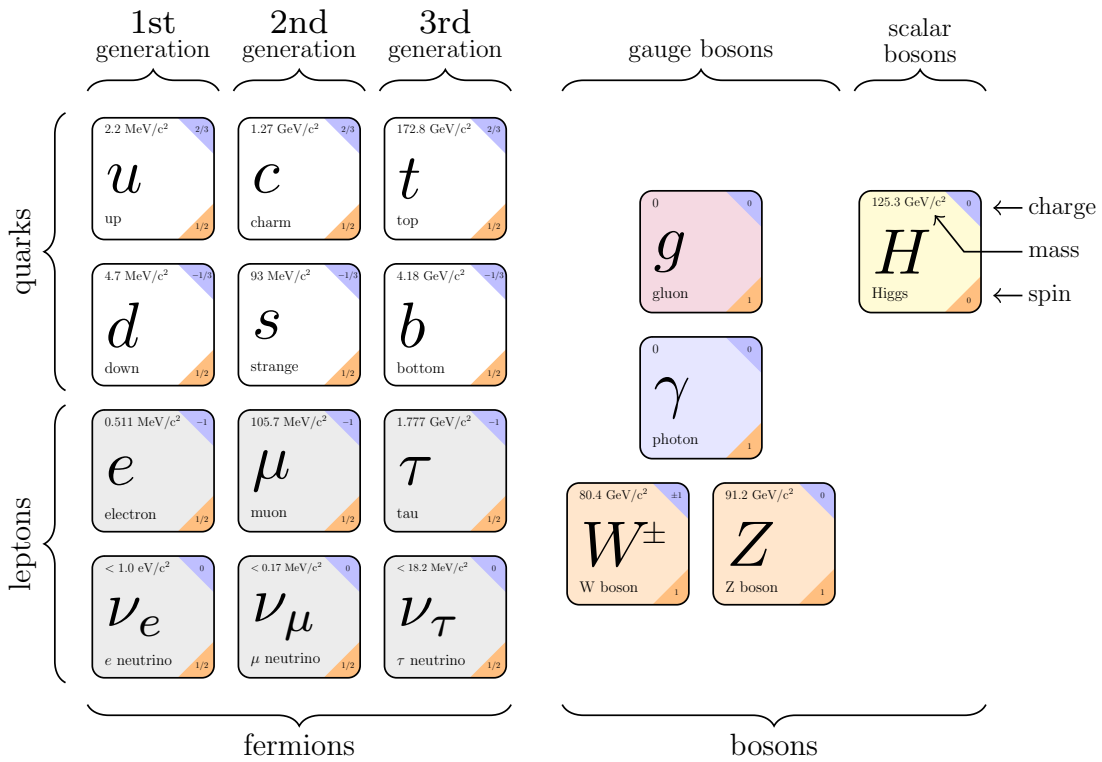


Figure 1.2: Elementary Particles of the Standard Model of Particle Physics, composed of 12 fermions and 5 bosons. Anti-fermions are not shown. Figure adapted from [44]. Data from [45].

The three generations of leptons are classified as: The electron (e), the muon (μ), and the tau (τ), each with a corresponding neutrino (ν_e , ν_μ , and ν_τ , respectively). While the μ and τ leptons decay into electrons, neutrinos do not decay and simply oscillate between

flavors instead. Neutrino flavor oscillations were predicted by Bruno Pontecorvo [46], with the first observations taking place in the early 1960s in solar neutrinos [47, 48]. Almost four decades later, more conclusive measurements of neutrino oscillations occurred at the Super-Kamioka Neutrino Detection Experiment (Super-K) [49], the Sudbury Neutrino Observatory (SNO) [50] and were confirmed by various other experiments [51].

On the other hand, bosons have integer spin values and are not subject to Pauli Exclusion. In the Standard Model framework, the fundamental forces are each mediated by the exchange of a spin 1 gauge boson between interacting particles possessing specific charges [52]. The fundamental forces and their respective bosons are the following: The weak nuclear force, which describes processes involving neutrinos (e.g. β -decay) is mediated by the Z and $W^\pm$ bosons, which govern the interactions between weakly charged particles. The electromagnetic force is mediated by the photon (γ); governing the interactions between electrically charged particles and is described by the theory of Quantum Electrodynamics (QED). Lastly, the strong nuclear force, responsible for binding quarks into hadrons and hadrons into atomic nuclei, is mediated by the gluon (g), which governs the interactions between particles carrying a color charge. The range of the strong nuclear force is over two orders of magnitude greater than that of the electromagnetic force and the weak interaction. The strong nuclear interaction is described by the theory of Quantum Chromodynamics and is addressed in detail in Section 1.2. Additionally included in the Standard Model is the recently discovered Higgs boson (H) [53, 54], which facilitates the generation of elementary particles masses [55, 56]. Unlike the γ , g , Z , and $W^\pm$, the Higgs boson has spin 0 and is thus a scalar boson. The right-hand side of Figure 1.2 depicts the bosons of the Standard Model stratified into gauge and scalar bosons; the different colors are used to distinguish the different processes mediated by each boson.

1.2 Quantum Chromodynamics (QCD)

Quantum Chromodynamics (QCD) is the quantum field theory delineating the fundamental dynamics of the strong interaction. The advent of QCD can be traced back to the early 1960's with the introduction of the quark model, independently by Murray Gell-Mann and Yuval Ne'eman [57, 58], as an organizational scheme for the onslaught of observed particles from cosmic ray and collider experiments at the time. A detailed account of the Gell-Mann description of the quark model, dubbed *The Eightfold Way*, is provided in Chapter 2. While this stand-alone model proved to be successful at the time, difficulties emerged when attempting to build a wavefunction of the experimentally observed Δ^{++} resonance – a bound state of three u quarks with identical spins – which violated the Pauli Exclusion Principle [52]. To remedy this, Oscar Wallace Greenberg proposed that quarks carry an additional quantum number [59], which became known as the *color charge*.

QCD represents the SU(3) in the $SU(3) \times SU(2) \times U(1)$ symmetry group of the Standard Model and is a non-abelian gauge theory describing the interactions between quarks and gluons via the color charge – the latter being an analog to the electric charge from QED. While similarities exist between QED and QCD, a notable difference is the fact that the QCD gluons themselves carry the color charge, whereas in QED, photons possess no electric charge. Moreover, contrasted to the single electric charge from QED, the QCD color charge exists in three different varieties: red (R), green (G), and blue (B).

Without loss of generality, applying the Feynman slash notation and assuming summation over repeated indices, the gauge invariant QCD Lagrangian is expressed as [41, 45]:

$$\mathcal{L}_{\text{QCD}} = \sum_q \bar{\psi}_{q,i} (i \not{D}_{ij} - m_q \delta_{ij}) \psi_{q,j} - \frac{1}{4} G_{\mu\nu}^a G_a^{\mu\nu}. \quad (1.1)$$

Here, $\psi_{q,i}$ are the spinors of the quark field for a given quark of flavor q and its pertaining

mass m_q . The indices i and j represent the color charge and range from 1 to $N_{\text{color}} = 3$. To reiterate, by definition $\not{D}_{ij} = \gamma^\mu D_{\mu ij}$, where γ^μ represents the Dirac γ -matrices and $D_{\mu ij}$ is the gauge covariant derivative given by $D_{\mu ij} = \partial_\mu \delta_{ij} - ig_s t_{ij}^c A_\mu^c$. Here, g_s represents the flavor-independent QCD coupling and $\mathcal{A}_\mu^c$ pertains to the gluon fields, which range from 1 to $N_{\text{color}}^2 = 8$. The eight generators of the SU(3) group are represented by t_{ij}^c and are generally expressed in a standard basis, such that $t^c = \frac{1}{2}\lambda^c$, where λ^c denotes the eight Gell-Mann Matrices given by:

$$\begin{aligned}\lambda^1 &= \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda^2 &= \begin{pmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda^3 &= \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \\ \lambda^4 &= \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}, & \lambda^5 &= \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}, & \lambda^6 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \\ \lambda^7 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}, & \lambda^8 &= \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}\end{aligned}$$

Lastly, the trailing $G_{\mu\nu}^a$ term in Equation 1.1 represents the gluon field strength tensor, analogous to the QED field strength tensor $F_{\mu\nu}^a$, and is given by:

$$G_{\mu\nu}^a = \partial_\mu \mathcal{A}_\nu^a - \partial_\nu \mathcal{A}_\mu^a + g_s f_{abc} \mathcal{A}_\mu^b \mathcal{A}_\nu^c \quad (1.2)$$

. Again, $\mathcal{A}_\mu^c$ are the gluon fields whose indices a , b , and c run from 1 to 8, while f_{abc} are the structure coefficients of the SU(3) group. The non-abelian structure of QCD emerges from the $g_s f_{abc} \mathcal{A}_\mu^b \mathcal{A}_\mu^c$ term, which reflects the gluon self-interaction.

1.2.1 Color Confinement

Gluons and quarks, unlike the electroweak gauge bosons (γ , Z , and $W^\pm$) and the leptons, are not experimentally observed as free particles. Instead, they bind into color-neutral states, namely, hadrons, which are composed of either three quarks or a quark-antiquark pair. Three quark states and quark-antiquark pairs are referred to as *baryons* and *mesons*, respectively. To form a color neutral object, one of each of the three quarks in a baryon must exclusively possess a red, green, and blue color charge. The same applies for antibaryons, except with the charges inversed. In the case of a meson, the constituent quark possess any of the three color charges, while the antiquark matches the color but has an opposite charge (i.e. anticolor charge).

This fundamental feature of QCD, known as *color confinement* [60], makes it such that the gluon field between two quarks generates a narrow string-like flux as they become separated, keeping a constant force per unit distance between the two quarks. In this framework, it is more energetically favorable for a new quark-antiquark pair to form rather than elongating the string-like flux between two quarks as they are separated. Such phenomenon is generally attributed to the self interaction between the gluon fields and leads to isolated quarks not being physically observable. The effective potential $V(r)$ between a quark-antiquark pair with a charge separation r can be properly described by phenomenological models, one of which is the funnel-like *Cornell Model* [61], expressed as [62]:

$$V(r) = -\frac{\kappa}{r} + \sigma r. \tag{1.3}$$

This potential is effectively a combination of a Coulomb-like term representing the exchange of a single gluon between the quark-antiquark pair, given by the parameter κ and

a linear term expressing the confinement portion of the potential, expressed by the parameter σ , representing the string tension between the quark-antiquark pair. Figure 1.3 shows an illustration of the total quark-antiquark potential as a function of the charge separation distance; the separate contributions from the Coulomb-like and confinement terms are also shown.

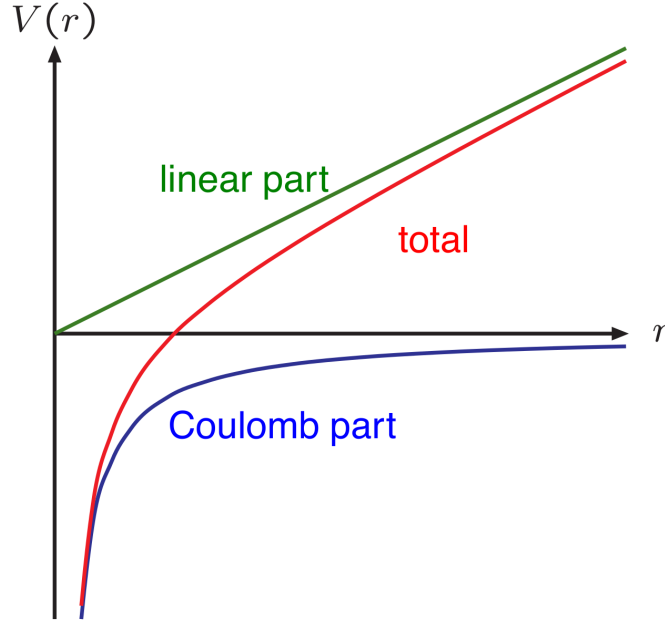


Figure 1.3: Static quark potential as a function of interquark separation. The linear contribution to the potential is shown in green while the Coulomb like contribution is shown in blue. The total effective Coloumb-linear potential is shown in red. Figure from [63].

The notion of quark confinement is in direct contrast with QED, where the creation of electron-positron pairs generates a dielectric property in the vacuum which in turn causes the strength of the electric charge to decrease at large distances. Figure 1.4 shows an illustrative depiction of the composition of color neutral (anti-)baryons and mesons via color and/or anti-color charges. This formalism of color neutral objects draws a heuristic from the manner in which a beam of white light can be generated from three overlapping beams of red, green, and blue light. Nevertheless, the reader is urged to exercise caution as quarks do not

possess color in the normal sense of the word [64]; the labelling exists simply to differentiate between the three color charges.

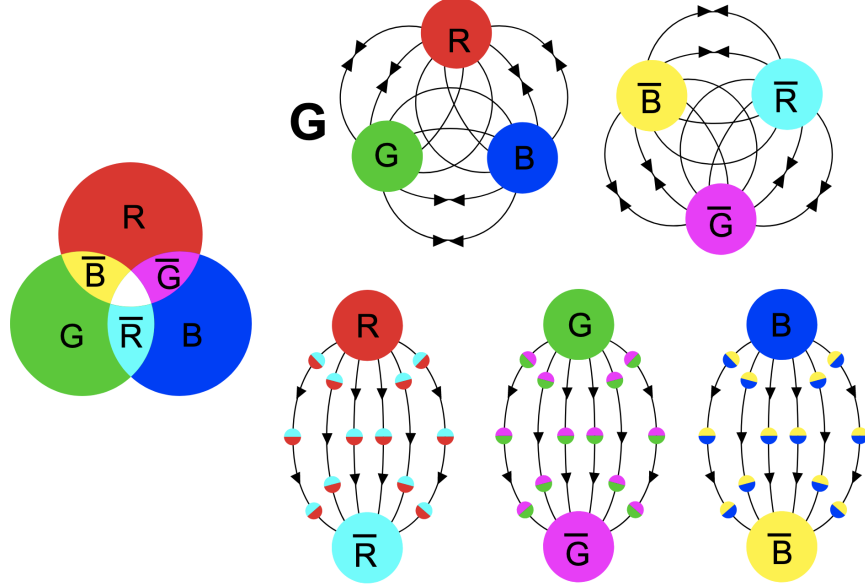


Figure 1.4: Color neutral combinations of quarks and anti-quarks via the color charge. Top: Three (anti-)quark states – (anti-)baryons. Bottom: Quark-antiquark pairs – mesons. The red, green, and blue color charges are colored accordingly and are labelled as R, G, and B, respectively. The anti-red, anti-green, and anti-blue anti-color charges are respectively colored in cyan, magenta, and yellow and are labelled as $\bar{R}$, $\bar{G}$, and $\bar{B}$, respectively. The QCD gluon field is shown in black and labelled as **G**. Figure from [65, 66]

1.2.2 Asymptotic Freedom

One of the most far-reaching contributions to QCD was made in 1973 by David Jonathan Gross, Hugh David Politzer and Frank Wilczek through the notion of *asymptotic freedom* [67, 68]. This phenomenon is typically expressed in the fundamental fermion representation by the coupling-dependent $\beta(g)$ function as [41]:

$$\beta(g) = -\frac{g^3}{(4\pi)^2} \left(\frac{11}{3}N - \frac{2}{3}n_f \right), \quad (1.4)$$

where the overall minus sign generally displays the asymptotic freedom of any non-abelian gauge theory of $SU(N)$ symmetry given a sufficiently small parameter n_f . Using Equation 1.4 and letting $\alpha = \frac{g^2}{4\pi}$, the value of the QCD running coupling constant α_s , at a given momentum scale Q^2 , can be expressed in terms of its value α at a different momentum scale Λ^2 in the following manner:

$$\alpha_s(Q^2) = \frac{\alpha}{1 + \frac{\alpha}{4\pi} \left(\frac{11}{3}N - \frac{2}{3}n_f \right) \log(Q^2/\Lambda^2)}. \quad (1.5)$$

As is the case in QCD, where $N = 3$, Equation 1.5 asymptotically approaches zero at a sufficiently large momentum scale for all values of $n_f \leq 16$ [69]. In this context, taking n_f to represent the total number of flavors in the theory, the aforementioned diminishing value of $\alpha_s(Q^2)$ at large Q^2 scales is expected. This phenomenon arises through an anti-screening effect generated by the non-abelian gauge bosons – in the case of QCD, the color charge-carrying gluons – opposite to the net screening from quarks and virtual gluon pairs. The former anti-screening effect has been shown to be an order of magnitude larger than the latter screening effect [70], enabling an amplification of the coupling constant at large distances. Conversely, this effect results in a reduction of the QCD coupling at either high energy scales or small separating distances, where the anti-screening effect diminishes.

Figure 1.5 shows various measurements of the QCD coupling, $\alpha_s(Q^2)$, as a function of the energy scale Q using particular perturbation orders in the different calculations, shown in the legend enclosed in parentheses. It is apparent that the value of $\alpha_s(Q^2)$ rapidly decreases as $Q \rightarrow \infty$ and increases quickly as $Q \rightarrow 0$. Thus, two separate regimes of the theory can be distinguished based on the value of the energy scale. For large values of Q , perturbative methods are valid, much like those used in QED. At smaller Q values, where

$\alpha_s(Q^2) \rightarrow 1$, perturbative calculations are no longer valid. In this non-perturbative regime, different method of calculation must be employed in order to adequately describe the theory.

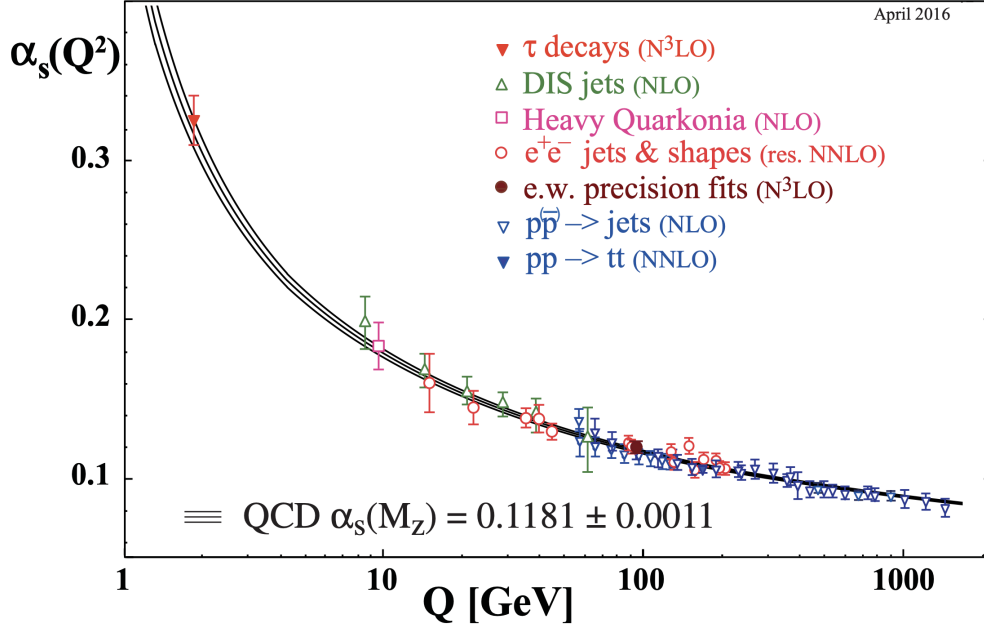


Figure 1.5: Coupling constant ($\alpha_s(Q^2)$) values measured as a function of different energy scales (Q) through different processes. Brackets indicate the respective order of QCD Perturbation Theory used to extract $\alpha_s(Q^2)$. The theoretical calculation of α_s is given by $\alpha_s(M_Z)$. For reference, NLO: next-to-leading order; NNLO: next-to-next-to leading order; res. NNLO: NNLO matched with resummed next-to-leading logs; N³LO: next-to-NNLO. Figure from [45].

Lattice QCD has become a well-established approach to dealing with such problems in the non-perturbative regime. Without loss of generality, Lattice QCD is a mathematically well-defined discretized formulation of QCD, rather than continuous, placed on a Euclidean grid (i.e. lattice) whose intersection points represent the quark fields. Each point is separated by a distance a , known as the *lattice spacing*, and is connected to its neighbors by the gluon fields. Usually, calculations are performed at finite lattice spacing and are extrapolated to the QCD continuum limit as $a \rightarrow 0$. In the last decade, Lattice QCD has proven its immense calculation power through various results such as the theoretical determination of the proton mass [71] as well as the prediction of the temperature range and order of the phase transition

from hadronic matter into deconfined quarks and gluons in the QCD phase diagram [72–74].

1.2.3 Chiral Symmetry

One of the global transformations that can occur in the QCD fermionic fields, the chiral transformation – also known as the pseudovector transformation – is responsible for rotating the chirality of a given particle. Chirality can typically be understood as the “handedness” of an object that may not be distinguished with a mirror image of itself. In the isospin notation, where $\psi = [u, d]$ for the light quark sector, the chiral transformation, $\mathcal{T}_\chi$, may be expressed in the following manner [75]:

$$\mathcal{T}_\chi : \psi \rightarrow e^{-i\gamma_5 \frac{\vec{\sigma}}{2} \vec{\vartheta}} \psi, \quad (1.6)$$

where γ_5 is the fifth Dirac matrix defined as $\gamma_5 \equiv i\gamma_0\gamma_1\gamma_2\gamma_3$, $\vec{\sigma}$ is the Pauli isospin vector and $\vec{\vartheta}$ represents the rotation angle of the transformation. The QCD Lagrangian is understood to be invariant under such a transformation, assuming vanishing quark masses – making it such that the pseudovector currents are *chirally symmetric*. Given that quarks are not massless, it is not the case that the pseudovector current is invariant under such a transformation and an explicit breaking of chiral symmetry occurs. This symmetry is understood to be approximately conserved in the light quark sector as the masses of these quarks are significantly lower relative to the QCD energy scale.

The chiral transformation can be used to describe the relationship between any two *chiral partners*, like a meson with total spin of 1 and odd parity (i.e. a *vector meson*), and a meson with total spin of 1 and even parity (i.e. a *pseudovector meson*). As a working example, the ρ^0 vector meson and a_1^0 pseudovector meson, respectively defined as $\vec{\rho}_\mu = \bar{\psi} \vec{\sigma} \gamma_\mu \psi$ and $\vec{a}_{1\mu} = \bar{\psi} \vec{\sigma} \gamma_\mu \gamma_5 \psi$ [76], can be rotated into one another through the chiral transformation as

follows:

$$\mathcal{T}_\chi \vec{\rho}_\mu = \vec{\rho}_\mu + \vec{\vartheta} \times \vec{a}_1. \quad (1.7)$$

Such relationship suggests the ρ^0 and a_1^0 mesons should in principle possess the same masses, given uniquely degenerate eigenvalues. In practice this is not the case; the neutral ρ^0 and a_1^0 meson masses have been experimentally measured as $775 \text{ MeV}/c^2$ and $1230 \text{ MeV}/c^2$, respectively [45]. Since the explicit breaking of chiral symmetry alone cannot explain such a mass difference resulting from non-zero quark masses, a spontaneous breaking of chiral symmetry is introduced [77]. This can be attributed to a non-zero vacuum expectation value of the chiral condensate, given by $|\langle \bar{\psi}\psi \rangle|$, which serves as an order parameter for the spontaneous breaking of chiral symmetry. In this formalism, for increasing values of baryon density and/or temperature, $|\langle \bar{\psi}\psi \rangle|$ approaches zero, where chiral symmetry is restored. In this restored phase the masses of chiral partners, like the ρ^0 and a_1^0 mesons, will degenerately converge. Figure 1.6 shows the temperature dependence of the spectral functions of both the ρ^0 vector meson and a_1^0 pseudovector meson ground states across a finite temperature range. Both spectral functions have different well-defined vacuum peak values, pertaining to the ground state masses of both mesons, which broaden and decrease as a function of increasing temperature. At temperature values above 150 MeV, the peaks begin to overlap and ultimately converge, suggesting a degeneracy in the masses of the chiral partners caused by a restoration of chiral symmetry above a specific temperature.

While Figure 1.6 does not explicitly show a temperature-dependent smooth transition into the chirally symmetric state, Figure 1.7 shows Lattice QCD results from the Wuppertal-Budapest Collaboration via the renormalized chiral condensate (left panel) and the subtracted chiral condensate (right panel) as a function of temperature, respectively given by $\langle \bar{\psi}\psi \rangle_R$ and $\Delta_{l,s}$ defined as [78]:

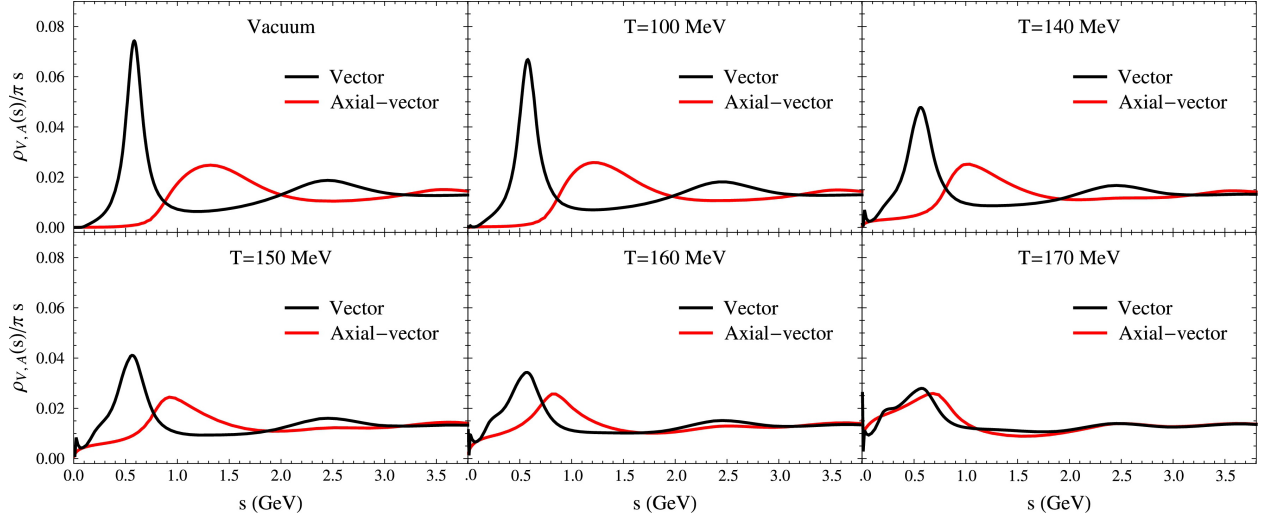


Figure 1.6: Spectral functions of vector and pseudovector meson ground states, respectively shown in black and red, in vacuum and at a finite temperature (T) values of 100, 140, 150, 160, and 170 MeV. The horizontal axis depicts the particle masses in GeV. Figure from [79].

$$\langle \bar{\psi}\psi \rangle_R = -[\langle \bar{\psi}\psi \rangle_{l,T} - \langle \bar{\psi}\psi \rangle_{l,0}] \frac{m_l}{X^4} \quad \text{for} \quad l = u, d \quad (1.8)$$

and

$$\Delta_{l,s} = \frac{\langle \bar{\psi}\psi \rangle_{l,T} - \frac{m_l}{m_s} \langle \bar{\psi}\psi \rangle_{s,T}}{\langle \bar{\psi}\psi \rangle_{l,0} - \frac{m_l}{m_s} \langle \bar{\psi}\psi \rangle_{s,0}} \quad \text{for} \quad l = u, d, \quad (1.9)$$

where the chiral condensate $\langle \bar{\psi}\psi \rangle_q$ for a given quark flavor $q = u, d, s$ of mass m_q is defined in terms of the partition function as:

$$\langle \bar{\psi}\psi \rangle_q = \frac{T}{V} \frac{\partial \ln Z}{\partial m_q}. \quad (1.10)$$

In the case of a real chiral phase transition, the chiral condensate serves as the corresponding order parameter. This is not the case with physical quark masses, as the phase transition is merely a cross-over. Nevertheless, the chiral condensate rapidly changes in the cross-over region and serves as an observable for such a transition. In Equation 1.8, the X term is any

quantity with a dimension of mass for which the pion mass, m_π , is used. Both panels in Figure 1.7 show a smooth chiral transition in this temperature region with inflection points centered about 150 – 170 MeV.

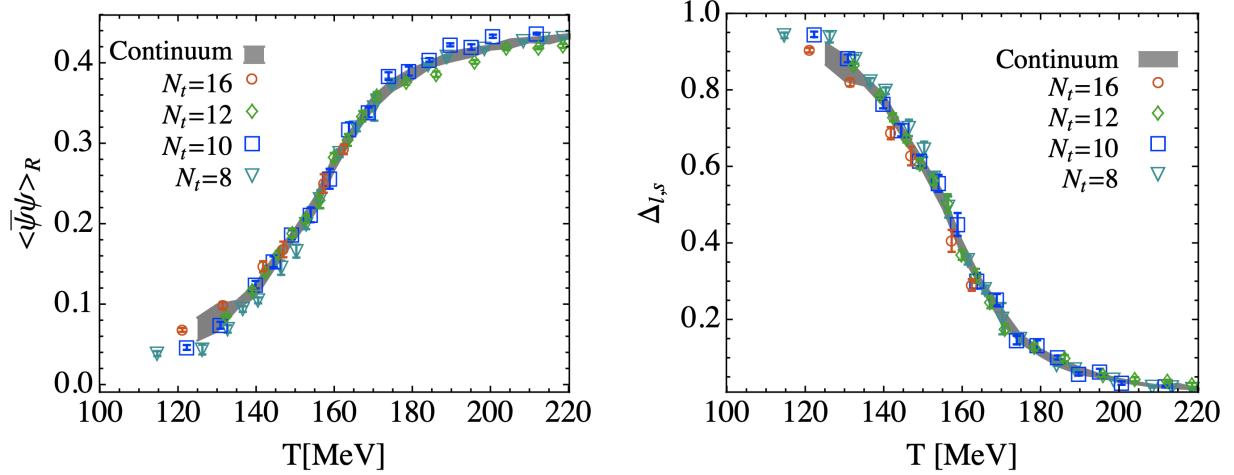


Figure 1.7: Left Panel: Renormalized Chiral Condensate ($\langle \bar{\psi}\psi \rangle_R$) as defined by Equation 1.8 for different temperature values. Right Panel: Subtracted Chiral Condensate ($\Delta_{l,s}$) as defined by Equation 1.9 as a function of temperature. The different symbols in each panel correspond to a different number of lattice sites (N_t) in the temporal dimension. The gray band represents the continuum extrapolation estimate. Figure from [78].

1.3 The QCD Phase Diagram and the Quark Gluon Plasma (QGP)

A direct consequence of asymptotic freedom is the notion that at high enough temperatures and/or densities, QCD matter transitions into a system of deconfined quarks and gluons – namely, the QGP. In this phase, where chiral symmetry is restored, quarks and gluons are the relevant degrees of freedom describing the dynamics of the system since they are no longer confined as hadronic matter [23, 24]. To understand this transition into a deconfined state of matter, a phase diagram of QCD matter can be constructed as a function of baryochemical potential (μ_B) and temperature. In principle, the value of μ_B represents the difference between the total number of baryons and anti-baryons in the system, namely, the net baryon density. Figure 1.8 shows an illustrative schematic of the QCD Phase Diagram

with various currently accepted characteristics labeled accordingly.

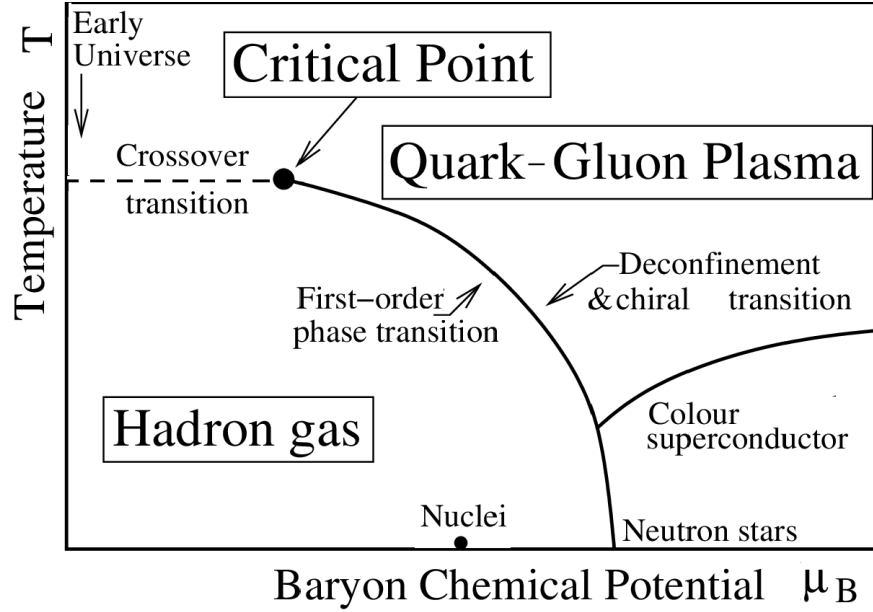


Figure 1.8: Schematic of the QCD Phase Diagram of matter as a function of temperature (T) and baryochemical potential (μ_B). Figure from [80].

The lower left-hand portion of 1.8, labeled as “hadron gas”, depicts the region of the diagram where QCD matter exists in a gaseous state of confined hadrons within the relevant T and μ_B ranges. As the temperature value $T \rightarrow 0$ MeV and μ_B increases, ordinary nuclear matter (labeled as “nuclei” in the figure) is formed around $\mu_B \approx 950$ MeV. In this low temperature range, as μ_B increases above the nuclear matter regime, a transition into a gaseous phase of degenerate neutrons takes place; conditions similar to those present in the cores of neutron stars. For higher μ_B values, the separating distance between neutrons diminishes and a gas of free roaming quarks forms and enters a predicted color superconducting regime [81]. On the other hand, for a vanishing baryochemical potential (i.e. $\mu_B \rightarrow 0$ MeV) when T increases to a high enough temperature, at approximately 150–170 MeV, the hadronic system undergoes a smooth transition [78] into the QGP phase. Around this temperature range, for increasing values of μ_B , the phase transition is considered to be of first order and suggests the existence

of a critical point in this vicinity [69]. This characterization of the QCD Phase Diagram is of utmost interest to the experimental heavy ion physics community and plays a crucial role in this dissertation.

As suggested by the QCD Phase Diagram, the recipe for creating a QGP consists of a system possessing a high enough temperature and/or baryon density. These two possible scenarios are independently sketched out in Figure 1.9. The top panel of the figure demonstrates QGP formation in a hadronic system at high temperatures through the progression of the temperature (given by T) as it approaches and exceeds a certain critical temperature (T_c); for temperature values $T > T_c$, the hadronic system “melts” into a state of deconfined quarks and gluons. Similarly, the bottom panel of the figure demonstrates QGP formation in a hadronic system at high density through the progression of the density (given by ρ) as it approaches and exceeds a certain critical density (ρ_c); for density values $\rho > \rho_c$, the hadrons in the system “dissolve” into a state of degenerate quark matter [33].

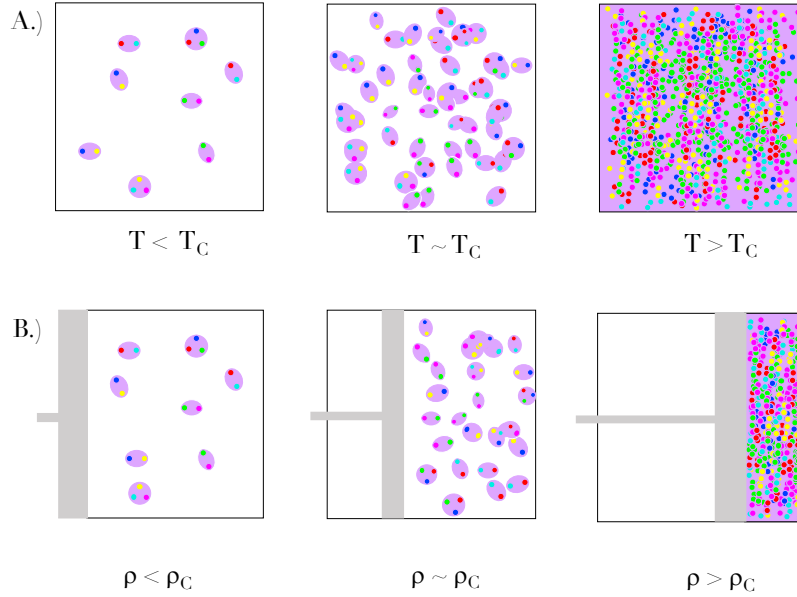


Figure 1.9: QGP formation at A.) high temperature (T) and B.) large baryon density (ρ). Figure adapted from [33].

The high temperature and high density recipes for QGP formation hint at three possible scenarios where QGP formation is to be expected. The first, which has extensively been described, is within the first microseconds after the Big Bang. The second would be in the center of astronomical objects, such as neutron stars. Lastly, the third is in the initial stages following a relativistic heavy ion collision. This dissertation focuses on the latter of the three, namely, using relativistic heavy ion collisions as the preferred method for experimentally observing the formation of the QGP as well as its transition back into hadronic matter. A simplified illustration of QGP formation via a heavy ion collision is shown in Figure 1.10 in two cases: (a) High temperature and vanishing baryochemical potential, where the colliding nuclei traverse through one another, and (b) high temperature and non-zero baryochemical potential, where the colliding nuclei stop one another. The former of these two cases is usually achieved when the center-of-mass collision energy per nucleon pairs ($\sqrt{s_{NN}}$) surpasses 100 GeV. It is the collective purpose of both the ALICE and STAR experiments to detect and measure traces of QGP formation at high temperatures for both vanishing and non-zero baryochemical potentials.

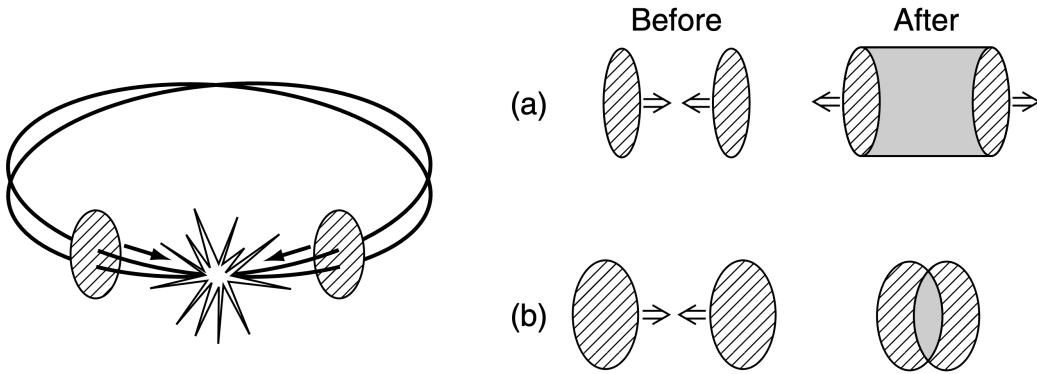


Figure 1.10: QGP formation via relativistic heavy ion collisions at (a) high temperature at large center of mass collision energies and (b) non-vanishing baryochemical potential (μ_B) at lower center of mass collision energies relative to (a). Figure from [33].

1.4 Heavy Ion Collisions

Relativistic Heavy Ion Collisions serve as a versatile tool for generating the conditions necessary for producing a QGP and investigating various aspects of the QCD Phase Diagram. The typical time scale of a heavy ion collision is of the order of $10 \text{ fm}/c$ ($\sim 10^{-23}$ seconds) with a corresponding length of 10 fm . Shortly after the initial collision, the generated *fireball* quickly expands and cools via radiative processes and transitions into a hadronic gas regardless of whether or not a QGP is formed. The evolution of a heavy ion collision is outlined by Figure 1.11 without (left) and with (right) the formation of the QGP.

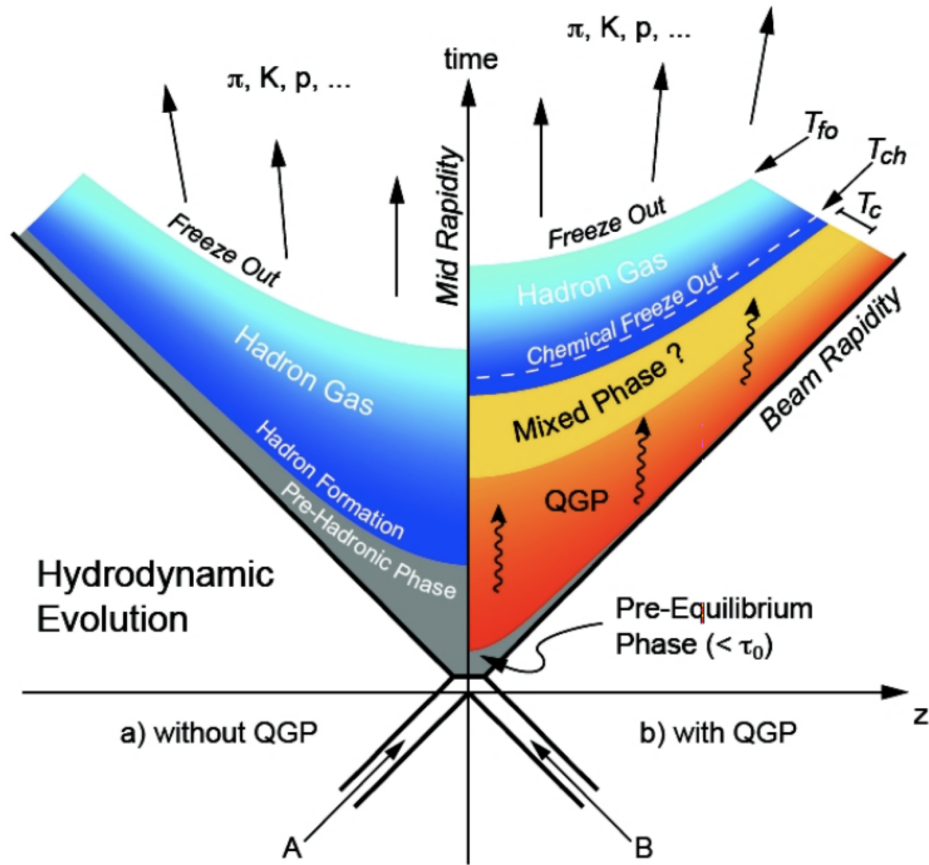


Figure 1.11: Depiction of the spacetime evolution of a heavy ion collision between ion beams A and B without (left) and with (right) QGP formation. Figure adapted from [82].

Generally speaking, heavy ions (e.g. Pb or Au nuclei) in a circular collider are accelerated

close to the speed of light and experience typical relativistic effects – the most noticeable being Lorentz contraction along the direction of motion, appearing as tortilla-like objects at the time of the collision. The two separate ion beams, one of which is accelerated in a clockwise manner and the other counter-clockwise, are injected into the interaction point where the experimental detector is situated and then collide into one another. As the ion beams intersect, the partons (i.e. quarks and gluons), which the colliding nuclei are composed of, undergo an initial hard scattering process. In the case where the fireball temperature exceeds the critical temperature, the out-of-equilibrium particles produced by the initial hard scattering interact with each other until thermal equilibrium is reached and a QGP is formed. The QGP system rapidly cools down below the critical temperature and hadronizes into a system of interacting hadrons and hadronic resonances. The hadrons and resonances in the expanding system continue interacting inelastically until the fireball reaches a threshold temperature, known as the *chemical freeze-out* temperature (T_{ch}), below which the produced hadronic species in the thermal bath and their relative yields remain unchanged. As the hadronic gas expands, the hadrons continue interacting elastically until the gas reaches the *kinetic freeze-out* temperature (T_{kin}); at this point the dynamical quantities of particles, such as transverse momentum (p_{T}), are fixed. After this point, the hadrons leave the system as free streaming particles and are then detected by the experimental apparatus.

1.4.1 Experimental QGP Signatures

Due to its short-lived lifetime, QGP formation cannot be directly measured experimentally in a heavy ion collision. Nevertheless, there exists a general consensus of the creation of a deconfined state of quarks and gluons in relativistic heavy ion collisions based on various experimental observables serving as potential signatures of the QGP [83]. This procedure is analogous to the methods used to probe the early Universe through remnants, such as the

cosmic microwave background and the abundance of chemical elements [33]. Some of the reigning experimental QGP signatures are delineated below.

Collective Anisotropic Flow: Hydrodynamic simulations of the QGP have made the assertion of it behaving as a near perfect fluid, with very low viscosity [84]. Flow as an observable provides an insight of key aspects of the QGP and its evolution following a heavy ion collision. As various particles are produced in a heavy ion collision, their respective energy and momentum distributions can be characterized through such an observable. Anisotropic flow measurements from collisions from off-center nuclei show evidence of a non-uniform spatial distribution of the produced medium in heavy ion collisions relative to the beam line. These effects between particles can be analyzed through the azimuthal distribution relative to the collision symmetry plane (Ψ_n) given by:

$$\frac{dN}{d\varphi} = \frac{1}{2\pi} \left(1 + 2 \sum_n \nu_n \cos(n(\varphi - \Psi_n)) \right), \quad (1.11)$$

where ν_n are the n^{th} harmonics of the distribution and φ represents the azimuthal angle. Each flow coefficient ν_n represents a specific class of collective motion. The viscosity of the produced medium can be determined from the ν_2 and compared to the values acquired from viscous hydrodynamics simulations.

Jet Quenching: During the initial hard partonic scattering following a heavy ion collision, high energy collimated particle beams, namely *jets*, are produced. In the case where a QGP is formed, as the outgoing jets strongly interact with the generated hot and dense medium, they experience a reduction in their overall energy mostly through

in-medium radiative and collisional energy losses. Hence, the measurement and reconstruction of jets produced in heavy ion collisions give insight to the properties of the QGP (e.g. temperature, volume and color density) based on the amount of *quenching* they undergo as they traverse the medium.

Heavy Quarkonia Production: Another resulting byproduct of the initial hard scattering from colliding heavy ions is the production of heavy quark-anti-quark pairs. In the presence of a QGP, where unbound color charges are freely available, the potential generated between heavy quarks is effectively screened, causing the heavy quark bound states to dissociate into the medium. This in turn would lead to a suppression in the production of heavy vector mesons. While this suppression effect is present at lower collision energies, it is not so prevalent at higher energies, where the heavy quarkonia have been observed to *regenerate*. Given that no heavy quarks are present prior to the initial collision, the production rate of heavy quarkonia should in principle scale as a function of the collision energy by a simple mass argument. Hence, for high enough energies, an increased probability of two independently produced heavy quarks hadronizing into a bound state should in principle emerge.

Photon Production: As the produced QGP from a heavy ion collision undergoes its evolution back into a confined hadronic state, it will generally expand and cool via radiative processes. Through these processes, the QGP will typically emit thermally generated photons which in turn produce dileptons (e.g. electron-positron pairs). Since leptons do not possess a color charge, they do not interact via the strong nuclear force and are free to traverse through the QGP medium unperturbed. In turn, the dilepton pairs can be experimentally detected and serve as an excellent probe for QGP

formation in high energy heavy ion collisions through the determination of the initial QGP temperature via thermal photon production.

Strangeness Enhancement: The enhanced production of strange quarks is historically one of the earliest predictions of the presence of the QGP resulting from colliding heavy nuclei. As is the case with heavy quarks, strange quarks do not constitute the initially colliding nuclei in heavy ion collisions; in principle, any strange quark should in principle be produced by the initial hard scattering of the collision. Nevertheless, given that the mass of a strange-anti-strange quark pair ($m_{s\bar{s}} \sim 186 \text{ MeV}/c^2$) is slightly above the transition temperature of the QGP, the production of such pairs should increase as a function of collision energy assuming they are thermally produced in the QGP medium.

Although the aforementioned signatures are all necessary and proper observables for warranting the presence of a deconfined state of quarks and gluons in a heavy ion collision, this dissertation focuses exclusively on the lattermost. This is choice addressed and justified accordingly in the body of this work.

2 SEQUENTIAL STRANGENESS FREEZE-OUT

This Chapter serves as the foundation for the underlying purpose of this dissertation, introducing the concept of *Sequential Strangeness Freeze-out* from a theoretical and phenomenological perspective.

In order to produce said foundation, I will first outline a brief historical account leading to the assertion of Gell-Mann’s quark model, which introduces the existence of the strangeness quantum number. I will touch upon the mathematical formulation of the so-called *Eightfold Way* as the means to classify the observed particles produced by cosmic rays and particle collisions. I will then recapitulate on strangeness enhancement as a signature of QGP production in heavy ion collisions as well as its consequences on strange particle production. I conclude the chapter by providing evidence of a potential flavor hierarchy at freeze-out in the crossover region of the QCD phase diagram; here, I suggest the use of statistical hadronization models as the means of investigating such phenomenon.

2.1 A Brief History of Strangeness

By the first half of the 20th century, a multitude of experimental efforts confirmed the presence of cascading ionized particles emanating from extraterrestrial sources, herein referred to as *cosmic rays*, penetrating and interacting with the Earth’s atmosphere [85]. Cosmic rays are typically high-energy protons and/or helium nuclei originating from the either the Sun or beyond the Solar System. Upon entering the atmosphere, cosmic rays in turn collide with atmospheric air molecules and produce a cascading shower of secondary particles, namely, hadrons, leptons, and/or electromagnetic radiation [86].

By the early 1950’s, observations from cloud chamber experiments suggested that the hadrons produced in cosmic air showers could possess significantly different masses spanning over an order of magnitude. A posteriori, these particles are now understood to be distinct

types of mesons and/or baryons composed of constituent (anti)quarks of specific flavors, e.g. pions, kaons, and so-called *hyperons* [87] like the Λ and Ξ baryons [88–94] – which at the time were simply referred to as *V-particles* given the topology of their decay tracks [95]. Soon after, some of the first collider experiments performed at particle accelerators, such as the Cosmotron at BNL and the Bevatron at LBNL, verified the existence of such hadronic entities [96,97], and facilitated phenomenological studies regarding their physical properties. Despite the presence of these findings, no generalized method existed which could organize experimentally observed hadrons systematically, as no distinction between elementary and composite particles was in place. A major leap in the development of a descriptive theory of hadronic entities came from the introduction of the hypercharge (Y), as a function of a particle’s average electric charge (Q), its elementary charge (e), and its third isospin component (I_3), given by [95]:

$$Y = 2 \left(\frac{Q}{e} - I_3 \right). \quad (2.1)$$

Table 2.1 shows the different known values of Q and I_3 for the most commonly known charge multiplets derived from the fundamental charge symmetry – namely, the nucleon doublet (neutron and proton), the pion triplet (π^- , π^0 and π^+) and the Δ resonances (Δ^- , Δ^0 , Δ^+ and Δ^{++}), which carry an angular momentum value of $J = 3/2$.

Table 2.1: Electric charge (Q) and isospin component (I_3) of the nucleon doublet, the pion triplet, and the Δ resonances.

	n	p	π^-	π^0	π^+	Δ^-	Δ^0	Δ^+	Δ^{++}
I_3	$-1/2$	$1/2$	-1	0	1	$-3/2$	$-1/2$	$1/2$	$3/2$
Q	0	e	$-e$	0	e	$-e$	0	e	$2e$

Using Equation 2.1, the Y values for the particles in Table 2.1 are readily attained as:

$$Y = 1 \quad \text{for} \quad n, p, \quad (2.2)$$

$$Y = 0 \quad \text{for} \quad \pi^-, \pi^0, \pi^+, \quad (2.3)$$

$$Y = 1 \quad \text{for} \quad \Delta^-, \Delta^0, \Delta^+, \Delta^{++}. \quad (2.4)$$

Hence, for the above cases, the Y value of each particle is identical to that of its baryon number (B), namely, $Y = B$. However, this is clearly not the case for the Λ nor the Ξ baryons. Given that the Λ baryon appears as a singlet state, namely as Λ^0 , assigning it a $I_3 = 0$ value by definition, renders the following, as per Equation 2.1:

$$Y = 0 \quad \text{for} \quad \Lambda^0. \quad (2.5)$$

In the case of the Ξ baryon, being a Ξ^0 and Ξ^- doublet, with $I_3 = 1/2$ and $I_3 = -1/2$ values, respectively, Equation 2.1 renders:

$$Y = -1 \quad \text{for} \quad \Xi^0, \Xi^-. \quad (2.6)$$

It is evident for both Equations 2.5 and 2.6, $Y \neq B$, given that both the Λ and Ξ baryons possess a $B = 1$ value, by definition. To remedy this, an additional quantum number, S , was introduced, defined as the difference between Y and B , namely:

$$S = Y - B. \quad (2.7)$$

While Nishijima referred to the number S as the “ η - charge” [98], Gell-Mann bluntly resorted to calling it *strangeness* [99]. Historically, the latter name persisted and thus is used accordingly for the remainder of this dissertation.

2.1.1 The Eightfold Way

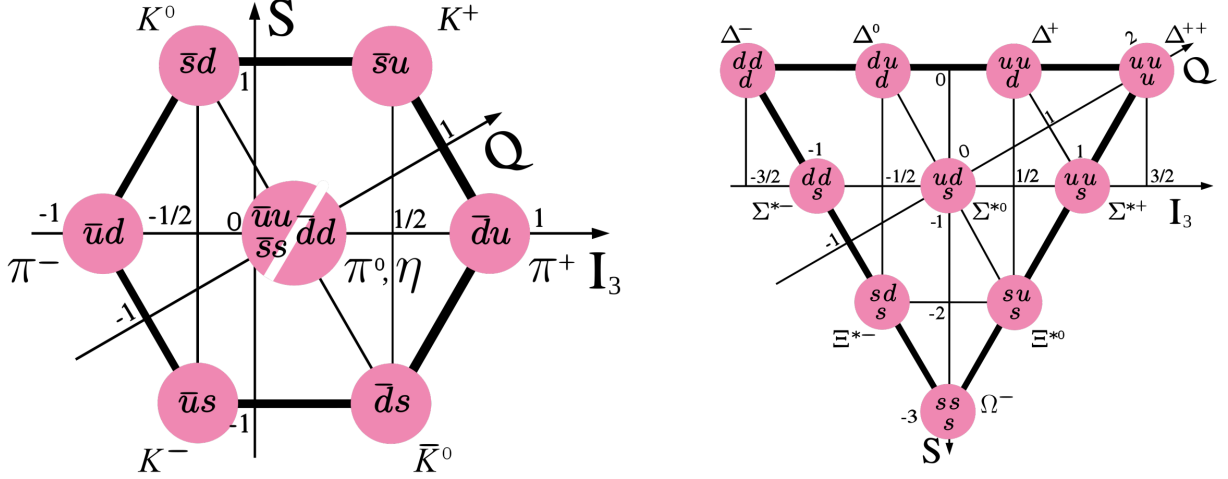


Figure 2.1: Meson octet (left) and baryon decuplet (right) as inspired by The Eightfold Way. For both diagrams, particles along the same diagonal line share the same elementary charge (Q), while all particles along the same vertical line possess identical values of the third isospin component (I_3). In this same manner, particles residing within the same horizontal line share indistinguishable strangeness content (S). As is convention, strange quarks carry values of $S = -1$ while anti-strange quarks have values of $S = 1$.

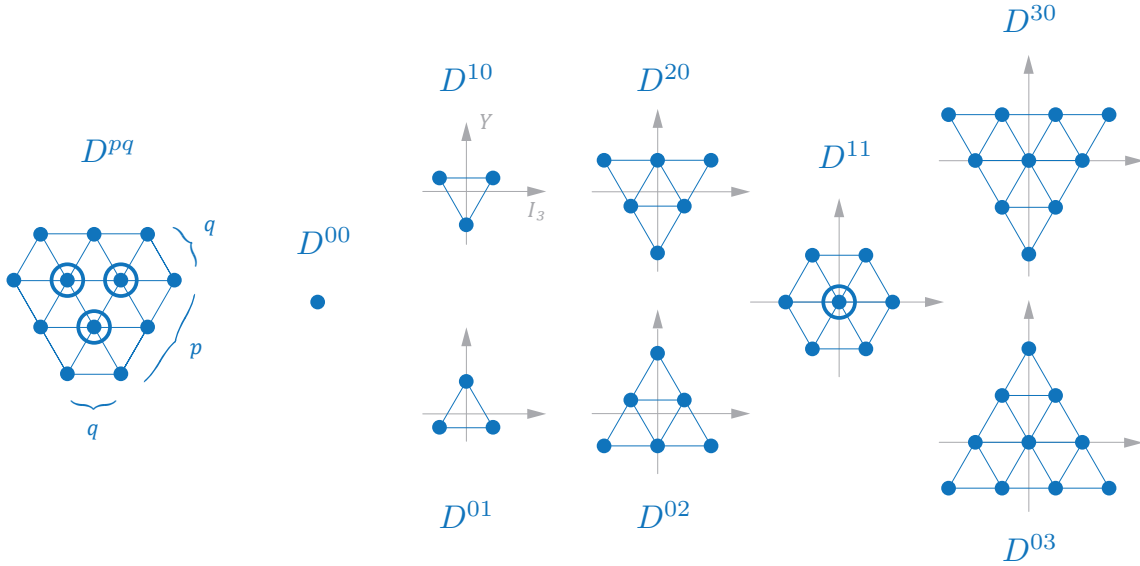


Figure 2.2: Generalized multiplets of the SU(3) group showing the generic construction of its irreducible representation, given by D^{pq} followed by its lowest-dimensional mathematical representations. The construction of the meson octet and baryon decuplet from Figure 2.1 are derived from D^{11} and D^{30} , respectively. Figure from [100].

The identification of S as a fundamental quantum number made it such that hadronic entities measured by cosmic ray experiments and early bubble chamber experiments could be neatly organized in a systematic fashion by the organizational scheme known as *The Eightfold Way* – named by Gell-Mann after the Noble Eightfold Path of Buddhism [101]. The Eightfold Way posited the existence of flavored quarks as elementary particles which in turn can be grouped to form mesons and baryons (Figure 2.1) based on the lowest-dimensional irreducible representations of the $SU(3)$ Lie Group (Figure 2.2).

A major prediction put forward by The Eightfold Way was the existence of a spin $3/2$ baryon consisting of three strange quarks, namely an sss state, known as the Ω^- baryon. It was in 1964 that the existence of the Ω^- baryon was confirmed in a bubble chamber experiment at Brookhaven National Lab from the collision of a negatively charged kaon beam onto a proton from the hydrogen gas in the chamber (Figure 2.3).

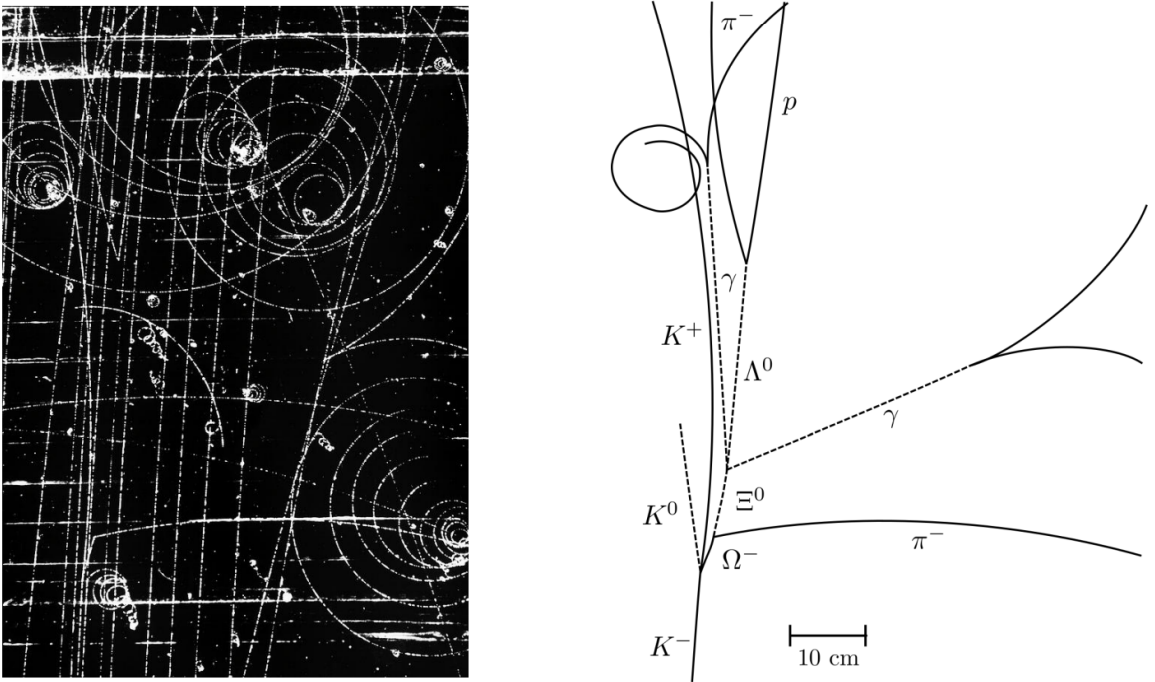
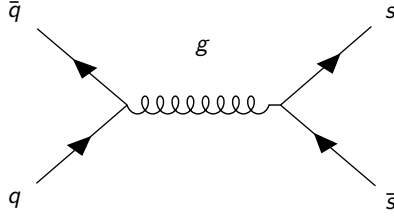


Figure 2.3: Left: Bubble chamber photograph of the Ω^- discovery at Brookhaven National Lab in 1964. Right: Line diagram of the Ω^- event decay photographed on the left-hand side of the figure. Figure adapted from [102].

The discovery of the Ω^- baryon provided the missing piece of baryon decuplet, as the other particles had already been experimentally verified, and paved the way for the validation of the quark model and the early developments of QCD.

2.2 Strangeness Enhancement as a Signature of the QGP

A.)



B.)

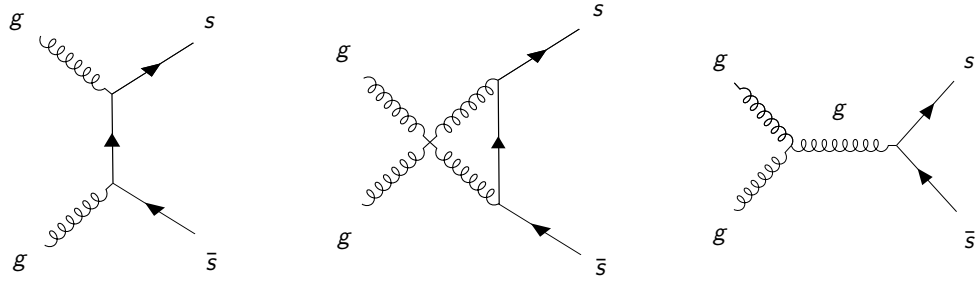


Figure 2.4: Lowest order QCD Feynman Diagrams representing the production of $s\bar{s}$ pairs for A.) $q\bar{q} \rightarrow s\bar{s}$ and B.) $gg \rightarrow s\bar{s}$ processes. For each diagram, time flows from left to right. Figure adapted from [34].

The enhanced production of strange quarks as a signature of the production of a QGP in relativistic heavy ion collision was first predicted by Rafelski [34, 35] due to the abundant production of $s\bar{s}$ pairs dominated by the fusion of gluon pairs freely available in the QGP medium as opposed to the $q\bar{q} \rightarrow s\bar{s}$ process (Figure 2.4) shown to contribute to less than 10% of the total production rate. Given that the mass of the $s\bar{s}$ pair sits just above the characteristic pseudocritical temperature, a sharp increase in the production of strange hadrons would suggest the presence of the QGP, as the QGP medium would thermally generate an abundant amount of additional $s\bar{s}$ pairs. This would in turn result in the enhanced production of

final state strange and multi-strange baryons (i.e. Λ , Ξ and Ω). Strangeness enhancement has been observed across various heavy ion experiments and has also been shown to persist in high-multiplicity proton-proton collisions [36]. Measurements of the abundant production of strange hadrons in heavy ion and elementary particle collisions provide an accessible method of investigating the reigning mechanisms leading to said production – particularly from a phenomenological point of view questioning the characteristic transition temperature from quark to hadronic degrees of freedom across particles of different species (i.e. different flavors).

2.3 Sequential Strangeness Freeze-out

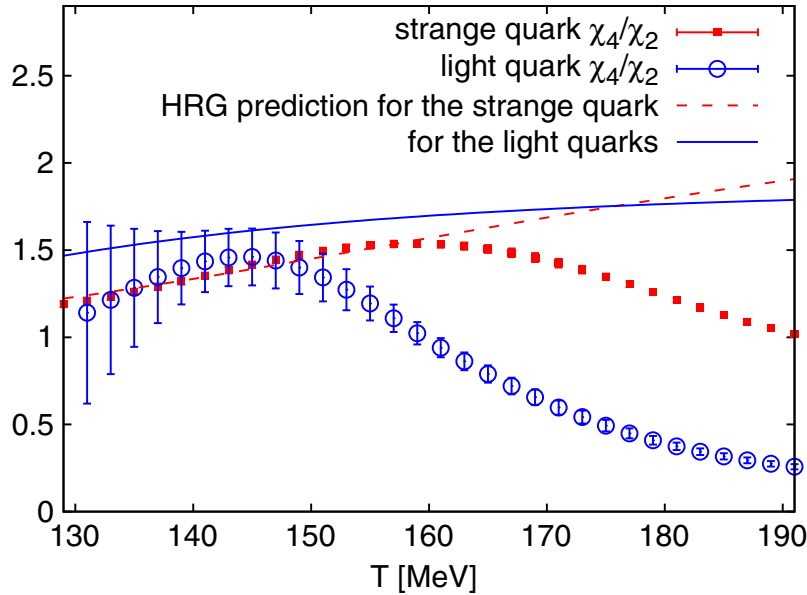


Figure 2.5: Temperature dependence of χ_4/χ_2 for light and strange quarks in the continuum limit, shown in blue open circles and red closed squares. The lattice data are also compared to their respective flavor-dependent HRG calculations, shown as a dashed red line for the strange quark and a solid blue line for the light quarks. Figure from [39].

Continuum extrapolated susceptibility calculations of single flavor quantum numbers on the

lattice [39, 103] show a significant difference in the determined freeze-out temperatures between flavors in the crossover region of the QCD phase diagram. In terms of thermodynamic quantities, susceptibilities can be calculated as the partial derivatives of the pressure (p) with respect to changes to the components of the chemical potential as:

$$\chi_{ijk}^{BSQ} = \frac{\partial^{i+j+k} p / T^4}{\partial(\mu_B/T)^i \partial(\mu_S/T)^j \partial(\mu_Q/T)^k}, \quad (2.8)$$

where i , j , and k are the pertaining orders of the susceptibility, B , S , and Q are the baryon number, strangeness content, and electric charge, respectively, and T is the temperature of the system. In terms of the partition function, Z , the susceptibilities take on the following form:

$$\chi_{ijk}^{BSQ} = \frac{1}{VT^3} \frac{\partial^{i+j+k} (\ln Z)}{\partial(\mu_B/T)^i \partial(\mu_S/T)^j \partial(\mu_Q/T)^k}, \quad (2.9)$$

The susceptibilities can in turn be related directly to experimentally measured distribution moments of conserved charges – i.e. the mean (M), variance (σ^2), skewness (S), and kurtosis (κ):

$$M = VT^3 \chi_1, \quad (2.10)$$

$$\sigma^2 = VT^3 \chi_2, \quad (2.11)$$

$$S = \frac{VT^3 \chi_3}{\chi_2^{3/2}}, \quad (2.12)$$

$$\kappa = \frac{VT^3 \chi_4}{\chi_2^2}. \quad (2.13)$$

In order to account for the system volume dependence, ratios of the susceptibility values are often employed:

$$\frac{\chi_2}{\chi_1} = \frac{\sigma^2}{M}, \quad (2.14)$$

$$\frac{\chi_3}{\chi_2} = S\sigma, \quad (2.15)$$

$$\frac{\chi_3}{\chi_1} = \frac{S\sigma^3}{M}, \quad (2.16)$$

$$\frac{\chi_4}{\chi_2} = \kappa\sigma^2. \quad (2.17)$$

While the even-odd ratios of the susceptibilities serve as a good observable for the determination of μ_B , even-even susceptibility ratios – to first order – are determined solely by the temperature at freeze-out and thus serve as an excellent observable for directly determining said freeze-out temperature [104].

A comparison of the susceptibility ratios χ_4/χ_2 in a flavor specific manner shows a deviation of the lattice and Hadron Resonance Gas (HRG) model calculations which coincides with the peaks of the lattice data (Figure 2.5). This deviation, which occurs at flavor-specific temperatures, differs by 15 – 20 MeV from light to strange quarks [39].

Additionally, SHM calculations from thermal fits to experimental data show a potential difference in the freeze-out temperature in the cross-over region of the QCD phase diagram may be present. Figure 2.6 shows three different thermal model calculations of experimental particle yields (dN/dy) produced in the most central (0 - 10%) PbPb collisions at ALICE at $\sqrt{s_{NN}} = 2.76$ TeV assuming a single universal temperature at chemical freeze-out. While the three models show exceptional agreement with the experimental dN/dy values over six orders orders of magnitude, the bottom panel of the Figure shows an inherent tension in the fit between the (anti-)protons and the multi-strange (anti-)baryons – where a difference of about 5σ is present between the two model calculations. This tension has been shown to increase to a difference close to 8σ in the single temperature thermal model fits to the ALICE

data performed in the most central (0 - 10%) PbPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [105].

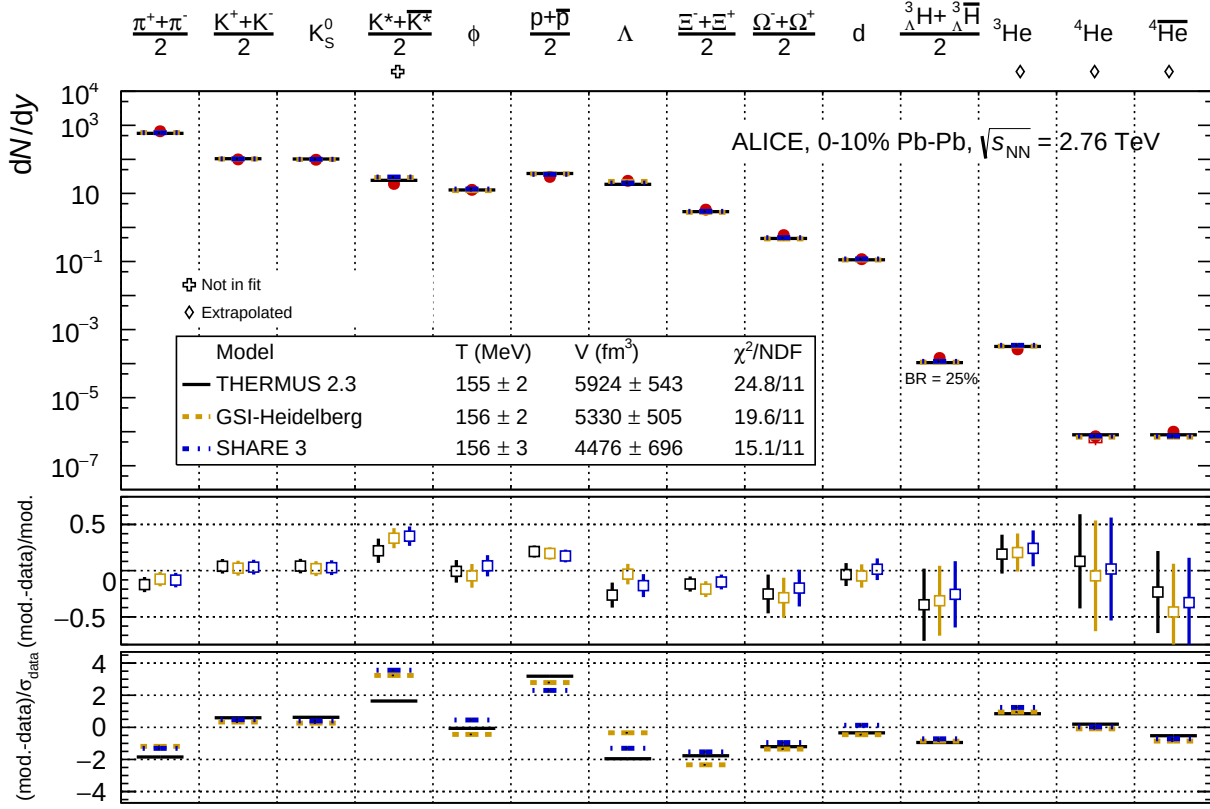


Figure 2.6: Top Panel: Thermal model calculations of experimental particle yields (dN/dy) produced in the most central (0 - 10%) PbPb collisions at ALICE at $\sqrt{s_{\text{NN}}} = 2.76$ TeV for three different models – THERMUS, GSI-Heidelberg, and SHARE 3 – respectively shown as solid black lines, dash-dotted yellow lines and dash-dotted blue lines. The experimental data is shown as solid red circles. The respective extracted chemical freeze-out temperatures (T), volumes per unit rapidity and goodness of fit measures for each model are enclosed in the box located in the lower left-hand corner. The middle panel shows the relative differences between each model yield calculation and the measured experimental yields, and normalizes them to the model yield calculations. The bottom panel depicts the same relative differences as the middle panel but instead normalizes them to the experimental errors. Figure from [106].

The theoretical and experimental results presented above give rise to the possibility of a flavor hierarchy in the extracted chemical freeze-out temperature in the cross-over region of the QCD phase diagram and provide grounds for the exploitation of the SHM framework

in order to probe said possibility. It is this very concept that generates the foundation for the main analysis presented in this dissertation. A simplified schematic of the spacetime evolution of a heavy ion collision following the creation of a QGP in which a flavor-dependent chemical freeze-out sequentially occurs is shown in Figure 2.7. In this schema, strange and light flavor hadrons reach chemical freeze-out at different temperatures, respectively given as T_S and T_L , which correspond to different points in time in the evolution of a heavy ion collision.

Following such schema, this dissertation will show further evidence of sequential hadronization in both an energy dependent and collision system size dependent manner via the SHM framework using final state experimental yields of produced particles measured by both ALICE and STAR.

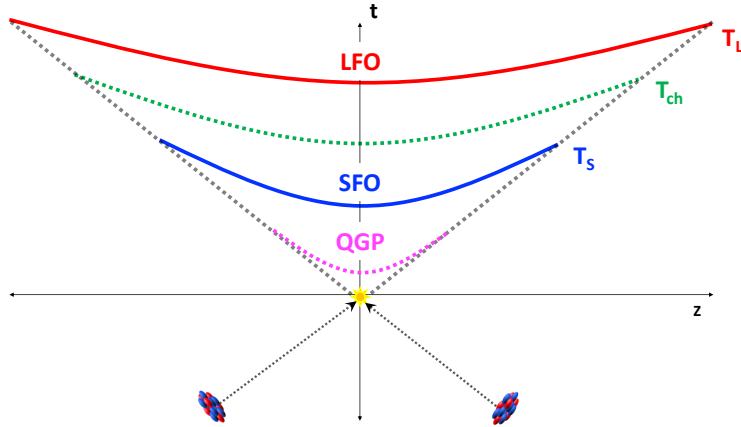


Figure 2.7: Modified depiction of the spacetime evolution of a heavy ion collision showing two separate temporal points at which flavor specific freeze-outs occur. The solid blue line indicates the point at which the strange flavor freeze-out (SFO) occurs, while the solid red line shows the light flavor freeze-out (LFO); the former has a corresponding freeze-out temperature T_S and the latter has the temperature T_L .

3 A LARGE ION COLLIDER EXPERIMENT

In this chapter, a brief overview of the LHC facility at CERN is presented. The ALICE detector is then described in detail, focusing on the subdetectors primarily used for the study of strange hadrons generated by the pp, pPb and PbPb collision systems at the LHC. The final section of the chapter draws similarities between the ALICE detector and supplementary experiments located at different laboratory facilities, namely the STAR detector at RHIC, and points out the benefits from comparing the results from both experiments from a phenomenological point of view.

3.1 The Large Hadron Collider

The LHC [107] is an accelerator complex built in its present form from 1998 to 2008 at CERN, situated in the French-Swiss border. It stands as the world's largest circular collider, consisting of an underground ring 27 km in circumference. It is comprised of superconducting magnets and additional structures used to boost and accelerate particles to speeds close to the speed of light. The accelerated particle beams are injected into one of two beam pipes running in opposite directions (one clockwise and the other counterclockwise) along the circumference of the main LHC ring and are then allowed to collide at four main interaction points. Four major experiments are located at each of these interaction points and are used to study various structures of fundamental matter accessible via the high energy particle collisions at the LHC.

The four main experiments at the LHC are the Compact Muon Solenoid (CMS), A Toroidal LHC Apparatus (ATLAS), LHC beauty (LHCb) and A Large Heavy Ion Collider Experiment (ALICE). Both CMS and ATLAS are general purpose detectors whose main experimental objective is to detect the Higgs boson via pp collisions and to probe novel physics beyond our current understanding of the Standard Model. The LHCb experiment is

primarily designed to investigate the decays of particles containing bottom (b) quarks and to elucidate the asymmetry between matter and antimatter. The ALICE detector is designed specifically for the purpose of understanding the structure of strongly interacting matter, namely the QGP, generated by heavy ion collisions.

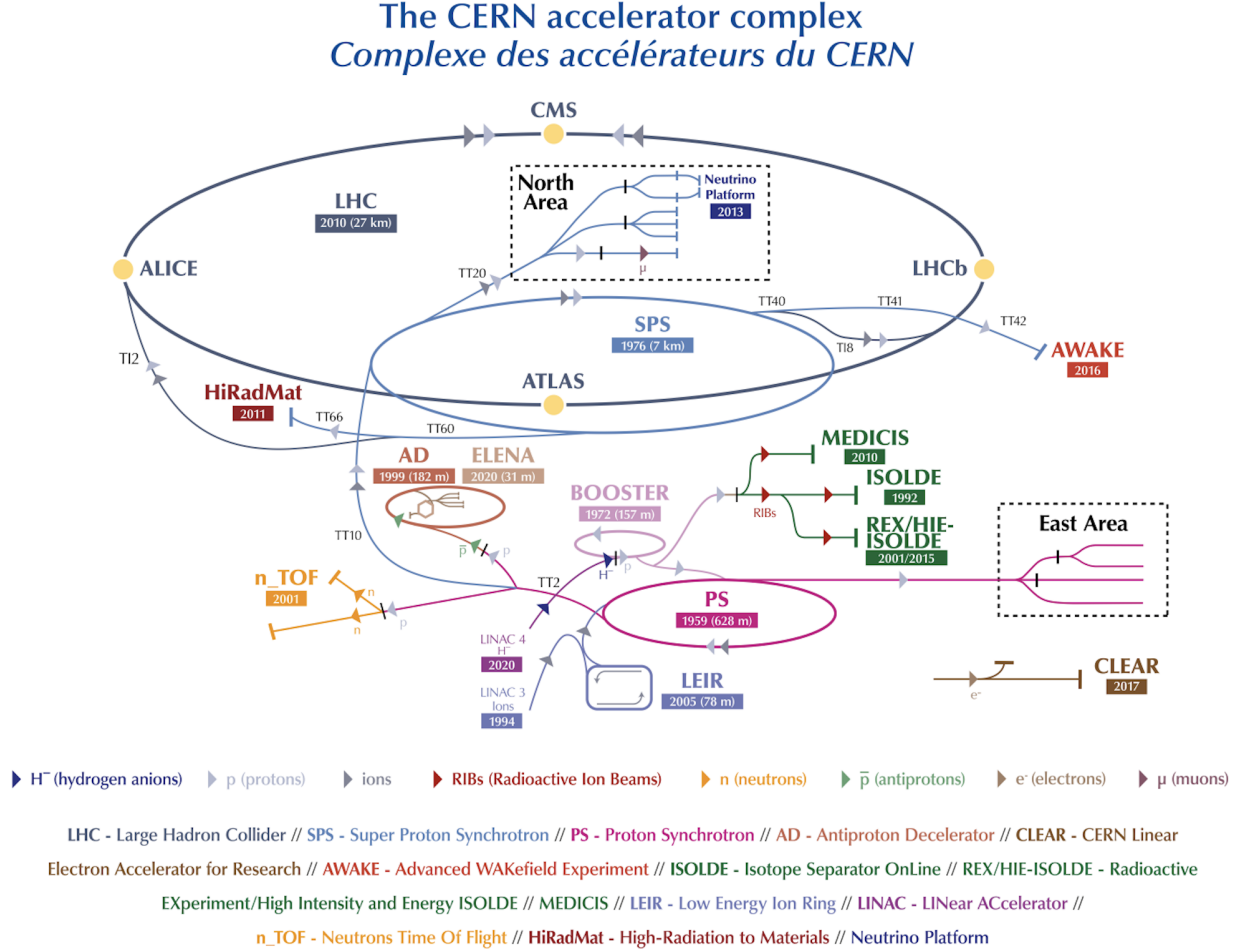


Figure 3.1: Schematic of the Large Hadron Collider complex at CERN depicting the accelerators, beam lines, and a selection of major detector experiments. The ALICE detector, located at Interaction Point 2, is shown on the left and labeled with a closed yellow circle. Proton sources are first accelerated by the Linear Accelerator 4 (Linac4), injected into the Proton Synchrotron Booster (PSB), further accelerated by the Proton Synchrotron (PS) and the Super Proton Synchrotron (SPS) before finally being injected into the two LHC beam pipes. On the other hand, heavy ions start at the Linac3 and are then collected by the Low Energy Ion Ring (LEIR) before following the same acceleration process as the protons. Figure from [108].

Figure 3.1 shows an overview of the LHC facility and depicts the acceleration chain of particles – both protons and heavy ions – as they are injected into the LHC beam pipes, where they are then allowed to interact at points where the experimental detectors are situated. By design, the maximum center-of-mass energy ($\sqrt{s}$) achieved by the LHC for proton on proton (pp) and lead on lead (PbPb) collisions are respectively 14 TeV and 5.5 TeV.

The Linear Accelerator 4 (Linac4), being the proton beam source at the LHC, accelerates negative hydrogen ions to 160 MeV prior to them being injected into the Proton Synchrotron Booster (PSB). During the injection of hydrogen ions from the Linac4 to the PSB, the ions have their two electrons stripped and remain simply as protons. The protons are then accelerated to 2 GeV and are prepared for further injection into the Proton Synchrotron (PS). The PS further accelerates the proton beams to 26 GeV and then relays them to the Super Proton Synchrotron (SPS). Here, the protons are then accelerated to 450 GeV. Lastly, the beams are injected into one of the two LHC beam pipes – one traveling in a clockwise fashion while the other circulates counterclockwise. Here the beams are further accelerated and finally reach their maximum energy of 6.5 TeV. The two beams are allowed to collide in the interaction points enclosed by CMS, ATLAS, LHCb and ALICE, where their total center-of-mass collision energy reaches 13 TeV. On the other hand, Pb ion do not originate from the Linac4, but rather the Linear Accelerator 3 (Linac3). A pure lead source is heated to a temperature around 800°C, causing it to vaporize. The lead vapor is ionized to a charged state and accelerated through a carbon foil, where it is then collected and further accelerated by the Low Energy Ion Ring (LEIR). The lead ions then follow the same acceleration process as the protons and then reach their final collision energy of 5.02 TeV per nucleon pair ($\sqrt{s_{\text{NN}}}$).

3.2 A Large Ion Collider Experiment (ALICE)

The ALICE detector [109] is one of four major experiments at the LHC, hosted by CERN, whose primary purpose is specifically to study the production of strongly interacting matter generated in heavy-ion collisions, where the collision temperatures and energy densities most likely generate a QGP. While the main purpose of the ALICE detector is to measure and analyze heavy ion collisions, in particular PbPb collisions, the detector is also capable of recording and analyzing data emerging from collisions of proton on lead nuclei (pPb) as well as pp collisions. The latter two are typically seen as providing a baseline reference for the events (i.e. collisions) generated by the larger collision system (i.e. PbPb collisions) and provide a broader picture emerging from the experimental data.

THE ALICE DETECTOR

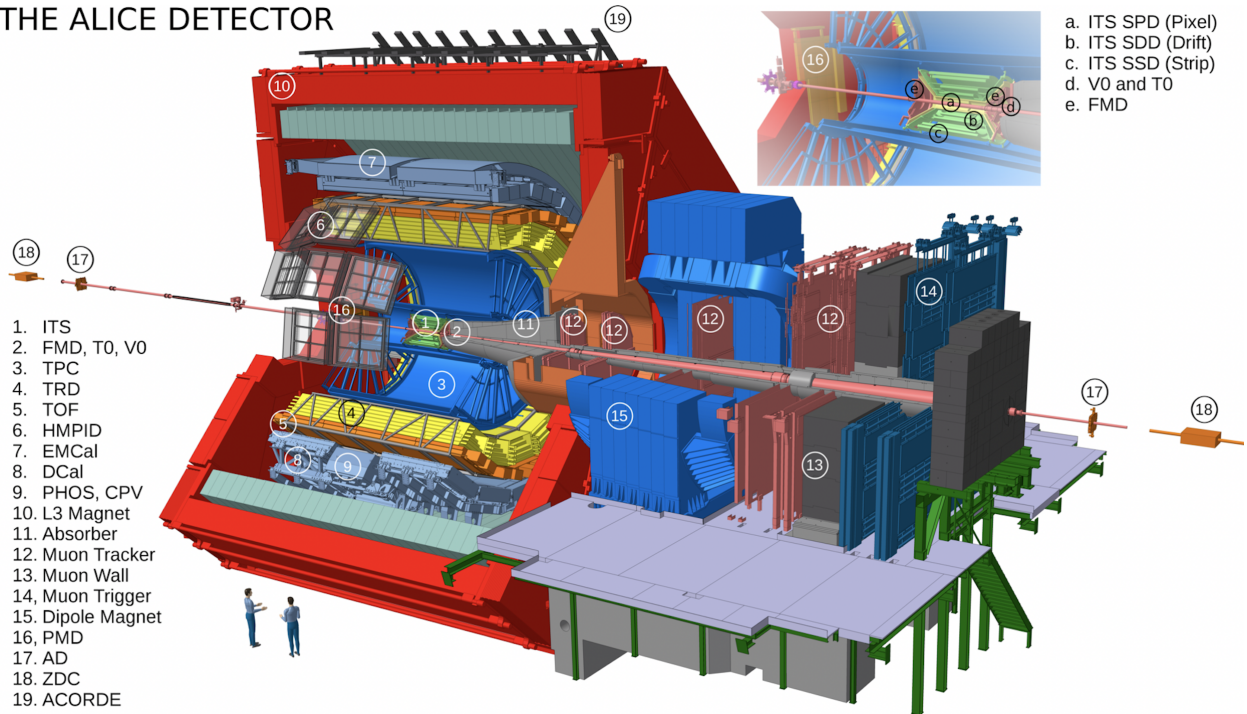


Figure 3.2: Schematic of the ALICE detector with a cutaway view of its inner detector components. Each sub-detector is labeled and listed accordingly on the left-hand side of the figure. The top right-hand sub-figure focuses in on the Inner Tracking System (ITS) and the detectors around its vicinity. Figure from [110].

The ALICE detector, measuring $26 \times 16 \times 26 \text{ m}^3$ and weighing 10 million kilograms, is composed of 19 primary subdetector components stratified into one of two separate groups: The Forward Muon Arm and the Central Barrel. Figure [110] shows a schematic of the ALICE detector with a cutaway view of its inner components. The Forward Muon Arm can be seen on the righthand side of the figure, while the Central Barrel is shown on the left.

The Forward Muon Arm detectors are specifically used for the study of heavy quarkonia through to their dimuon decay channel topology and cover the $2.5 < \eta < 4$ pseudorapidity range, where the pseudorapidity η is defined by:

$$\eta = -\ln \left[\tan \left(\frac{\theta}{2} \right) \right]. \quad (3.1)$$

In the above equation, θ denotes the polar angle between the momentum of a particle and the positive direction of the beam pipe. By convention, θ increases in a counterclockwise manner relative to the beam pipe. Given that heavy quarkonia measurements are outside of the scope of this dissertation, a detailed emphasis on the Forward Muon Arm subdetectors is omitted. Further details on these subdetectors can be found in [109].

On the other hand the Central Barrel subdetectors – enclosed by a large solenoid magnet whose uniform magnetic field along the direction of the beam reaches values up to 0.5 T – cover the $|\eta| < 0.9$ pseudorapidity range. The main detectors within the central barrel cover the full azimuthal angle φ and are primarily used for the measurement of strange hadron production. These detectors are the Inner Tracking System (ITS), the Time Projection Chamber (TPC), the Time of Flight (TOF) detector and the V0 Scintillators. These four detectors are described in detail following this section. Other subdetectors within the Central Barrel, such as the High Momentum Particle Identification Detector (HMPID), the Transition Radiation Detector (TRD), the Photon Spectrometer (PHOS) and the Electromagnetic Calorimeter (EMCal), are not further emphasized, given that they are not used in

the measurements of interest for this dissertation (i.e. strange hadrons). For further details on these subdetectors, the reader is directed to [109].

Charged particle tracks traversing the bulk of the detector are reconstructed from track points in the ITS and TPC via a Kalman Filter statistical estimation technique [111, 112], which splits the volume of the TPC into 36 trapezoidal sectors. The filter groups three hits within each sector – approximately linearly arranged in a straight line – and links them together to other triplet groups. The trajectory of the particles is then extrapolated through the TPC volume by merging track segments from individual sectors which are then fitted into a final track. The p_T of a track is calculated based on the charge of the track (q), the radius of curvature of the track (R) perpendicular to the magnetic field and the strength of the magnetic field B by:

$$p_T = qRB. \tag{3.2}$$

The time of the collision is measured by additional detectors for both low and high multiplicity events. For low multiplicity events, the T0 detector is used. Without loss of generality, the T0 detector can achieve a time precision of 25 ps and is composed of an array of 12 Cherenkov counters. Conversely, for high multiplicity events, the TOF detector is used. Lastly, the V0 detector – described in Section 3.2.4 – is responsible for the rejection of possible background events as well as for the Minimum Bias (MB) trigger based on simultaneous signals from both the V0-A and V0-C detectors.

3.2.1 Inner Tracking System (ITS)

The ALICE Inner Tracking System is responsible for the precise determination and reconstruction of primary vertices [109, 113]. Its resolution is capable of reconstructing tracks down to 100 μm and possesses the ability to reconstruct secondary vertices from long lived

particles. In addition, the ITS is capable of identifying low transverse momentum (p_T) particles down to $p_T < 0.2$ GeV/c and collectively covers the $|\eta| < 0.9$ pseudorapidity range. Surrounding the central LHC beam pipe, its inner and outer layers have radii of 4.0 cm and 43.6 cm, respectively. The ITS is composed of six silicon layers, each comprised of three different types of silicon technologies. The two inner ITS layers are Silicon Pixel Detectors (SPD), the third and fourth layers are Silicon Drift Detectors (SDD), and the outer two layers are Silicon Strip Detectors (SSD). A schematic of the ITS is depicted in Figure 3.3, showing the positioning of the SPD, SDD and SSD relative to the LHC beam pipe. Unlike the SDD and SSD, the SPD possesses an extended pseudorapidity coverage range of $|\eta| < 1.98$ and provides a continuous coverage for the measurement of charged particle multiplicity along with the Forward Multiplicity Detector (FMD).

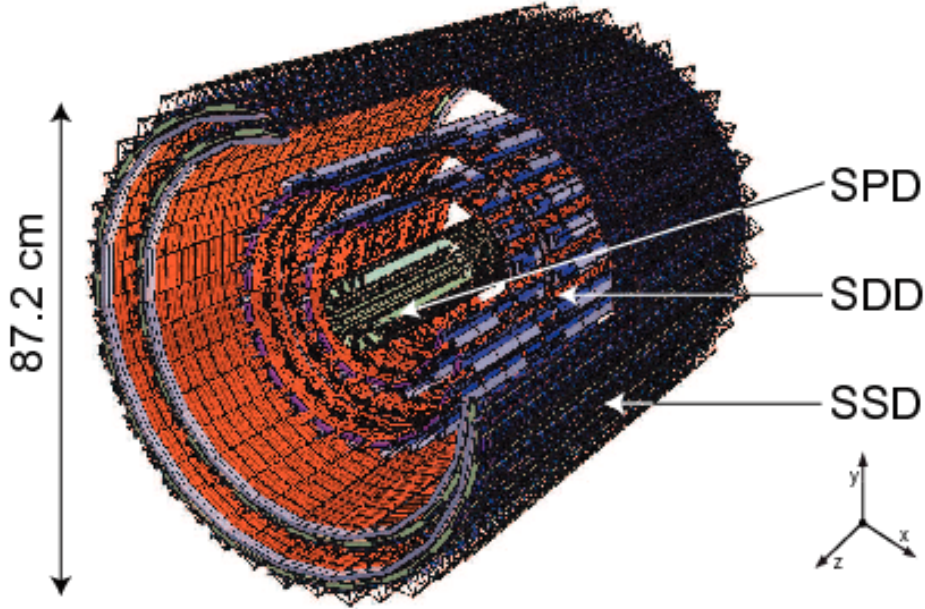


Figure 3.3: Schematic of the ALICE Inner Tracking System (ITS) depicting the three different varieties of silicon detectors composing the ITS. Starting from the inside, closest to the beam pipe, are two layers of Silicon Pixel Detectors (SPD). In the middle are two layers of Silicon Drift Detectors (SDD), while the outermost two layers are Silicon Strip Detectors (SSD). Figure from [114].

The SPD is composed of hybrid silicon pixels with dimensions of $425\ \mu\text{m} \times 50\ \mu\text{m}$ and contains 9.8×10^6 cells. The track density around the vicinity of the SPD is about $50\ \text{cm}^2$. On the other hand, the particle density around the SDD region is around $7\ \text{cm}^2$. The SSD is responsible for joining the ITS tracks to the TPC. The four outer layers (SDD and SSD) have analogue readout, which enables both tracking and low- p_T PID through the specific energy loss (dE/dx) of a passing track. The tracking resolution for low- p_T tracks is governed by multiple scattering interactions between the bulk of the detector material and the crossing track; for this reason, the detector material budget is kept to an absolute minimum. Both the SSD and SDD layers are thus required to have a minimum thickness of $300\ \mu\text{m}$ in order to minimize the detector signal to noise ratio.

3.2.2 Time Projection Chamber (TPC)

The TPC is the main detector of the ALICE Central Barrel, providing exceptional vertex reconstruction and PID as well as tracking charged particles, and provides up to 160 spatial points for every track [109, 115]. It is capable of resolving the transverse momentum of a particle from the curvature of its track inside the magnetic field over a wide p_T range ($0.1 < p_T < 100\ \text{GeV}/c$). The ALICE TPC is a large cylinder which envelops the ITS and measures 510 cm in length and 250 cm in diameter, also covering the $|\eta| < 0.9$ pseudo-rapidity range. Initially a 90-10-5 Ne-CO₂-N₂ gas mixture filled the TPC field cage, and was later changed to 90-10 Ne-CO₂. Prior to the start of Run 2 in 2015, the gas mixture was changed to 88-12 Ar-CO₂ for operational stability. The TPC field cage, containing two potential dividers as well as a central high voltage (100 kV) electrode, generates a uniform electric field of 400 V/cm which spawns a $2.7\ \text{cm}/\mu\text{s}$ drift velocity and results in a $100\ \mu\text{s}$ maximum drift time. The TPC field cage is surrounded by an additional concentric cylindrical containment vessel generating a CO₂ filled gap. Both ends of the TPC cylinder are

sealed by end plates containing a series of 18 radially segmented Multiwire Proportional Chambers (MWPCs) split into two separate Outer (OROC) and Inner Readout Chambers (IROC) – totalling 36 readout chambers at each end cap. Figure 3.4 shows a schematic of the ALICE TPC with a cutout depicting the inside of the bisected outer field cage; the wired readout chambers can be seen at both end caps from both inside and outside the TPC.

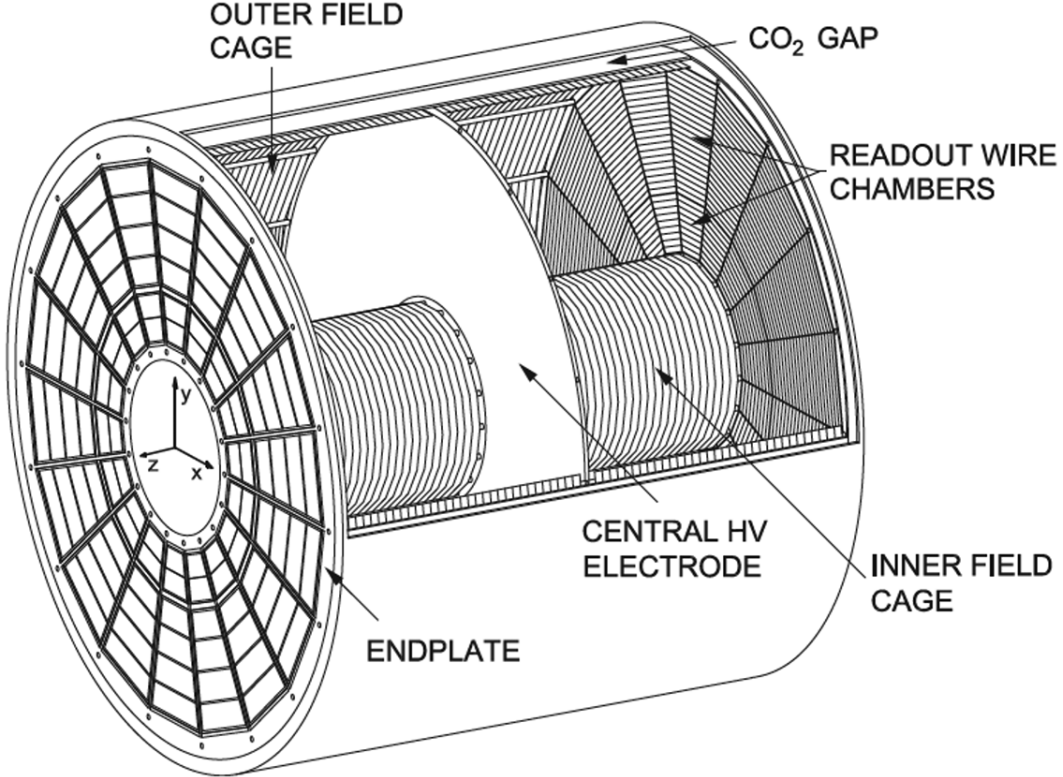


Figure 3.4: Schematic of the ALICE Time Projection Chamber (TPC) depicting the bisected field cage and the 18 sectors in both end plates, each consisting of 36 readout chambers. Figure from [116].

While the radial range of the IROCs is from 84.8 cm to 132.0 cm, the radial range of the OROCs is from 134.6 cm to 246.6. Each readout chamber contains 560,000 readout pads of size $4.0 \text{ mm} \times 7.5 \text{ mm}$ for the IROCs, and $6 \text{ mm} \times 10 - 15 \text{ mm}$ for the OROCs. The track density at which MWPCs can operate is limited by the pile-up effects generated by the relative motion of positive ions which produce a long tailed current signal of about 50

μs . On the front end side of the chamber are rows of up to 25 Front End Cards (FECs), each composed of 128 channels; within every channel, the signal of the deposited current on the pads traverses through a charge sensitive amplifier, a 25 MHz analog-to-digital converter and then a shortening filter which performs a baseline subtraction on the signal and cancels the aforementioned tail. A Readout Control Unit (RCU) couples each row of FECs to the Detector Control System, the Data Acquisition (DAQ) System and the global event trigger.

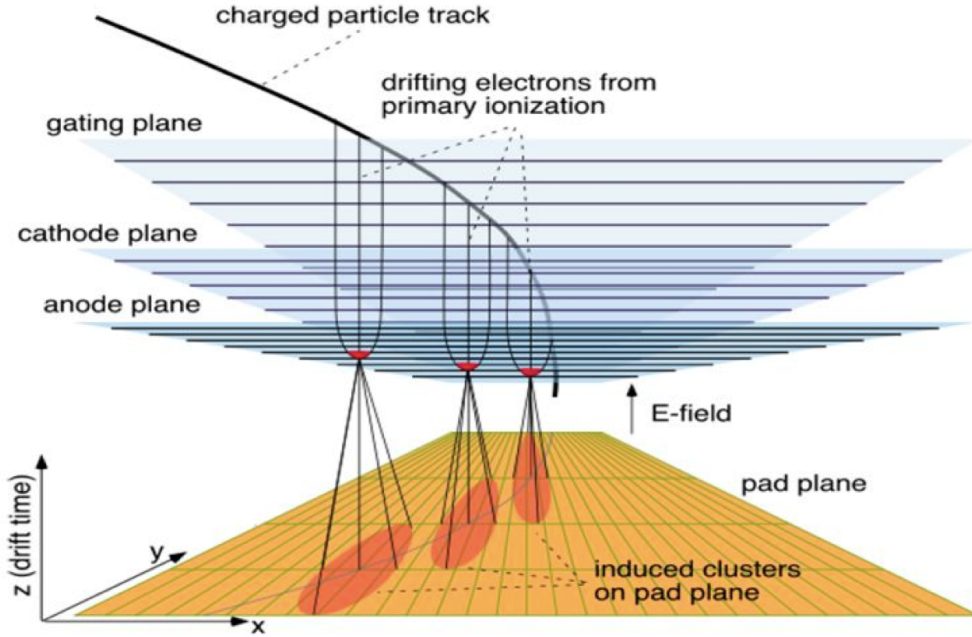


Figure 3.5: Working principle of the ALICE Multiwire Proportional Chamber (MWPC) TPC showing (from the top down) the internal gating grid, cathode and anode wires, and pad readout plane. Figure from [117].

The working principle behind the ALICE TPC is rather straight forward. The charged particles generated in a collision ionize the inert gas inside the bulk of the detector as they traverse through it. Due to the applied electric field in the detector, the released electrons then drift towards the end plates of the TPC. In the MWPC setup, a series of wires – running azimuthally perpendicular to the beam axis – is set-up leading to the pad readout plane; their order from the inside-out is: Gating wire grid, cathode wire grid and anode wire

grid. Due to a large electric field gradient, the drifting electrons generate an amplification avalanche near the vicinity of the anode wires. This results in the production of additional secondary ions and electrons which are proportional to the number of primary drift electrons and ionizing energy loss of the initial traversing charged particle. The resulting electron avalanche continues making its way towards the pad plane, where their induced clusters are detected. The cathode wire is responsible for collecting the ions produced in the amplification avalanche. This in turn provides a two dimensional signal along the spatial transverse plane. The gating grid can be turned from the open to closed position by alternating its voltage between ± 100 V. When the gate is in the closed position, the gate wires collect the drifting electrons resulting from the primary ionization of the gas and prevent them from further traversing toward the amplification region. In this closed position, the gating grid also prevents ions generated in the electron avalanche from prior events from traveling back into the TPC drift volume. The gating grid only opens in allowable intervals matching the maximum electron drift time of $90 \mu\text{s}$. The two-dimensional position of the charged particle track can then be measured. An additional spatial dimension can be extracted given that both the drift time and velocity are known a priori. The entire electron avalanche process is illustrated by Figure 3.5.

In order to accurately identify the species of a charged particle traversing through the bulk of the TPC, both the dE/dx and p_T of the particle's track are measured. In order to describe the mean dE/dx value of a charged relativistic track, the Bethe-Bloch function is typically used to provide a bona fide parameterization of known charged particles. This function is given as:

$$-\left\langle \frac{dE}{dx} \right\rangle = K z^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\frac{1}{2} \ln \frac{2m_e c^2 \beta^2 \gamma^2 T_{max}}{I^2} - \beta^2 - \frac{\delta(\beta\gamma)}{2} \right], \quad (3.3)$$

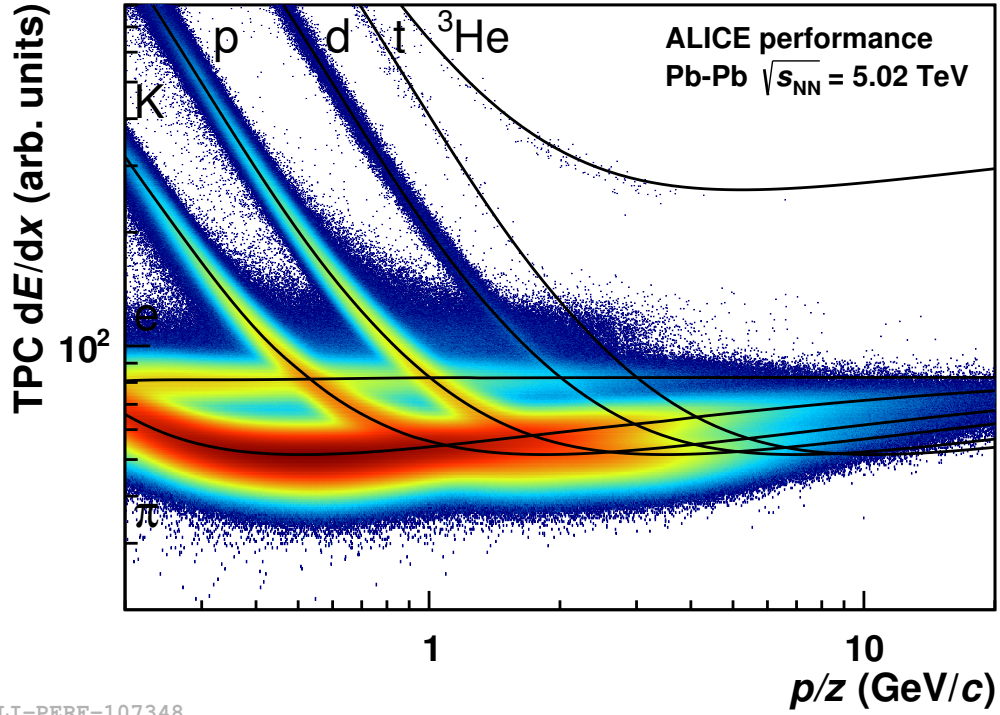


Figure 3.6: Specific ionization energy loss (dE/dx) for positive tracks in the TPC as a function of rigidity (p/z) for ALICE PbPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The solid black lines represent the parameterized Bethe-Bloch function values for each identified particle. Figure from [110].

where K is a constant, such that $K = 4\pi N_A r_e^2 m_e c^2 = 0.307075 \text{ MeV mol}^{-1} \text{ cm}^2$, z is the charge of the incident particle, A and Z are respectively the atomic number and atomic mass of the medium, $\beta = v/c$, $\gamma = (1 - \beta^2)^{-1/2}$, m_e is the mass of the electron, c is the speed of light, $T_{max} = 2m_e c^2 \beta^2 \gamma / [1 + (2\gamma m_e / M) + (m_e / M)^2]$ (where M is the Mass of the incident particle), I is the mean excitation energy of the medium and $\delta(\beta\gamma)$ is the density effect correction to ionization energy loss. Figure 3.6 illustrates the measurement of the distribution of energy loss as a function of the particle rigidity (p/z) in PbPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV for a series of identified positively charged particles by the ALICE TPC. The solid black lines represent the parameterized Bethe-Bloch function values for each identified particle.

3.2.3 Time of Flight Detector (TOF)

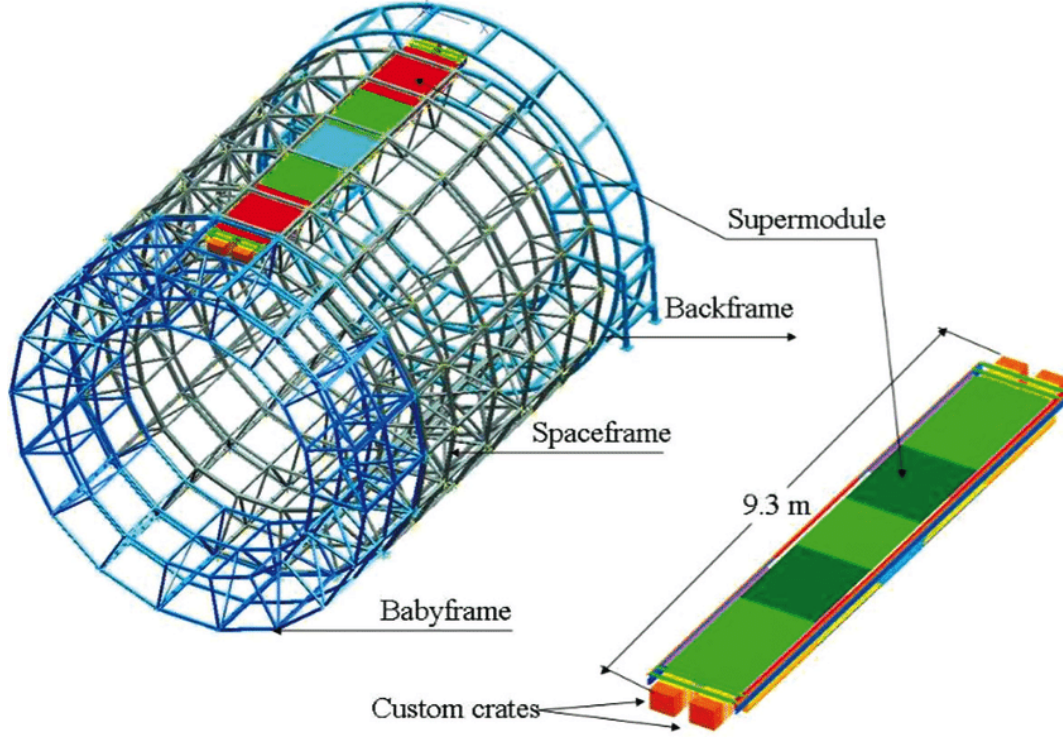


Figure 3.7: Schematic of one of 18 supermodules (right) of the ALICE Time of Flight (TOF) detector inside the ALICE space frame (left). Figure from [118].

The ALICE TOF detector is a Multigap Resistive Plate Chamber (MRPC) detector built as a cylindrical shell with outer and inner radii of 399 cm and 370 cm, respectively [109]. It is responsible for charged particle PID via charged particle track velocity determination in the intermediate transverse momentum range ($0.5 < p_T < 2.5$ GeV/c). The TOF detector is composed of two stacks consisting of five separate gas gaps – filled with a mixture of 90% $C_2H_2F_4$, 5% $i-C_4H_{10}$, and 5% SF_6 – along the z direction and 87 modules grouped into 18 super modules in the azimuthal direction, resulting in a pseudorapidity acceptance range of $|\eta| < 0.9$. Figure 3.7 shows a schematic of the ALICE TOF detector space frame as well as a single supermodule.

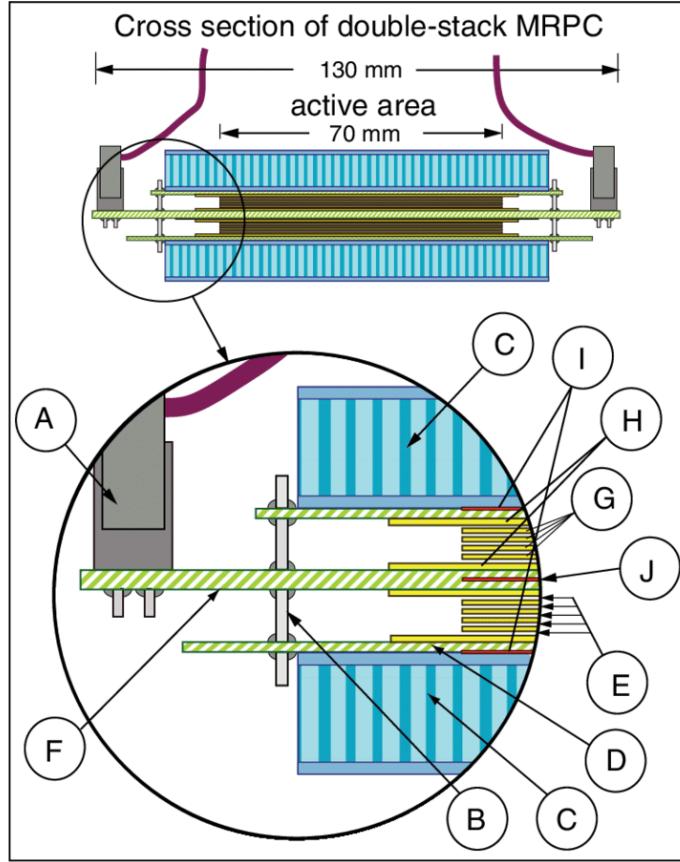


Figure 3.8: Double-stack Multigap Resistive Plate Chamber (MRPC) in the ALICE Time of Flight (TOF) detector. (A) Flat cable connector relaying differential signal from the MRPC to the front-end electronics. (B) Pins responsible for bringing signal from the cathode pick-up pads to central printed circuit board. (C) One centimeter thick honeycomb panel. (D) Printed circuit board containing cathode pickup pads. (E) Gas gaps of a $250\mu\text{m}$ width. (F) Central printed circuit board with anode pick-up pads. (G) Intermediate glass plates $400\mu\text{m}$ in thickness. (H) Resistive paint coated $550\mu\text{m}$ thick outer glass plates with outer surface coated with resistive paint. (I) Cathode pick-up pad; (J) Anode pick-up pad. Figure from [119].

A large and highly uniform electric field is applied across each module. This guarantees that any electron produced via the ionization process will immediately avalanche with negligible drift time. The sum of all the signals across every gap produced by the electrons and ions in the avalanche is equal to the induced signal on the pickup electrodes. This makes it such that there is no late tail in the time response signal. The time resolution of the ALICE TOF has been demonstrated to be below 50 ps and an efficiency close to 100% has been

shown. A detailed schematic of the ALICE TOF MRPC double stack is shown in Figure 3.8.

Unlike the PID determination via dE/dx measurements used by the ITS and the TPC, the TOF PID technique uses the mass and charge of a traversing particle in order to identify said particle. While the charge of a particle is directly measured, the mass itself is not. Instead, the mass (m) of a particle can be calculated based of its measured time-of-flight (t) over a known trajectory of length L – rendering the particle velocity $v = L/t$ – by the following:

$$m^2 = \frac{p^2}{c^2} \left(\frac{c^2 t^2}{L^2} - 1 \right). \quad (3.4)$$

Again, here p represents the momentum of the particle in question. In the case where two particles of different masses (m_1 and m_2) share the same momentum p , the difference between their time-of-flight ($t_1 - t_2$) is given by:

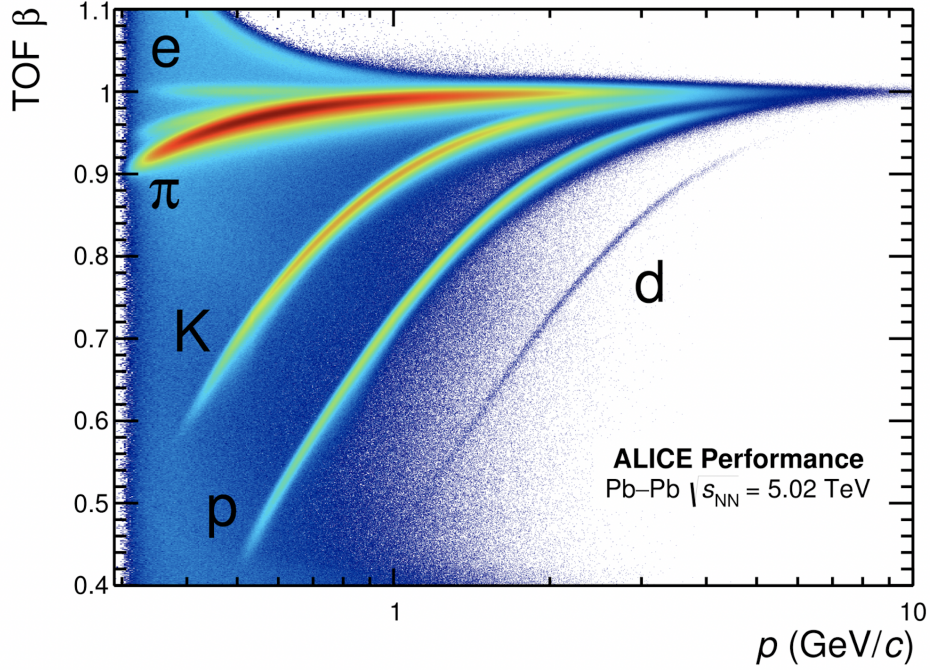
$$t_1 - t_2 = \frac{L}{2c} \left(\frac{m_1^2 c^2 - m_2^2 c^2}{p^2} - 1 \right). \quad (3.5)$$

The total number of standard deviations, n_σ , in the time-of-flight difference between two particles strictly limits the PID capability of the TOF, given a time resolution δt :

$$n_\sigma = \frac{t_1 - t_2}{\delta t}. \quad (3.6)$$

Assuming a particle track length $L = 3.70$ m and a time resolution of 80 ps, an electron to pion separation larger than 3σ is achieved only at low momentum ($p < 700$ MeV/c). On the other hand, pion to kaon and kaon to proton separation larger than 3σ is achieved for $p < 2.4$ GeV/c and $p < 4.0$ GeV/c, respectively. Figure 3.9 shows the ALICE TOF velocity measurements, given by $\beta = v/c$, of various known particles as a function of momentum in

PbPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV.



ALI-PERF-106336

Figure 3.9: TOF measurement of particle velocity, given by $\beta = v/c$ as a function of momentum for ALICE PbPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. Figure from [120].

3.2.4 V0 Scintillators

The VZero (V0) detector is composed of two scintillator arrays, labeled as VZero-A (V0-A) and VZero-C (V0-C), located towards the ATLAS and CMS detectors, respectively. The V0-A detector is positioned at 329 cm from the interaction point while the V0-C is located at 88 cm on either side of the collision vertex. They respectively cover the pseudorapidity range of $2.8 < \eta < 5.1$ and $-3.7 < \eta < -1.7$. Figure 3.10 shows the relative position of the V0 detectors within the bulk of the ALICE central barrel (left) and the collision interaction point (right).

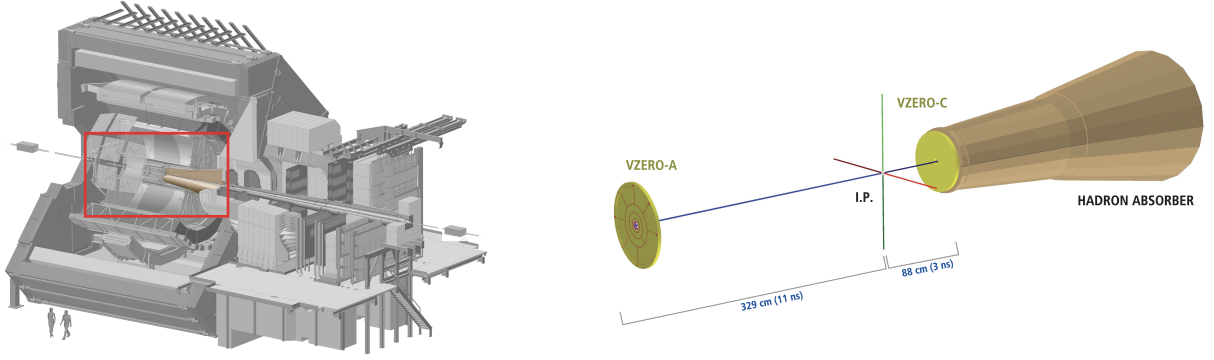


Figure 3.10: Left: Grayed out schematic of the ALICE detector highlighting (encompassed within the red rectangle) the position of the V0 scintillator arrays within the innermost structure of ALICE. Right: Relative positions and distances of the V0-A and V0-C scintillator arrays with respect to the point of interaction within the ALICE detector. Figure adapted from [121].

As depicted in the left and middle panels of Figure 3.11, the two V0 arrays are each radially segmented into four rings and further divided into eight sections along the azimuthal direction. The top and bottom right portions of the figure show the designs for the individual cells of the V0-A and V0-C arrays. The V0 detectors are primarily used as an indicator for the collision centrality determination and serves as a MB trigger for the central barrel detectors, requiring a coincident signal from both the V0-A and V0-C arrays from both primary and secondary particles. It possesses two centrality triggers for heavy ion collisions, controls the muon trigger signal and accurately measures the beam luminosity. The basic working principle of the V0 detector rests on the assumption that all charged particles share identical average response given that they all share the same velocity of travel ($\beta = v/c \approx 1$). This in turn renders a correlated total signal to the total particle multiplicity (i.e. total number of particles) in the detector. In order to extract centrality for each event, the V0 response amplitude is fitted with a Monte Carlo Glauber model. This procedure is outlined in the following subsection.

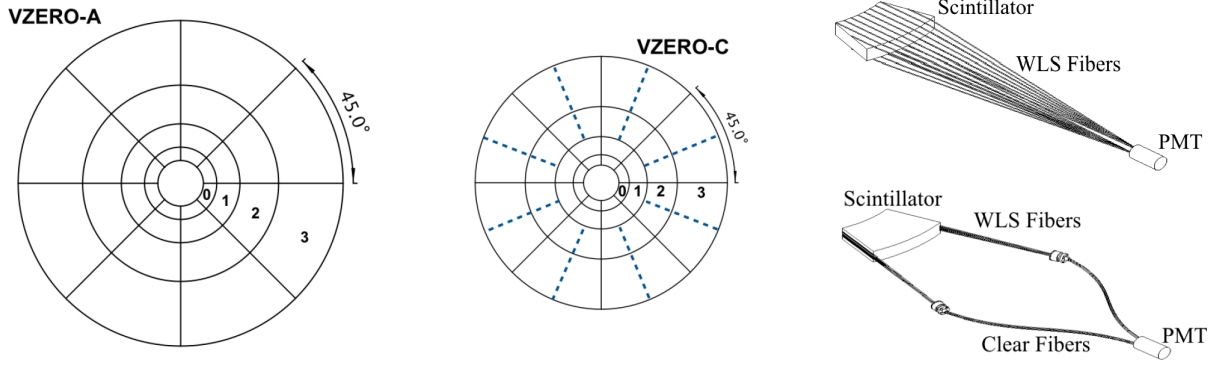


Figure 3.11: The left and middle schematics respectively illustrate the segmentation of the V0A and V0C arrays. Each scintillator has corresponding thickness of 2.5 and 2.0 cm. The segments of the V0-C scintillator on either side of the dashed lines – namely rings 2 and 3 – are both connected to the same photomultiplier tube. The top right schematic depicts the designs for the individual cells of the V0-A, while the bottom one shows the design for rings 0 and 1 of the V0-C. For rings 2 and 3 of the V0-C, a set of scintillator and wavelength shifting fibers are attached by four clear fiber beams to a single photomultiplier tube. Figure from [122].

3.2.5 Centrality Determination

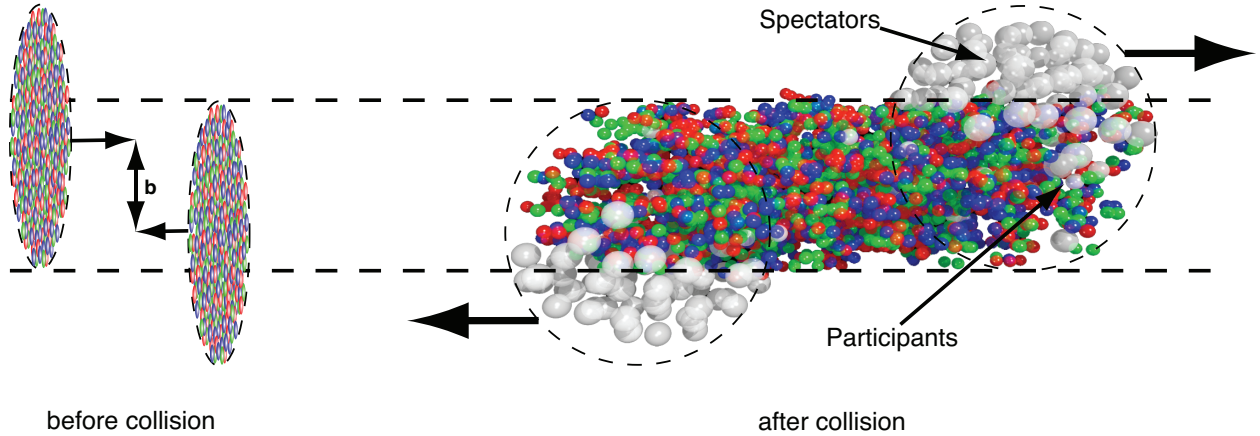


Figure 3.12: Illustration of a relativistic heavy ion collision between two heavy nuclei. The impact parameter (b) represents the offset distance between the centers of the two colliding nuclei and defines overlap region between them. Spectators represent the partons outside the overlap region while participants (i.e. N_{part}) represent the partons within said region. Figure from [123].

The geometry of a heavy ion collision plays a crucial role in the determination of the final

state particle yields produced by each collision event. It is thus important to emphasize how the characterization of each collision is defined. The overlap region between two incoming nuclei, referred to as the *impact parameter* (b), is defined as the distance between the center points of each nucleus perpendicular to the beam axis. Participating nucleons (N_{part}) are any nucleon undergoing at least one inelastic collision (N_{coll}) during the interaction, while non-participating nucleons are simply known as *spectators*. As $b \rightarrow 0$, the degree of overlap between two nuclei increases as well as N_{coll} and N_{part} . An illustration of a relativistic collision between two heavy nuclei, showing participating nucleons, spectating nucleons and b , is depicted in Figure 3.12.

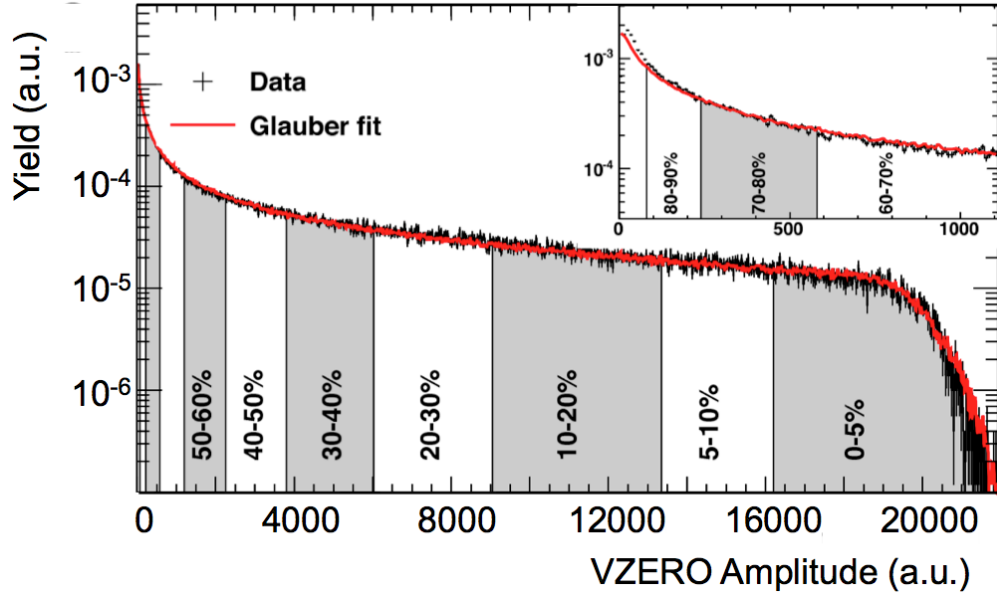


Figure 3.13: Sum of amplitudes distribution in both the two V0A and V0C arrays at ALICE, shown in black, for PbPb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. The Glauber model fit to the data is shown in red. The different centrality classes in the collision system are shown in the shaded areas and are stratified by vertical lines. The shaded areas define the different centrality classes of hadronic collisions. The inset in the top right corner depicts the left-most portion of the distribution pertaining to lower amplitudes. Figure from [122].

Since neither b nor N_{part} are measured directly, an event centrality – expressed in terms of percentiles of the hadronic cross section – is generally used to characterize each collision,

assuming the charged particle multiplicity (N_{ch}) monotonically increases as a function of b . A nuclear density profile T_{AA} is assigned by the Glauber Monte Carlo as it is extracted from N_{coll} from a mapping of the inelastic nucleon-nucleon cross section ($\sigma_{\text{NN}}^{\text{inel}}$). The V0 amplitude is then integrated from high to low intervals, from which the centrality bin measurements are made. In this manner, the 90 – 100% centrality class characterizes 10% of all events pertaining to the low V0 amplitude interval (generally referred to as *peripheral events*) and represents the events with the lowest multiplicity. Conversely, the 0 – 10% centrality class characterizes 10% of all events pertaining to the high V0 amplitude interval (generally referred to as *central events*) and represents the events with the highest multiplicity. Hence, the V0 amplitudes have a one-to-one correspondence with the mean charged particle multiplicity density ($\langle dN_{\text{ch}}/d\eta \rangle$) which relates the mean number of charged particles produced per event to the event centrality and/or multiplicity classes. Figure 3.13 shows the V0 amplitudes for ALICE PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV fitted with a Glauber Model functional form stratified according to their pertaining centrality classes.

3.3 ALICE TPC Upgrade

The foreseen operation of the LHC for the upcoming Runs 3 and 4 would make it such that the interaction rate experienced by the ALICE detector would rise to 50 kHz due to an increase of the overall luminosity of the LHC beams. The previously described MWPC setup of the ALICE TPC possesses an intrinsic limitation of about 3 kHz imposed by the active ion gating grid concept. Above the allowable 3 kHz interaction rate, the ions collected by the gating grid would no longer be kept from drifting back into the TPC drift volume and would lead to large spacetime distortions of the drift field as well as an accumulation of the ion tail and space-charge effects at the accumulation wires. This makes it such that triggered operation via an active gate is insufficient given that 5 collision events would pile

up on average within the $100\ \mu\text{s}$ readout time window described earlier in this chapter. This is necessary to fully upgrade the entire TPC MWPC-based readout system with one capable of providing continuous readout in order to cope with the increased interaction rate emerging from the LHC high luminosity upgrade.

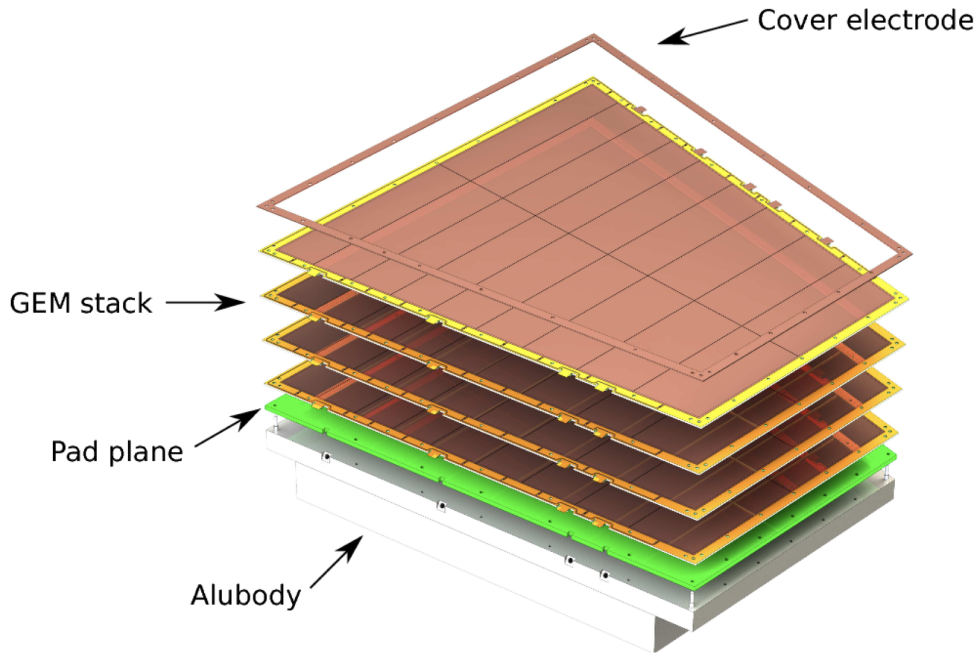


Figure 3.14: Single Chamber GEM Stack Schematic. Figure from [124].

Gas Electron Multipliers (GEMs) [125] have been identified as a viable candidate capable of withstanding the requirements of a continuous readout of the TPC at an interaction rate of 50 kHz. GEMs are thin ($50\ \mu\text{m}$) gaseous detectors made of polyimide foils consisting of a copper layer 2 to $5\ \mu\text{m}$ in thickness on either side. They are perforated in a hexagonal pattern made by double-conical holes by a photolithographic processing method with outer and inner diameters of 70 and $50\ \mu\text{m}$, respectively. GEMs are capable of being stacked atop one another and allow a separation over several amplification stages. A schematic of a single GEM readout chamber is shown in Figure 3.14. Both efficient ion back flow blocking and exceptional particle identification is achieved by stacking four GEM foils with a large (L, 280

μm) and standard (SP, $140\ \mu\text{m}$) hole pitches arranged in a S-LP-LP-S configuration. The purpose of the ALICE TPC upgrade is to replace the MWPC-based TPC with new GEM readout chambers, as well as supplement said chambers with a new high voltage system and new front end electronics. An in-depth description of the ALICE TPC upgrade can be found in [126, 127].

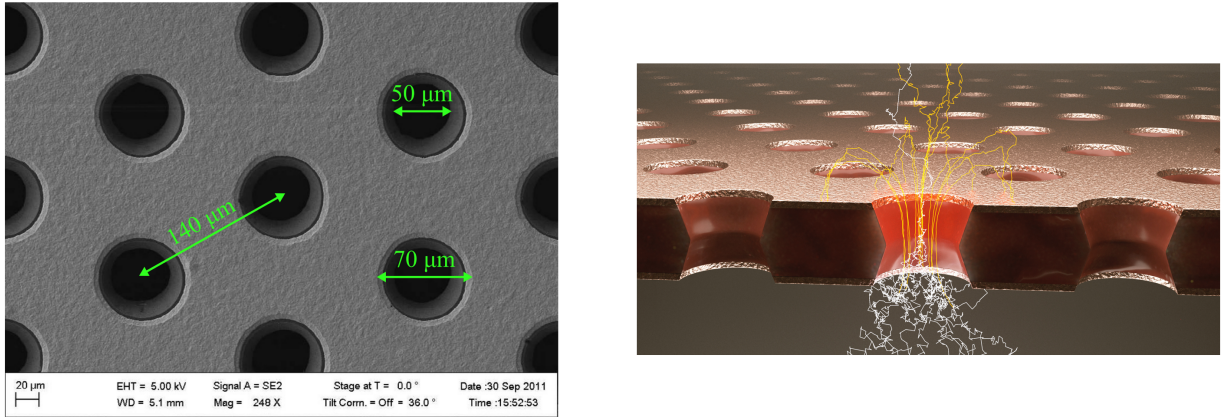


Figure 3.15: Left: Electron microscope image of a GEM foil with standard hole pitch; outer and inner diameters of the conical holes are shown. Right: Garfield/Maxboltz simulation of electron avalanches in a GEM hole. Electrons are shown in crimson, ions in yellow and electron-ion pairs as dark green dots. Figures from [128, 129].

Throughout a portion of my doctoral studies, I took part in the upgrade to the ALICE TPC detector. In particular, I participated in the basic Quality Assurance (QA) of the GEM foil production at CERN.

In the basic QA phase, GEM foils were first subjected to optical inspection and a measurement of the leakage current (I_{leak}) was performed. This was done in order to identify and reject defective foils as early as possible. During the optical inspection, the foils were placed atop a lit table; any obvious defects – such as missing holes, over-etched holes, scratches, etc. – were input into a database. Foils with cuts and/or pieces of copper jutting out from the foil were immediately rejected. Each resistor along the perimeter of the GEM segments was inspected and checked to be properly soldered onto the GEM foil.

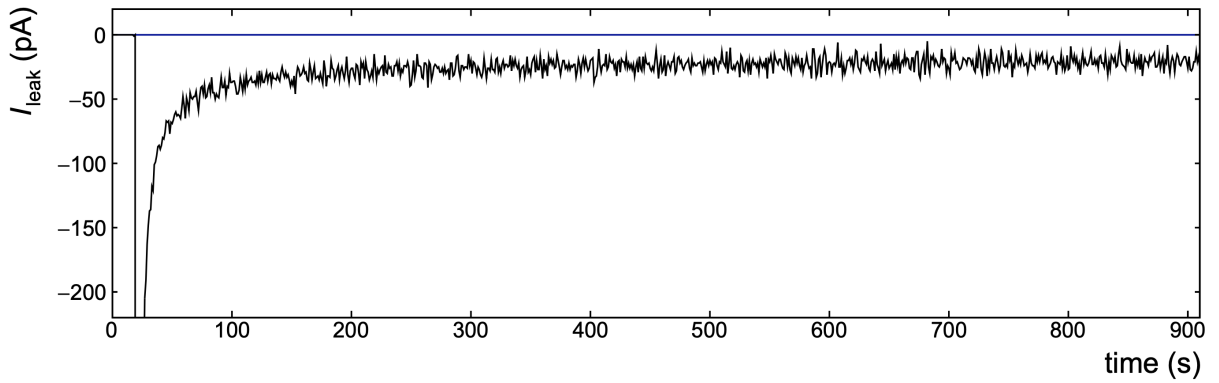


Figure 3.16: Leakage current (I_{leak}) measurement of a typical GEM foil segment. Figure from [127].

A single-channel high voltage source and multichannel pA meter system were used to conduct the leakage current measurement, which was performed in a sealed polycarbonate box filled with a dry nitrogen gas. Spring-loaded pins were used to apply high voltage to each separate GEM foil segment. In order to increase the flow of current across residual dust particles, which may be present in the GEM holes, 550 V were initially applied to the foils as a form of high voltage cleaning. The voltage was then decreased to 500 V and a measurement of I_{leak} was performed for 15 minutes. Any foil with an average leakage current above 0.1 nA in the [450 s, 900 s] time interval was removed from the sealed box, cleaned with a sticky roller, and retested. Figure 3.16 shows a typical leakage current measurement across a single GEM foil segment. Any foils continuously failing this step in the QA process were returned to the producer for cleaning. In the case where the additional cleaning did not reduce the I_{leak} below the allowable threshold, the foil is then discarded. The spark rate of each foil was monitored in addition to the leakage current. Any foil experiencing continuous sparking during the I_{leak} test was immediately removed from the sealed box and cleaned with a tacky roller. In the case of further anomalies, the foil was returned to the producer for cleaning and the aforementioned steps were then repeated. A full description of the Basic

QA protocols and procedures can be found in [130].

At the time of writing of this dissertation, the ALICE TPC upgrade has been long completed, along with other detector upgrades, and is fully operational for the high luminosity LHC runs.

3.4 Complementary Experimentation

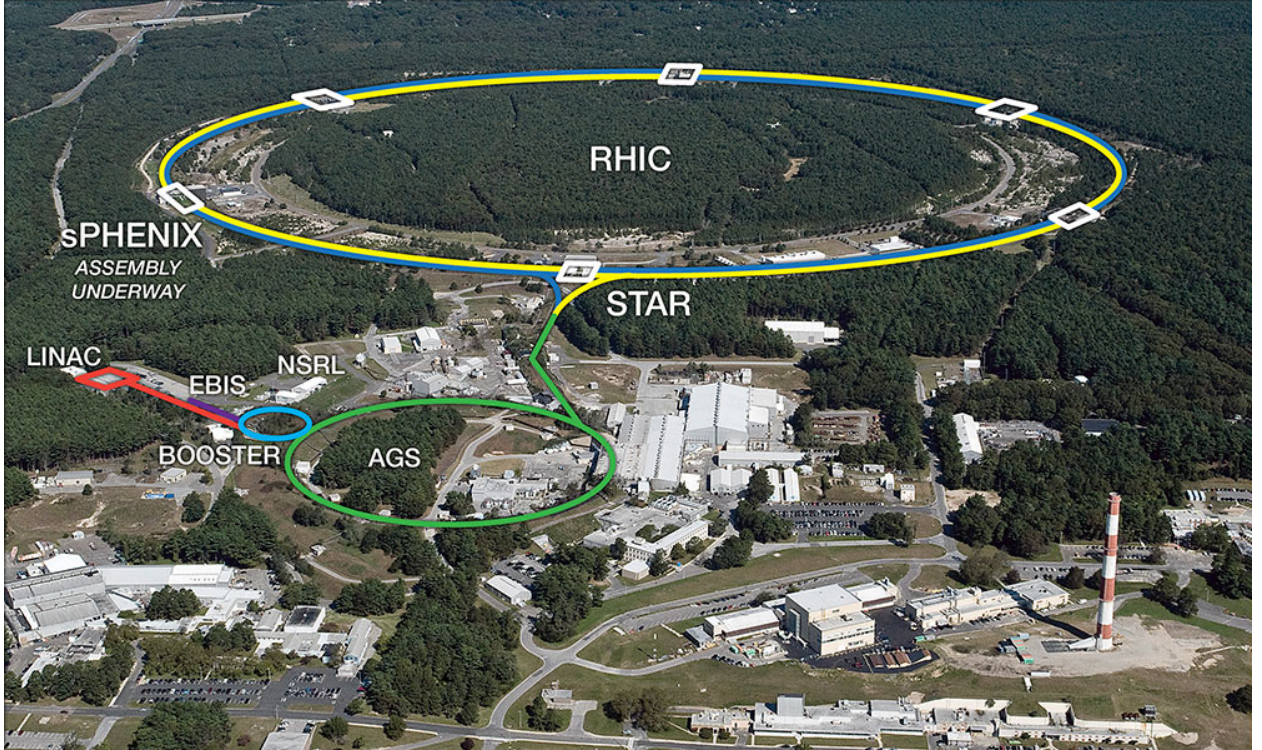


Figure 3.17: RHIC accelerator facility as seen from above. The position of STAR detector, along with additional experiments and facilities at RHIC are shown. Figure from [131].

As stated in Chapters 1 and 3, the ALICE detector is not the only detector around the world capable of generating and probing the QGP via relativistic heavy ion collisions. The STAR detector [132] at RHIC serves as a complementary detector to ALICE insofar as its capacity to detect particles generated by relativistic heavy ion and elementary particle collisions matches that of ALICE, but operates at lower beam energies.

The STAR detector is situated at an interaction point, referred as the *6:00 o' clock*

position, along the RHIC accelerator facility as shown in Figure 3.17. Like the LHC, RHIC is a circular collider, and measures 3.8 km in circumference. The difference in length of the accelerating ring between RHIC and the LHC makes it such that the top center-of-mass energies attained by heavy ion collisions at RHIC – typically gold on gold collisions (AuAu), ranging from $\sqrt{s_{\text{NN}}} = 5 - 200$ GeV – are at least one order of magnitude lower than those achieved by the LHC. This is by no means a disadvantage, given that at lower center-of-mass collision energies, as was shown in Figure 1.10, QGP formation occurs at high temperatures and a non-zero baryochemical potential; unlike at the LHC, where there is a vanishing baryochemical potential. This key difference makes it such that experimental measurements performed by STAR across the 5 - 200 GeV center-of-mass energy range provide phenomenological information lying in a region of the QCD phase diagram that is otherwise inaccessible by measurements made by ALICE alone.

The simplified acceleration chain of Au ion beams at RHIC is as follows [133]: The Electron-Beam Ion Source (EBIS) provides the Au ion sources, which are then transferred to and accelerated by the Radio Frequency Quadrupole (RFQ) and Linear Accelerator (Linac) before they are injected into the booster. The booster then transfers the accelerated beams into the AGS at around 2 MeV per nucleon, which are then focused along both the horizontal and vertical directions by an alternating field gradient. The beam is then transferred by the AGS into the *blue* and *yellow* beam pipes traveling in opposite directions along the RHIC ring. These two beams can then collide at the 6:00 o’ clock position, where the STAR detector is located.

Figure 3.18 shows a cutaway schematic of the STAR detector. The STAR detector weighs approximately 1.2 million kg and like ALICE is enclosed by a large solenoid magnet, which operates at up to 0.5 T. It includes a TPC, TOF and Silicon Vertex Tracker (SVT) which are used for the identification of produced particles.

STAR Detector

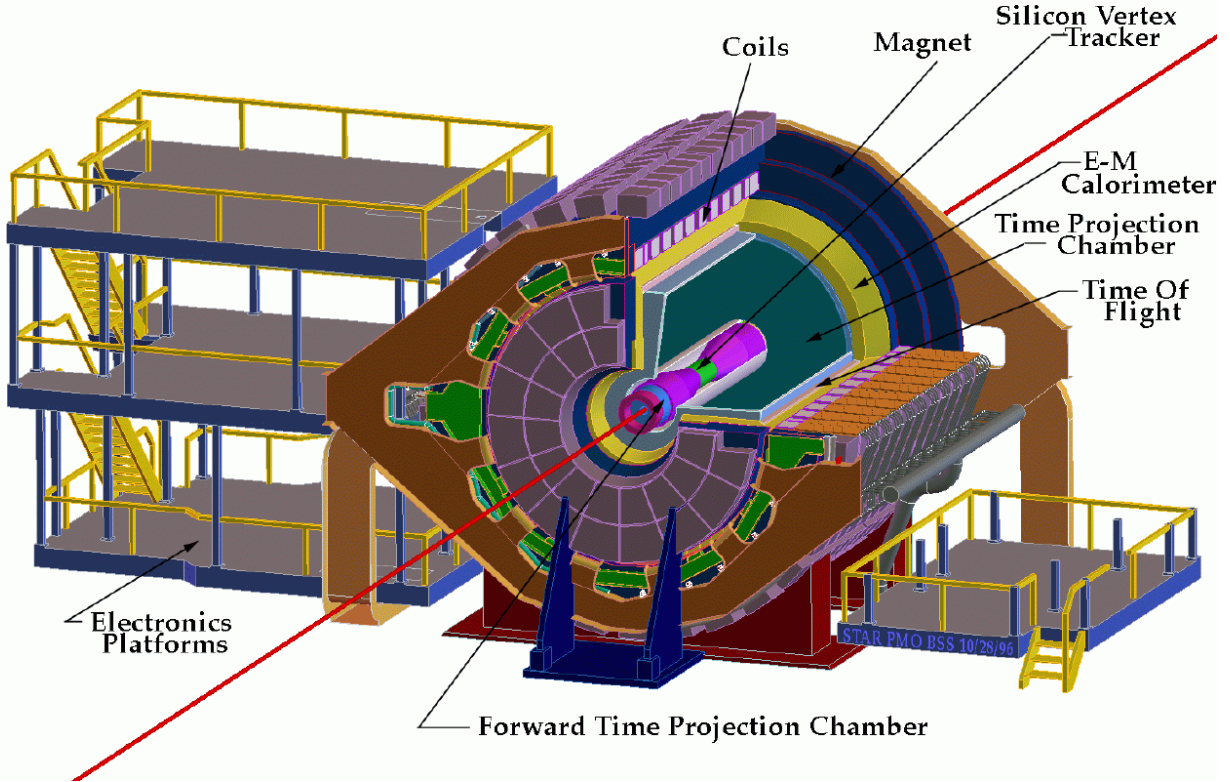


Figure 3.18: Cutaway schematic view of the STAR detector showing its inner detector components. Figure from [132].

The STAR TPC is a cylinder measuring 4.0 m in diameter and 4.2 m in length, covering the $|\eta| < 1.0$ range, and consisting of a MWPC readout system. The STAR TPC measures charged particles over a 0.1 to 30.0 GeV/c momentum range. Global tracks are determined from the fitting of helical tracks onto hits registered by the TPC readout pads. The trajectories of the tracks are then extrapolated onto the beam axis and are used to determine the position of the originating primary vertex for all particle trajectories. In the case where more than one primary vertex is determined, the one with the highest probability is used. Any global tracks with a distance of closest approach less than 3 cm in length are considered

primaries, while the rest are labeled as secondaries. For determining the PID, the STAR TPC also makes use of Equation 3.6 and can identify pions, protons and kaons within the 0.1 to 0.7 GeV/c momentum range. In addition to this, the STAR TPC is also used for the determination of event centrality. The STAR TOF detector covers the $|\eta| < 0.94$ pseudo-rapidity range and possesses the same MRPC working principle as the ALICE TOF. It has the capability of identifying pions and kaons to 1.8 GeV/c in momentum and protons with momentum up to 3.0 GeV/c. The STAR SVT extends the acceptance of the TPC to low momentum particles which do not traverse the bulk of the TPC volume. This is generally useful for detecting short-lived particles, such as strange and multi-strange baryons decaying close to the interaction point.

4 STATISTICAL HADRONIZATION MODEL ANALYSES

As discussed in Chapter 1, the chemical freeze-out stage in the evolution of a heavy ion collision provides the best window from which physical observables can be used to deduce properties of the system, given that from this stage the chemical composition of the system remains fixed. The fact that this stage occurs within the vicinity of the QCD transition temperature allows for the use of final state particles in the scope of statistical hadronization as accurate proxies for the determination of freeze-out parameters of a heavy ion collision, as well as those from smaller collision systems.

In this chapter, I will describe the statistical bootstrap and its modern use as the means for the determination of common freeze-out parameters extracted from particle yields emerging from heavy ion and elementary particle collisions. I will introduce the Hadron Resonance Gas (HRG) model as the primary tool for statistical hadronization modeling and the role it plays in the development of a thermal model. I will then describe the specific thermal model used in my analysis and outline the manner in which the analysis was set up.

4.1 The Hadron Resonance Gas (HRG) Model

Prior to the development of Gell-Mann's Eightfold Way, when a vast number of particles emerged from cosmic ray and bubble chamber experiments, Rolf Hagedorn entertained the idea that the observed particles possessing heavier masses were themselves composed of lighter particles consisting of even lighter particles in a continuous fractal manner (namely, *as a fireball which in turn consists of fireballs, and so on...*) [23]. The fact that the system was seen as being similar to itself at every stage in its evolution made it such that the distribution of the hadron masses can be determined by the manner in which particles of light mass could

be arranged to form into heavier particles, etc. The total mass of the composite system is thus determined from its total energy composed of its interacting substates. A statistical bootstrap equation with a composition volume V_0 corresponding to the system size can be expressed in the following manner [134]:

$$\rho(m, V_0) = \delta(m - m_0) + \sum_N \frac{1}{N!} \left[\frac{V_0}{(2\pi)^3} \right]^{N-1} \int \prod_{i=1}^N dm_i \rho(m_i) d^3 p_i \delta^4 \left(\sum_i p_i - p \right), \quad (4.1)$$

where the mass of the pion is typically used for the m_0 value. The analytic solution to the statistical bootstrap equation takes on the form [135]:

$$\rho(m, V_0) = C m^{-3} \exp \left[\frac{m}{T_H} \right], \quad (4.2)$$

where C is a constant and T_H is referred to as the *Hagedorn Temperature* determined by:

$$\frac{V_0 T_H^3}{2\pi^2} \left[\frac{m}{T_H} \right]^2 K_2 \left[\frac{m}{T_H} \right] = 2 \ln 2 - 1. \quad (4.3)$$

K_2 is a modified Bessel function of the second kind. Assuming $m_0 = m_\pi$, $V_0 \approx (4\pi/3)m_\pi^{-3}$, Hagedorn calculated $T_H \approx 150$ MeV. Letting r_h denote the range of the strong interaction renders $V_0 \approx (4\pi/3)r_h^3$ in the $m_\pi \rightarrow 0$ chiral limit. Setting $r_h = 1$ fm thus results in $T_H \approx 200$ MeV, which reaches an excellent estimation of the QCD transition temperature from rather simple statistical assumptions.

Moreover, Hagedorn suggested that a system of non-interacting hadrons and hadronic resonances could be used to describe an interacting ground state hadron gas. This notion led to the development of the Hadron Resonance Gas in its modern form, which posits that thermally produced hadronic resonances can themselves account for the interactions within a ground state hadronic system. The Grand Canonical Partition function for such a system

can be expressed as:

$$\mathcal{Z}(T, V) = \sum_N \frac{1}{N!} \left[\frac{V}{(2\pi)^3} \int d^3p \exp\left(-\sqrt{p^2 + m_0^2}/T\right) \right]^N, \quad (4.4)$$

which then renders

$$\ln \mathcal{Z}(T, V) = \frac{VTm_0^2}{2\pi^2} K_2\left(\frac{m_0}{T}\right). \quad (4.5)$$

In the case of an interacting gas of ground state hadrons and assuming the interaction itself results from the formation of hadronic resonances, the interacting medium can be expressed in terms of a non-interacting gas containing all possible resonances given by the partition function:

$$\ln \mathcal{Z}(T, V) = \sum_i \frac{VTm_i^2}{2\pi^2} \rho(m_i) K_2\left(\frac{m_i}{T}\right). \quad (4.6)$$

Beginning the sum with the stable ground state m_0 and summing up all of the possible resonant states m_i , where $\rho(m_i)$ denotes the relative weights with respect to m_0 , one can then account for a discrete mass spectrum including all experimentally observed hadrons and hadronic resonances. In this case, the continuum partition function can be expressed by replacing $\rho(m_i)$ by the spin degeneracy factor g_i for the i^{th} particle species taking the sum across all known resonances listed in [45]. Thus, the HRG partition function expressed as the sum of single particle partition functions, Z_i^1 of all hadrons and hadronic resonances takes on the form:

$$\ln Z(T, V) = \sum_i Z_i^1(T, V). \quad (4.7)$$

In the Boltzmann approximation, the single particle partition function Z_i^1 for particles of

mass m_i and spin degeneracy g_i reads

$$Z_i^1(T, V) = g_i \frac{VTm_i^2}{2\pi^2} K_2\left(\frac{m_i}{T}\right) \quad (4.8)$$

The pressure P and energy density ε of a hadron resonance gas with a discrete mass spectrum can be expressed as the sum over all single particle contributions due to the factorization of the partition function in Equation 4.7:

$$P = \sum_i P_i^1, \quad \varepsilon = \sum_i \varepsilon_i^1, \quad (4.9)$$

such that

$$\frac{\varepsilon_i^1}{T^4} = \frac{g_i}{2\pi^2} \left(\frac{m_i}{T}\right)^3 K_2(\beta m_i), \quad (4.10)$$

and

$$\frac{\varepsilon_i^1}{T^4} = \frac{g_i}{2\pi^2} \left(\frac{m_i}{T}\right)^3 \left[\frac{3K_2(\beta m_i)}{\beta m} + K_1(\beta m_i) \right], \quad (4.11)$$

where K_1 and K_2 are modified Bessel functions of the first and second kind, respectively, and $\beta = 1/T$. For the case of a finite temperature T and a vanishing chemical potential, the energy density ε can be related to the pressure P and entropy density s by

$$\varepsilon = sT - P. \quad (4.12)$$

In the case of a non-vanishing chemical potential $\vec{\mu}$, Equation 4.7 takes on the form

$$\ln Z(T, V, \vec{\mu}) = \sum_i Z_i^1(T, V, \vec{\mu}), \quad (4.13)$$

such that $\vec{\mu} = (\mu_B, \mu_S, \mu_Q)$, where μ_B , μ_S and μ_Q are the chemical potentials associated

with the conservation of baryon number, strangeness and electric charge, respectively. For any particle i possessing baryon number B_i , strangeness S_i , electric charge Q_i and spin degeneracy factor g_i , its single particle partition function reads as

$$Z_i^1(T, V, \vec{\mu}) = \frac{g_i V T m_i^2}{2\pi^2} K_2\left(\frac{m_i}{T}\right) \exp\left(\frac{B_i \mu_B + S_i \mu_S + Q_i \mu_Q}{T}\right). \quad (4.14)$$

Equation 4.8 is recovered in the case where $\vec{\mu} = 0$. In order to calculate a particle density n_i of a specific particle i , Z_i^1 in Equation 4.13 is replaced by $\gamma_i Z_i^1$ and a partial derivative is applied with respect to the particle fugacity γ_i

$$n_i = \left. \frac{\partial \ln Z}{\partial \gamma_i} \right|_{\gamma_i=1}. \quad (4.15)$$

Ultimately, $n_i = Z_i^1/V$ where Z_i^1 is defined by Equation 4.14.

4.2 Analysis Procedures

The procedures shown in this dissertation are based on studies which have been previously submitted for publication or accepted for publication and are both available on the arXiv preprint server [136, 137]. They are each separately described in detail in the last two subsections of this chapter.

The main analysis and results in this dissertation were split into two separate procedures – respectively described in Sections 4.2.1 and 4.2.2. The former of the two focused on the application of the SHM framework exclusively to experimental results from heavy ion collisions from both ALICE and STAR across a wide range of center-of-mass collision energies in the 0 - 10% centrality class. On the other hand, the latter of the two procedures focused on the application of the thermal model onto ALICE data from pp, pPb and PbPb across all available multiplicity classes for all three collision systems.

Both analyses were performed using the open source Thermal FIST (‘The FIST’) thermal model package, which is a user-friendly package within the family of HRG Models providing a wide range of options for the HRG model within its framework. Within ‘The FIST’, one can perform thermal model calculations assuming both an ideal and non-ideal hadron resonance gas – the former excluding interactions, while the latter can admit interactions (e.g. van der Waals). Calculations can be made within a Grand Canonical Ensemble, where baryon number, electric charge and strangeness are conserved on average, a Mixed Canonical Ensemble (e.g. a Strangeness Canonical Ensemble in which strangeness is explicitly conserved) and a full Canonical Ensemble with explicit conservation of baryon number, electric charge and strangeness. Thermal fits within ‘The FIST’ are employed by a minimization of χ^2/dof , via the MINUIT package [138], such that

$$\frac{\chi^2}{dof} = \frac{\chi^2}{dof} \sum_{i=1}^N \frac{(N_i^{\text{exp}} - N_i^{\text{HRG}})^2}{\sigma_i^2}, \quad (4.16)$$

where the number of degrees of freedom (dof) is given by the difference between the total number of particle yields and the number of free parameters used in the fit, N_i^{exp} are the experimental particle yields, N_i^{HRG} are the HRG-calculated particle multiplicities and σ_i^2 is the sum of the square of the systematic and statistical errors from the experimental yields. The minimization routine for the χ^2/dof value is accomplished via an iterative parabolic technique employing the first and second derivative tests to determine a local minimum. In principle, lower χ^2/dof values represent a better quality of the performed fit. A detailed description of ‘The FIST’ as well as its open source code can be found in [139] and [140], respectively.

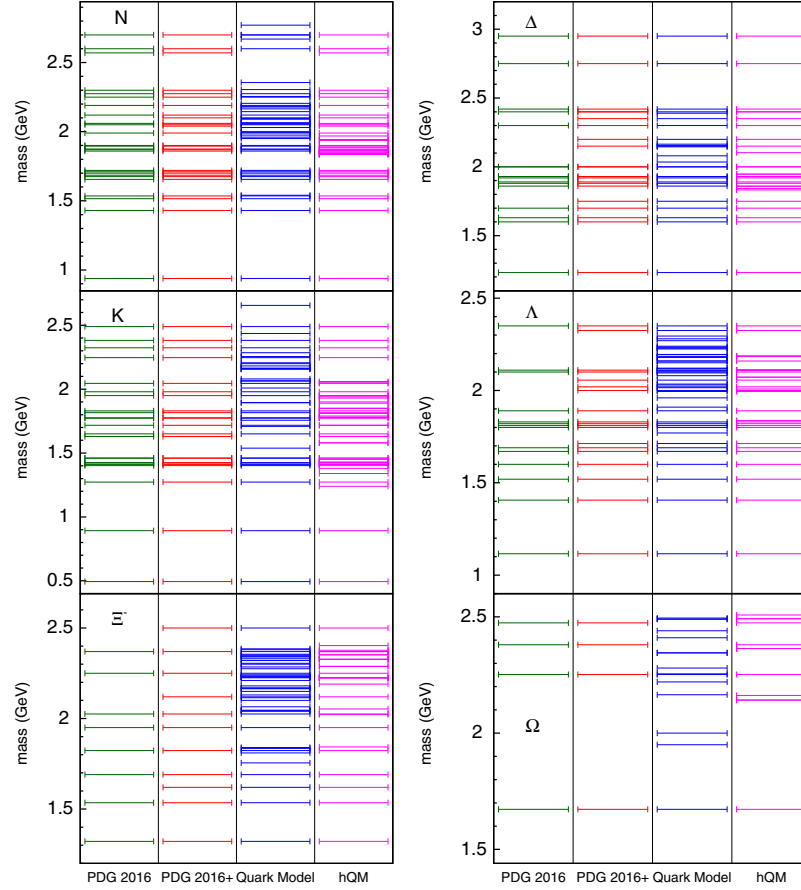


Figure 4.1: Population for different particle species in the PDG2016, PDG2016+, Quark Model and Hypercentral Quark Model (hQM). Figure from [141].

As the HRG input list, all my calculations in ‘The FIST’ used the PDG2016+ hadronic spectrum [141]. As shown by [142–144], deviations of the HRG calculations from the lattice curves of flavor specific susceptibilities at flavor specific temperatures in the QCD crossover region may be affected by the inclusion of certain resonant states. This made it such that a realistic determination of the underlying hadronic spectrum was key to the results presented in this dissertation. The PDG2016+ hadronic spectrum has been shown to be an optimized compromise between too few (experimentally verified) and too many (from a simple Quark Model) excited states when compared to a large number of predictions from Lattice QCD [141]. The PDG2016+ hadronic spectrum includes a total of 738 states consisting of *, **, *

*** and **** states from the 2016 Particle Data Group Data Book [145]. Here, **** states represent well-established states from experimental results, while * states are those with least experimental confirmation, having incomplete decay channels and/or branching ratios. Figure 4.1 shows a comparison of the population differences across several particle species between the PDG2016, PDG2016+, Quark Model and Hypercentral Quark Model (hQM).

4.2.1 Energy and Flavor Dependence of Chemical Freeze-out at ALICE and STAR

Although there exists a wide range of options for the HRG model within ‘The FIST’ framework, this portion of the analysis was restricted to the default – namely, modeling an ideal non-interacting gas of hadrons and resonances within a Grand Canonical Ensemble (GCE), where baryon number, electric charge and strangeness are conserved on average. Yield data for π^+ , π^- , K^+ , K^- , p , $\bar{p}$, Λ , $\bar{\Lambda}$, Ξ^- , $\bar{\Xi}^+$, Ω^- , $\bar{\Omega}^+$, K_S^0 , and ϕ for ALICE PbPb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [146–149] and preliminary results at 5.02 TeV [105] in the 0 - 10% centrality class, as well as STAR AuAu collisions at $\sqrt{s_{NN}} = 11.5, 19.6, 27.0, 39.0, 62.4$ and 200 GeV [150–155] were used. Yield measurements from AuAu collisions at $\sqrt{s_{NN}} = 7.7$ GeV were excluded from this analysis due to the wider centrality binning of the data, particularly for multi-strange baryons.

For the sake of brevity, a shorthand notation when naming the fits was introduced, such that both particle and anti-particle species are coupled into a single character symbol (e.g. Ω refers to both Ω^- and $\bar{\Omega}^+$, etc.). This shorthand is used for the remainder of this dissertation.

(Anti-)proton yields for the STAR data in Refs. [150, 153, 154] are all inclusive. In order to correct for weak-decay feed-down contributions from $(\bar{\Lambda})\Lambda$, an interpolation of the contributions to (anti-)proton yields based on the results Ref. [156] was used – where the weak decay feed-down contribution to (anti-)proton yields was shown to be asymptotic around 35% as a

function of beam energy based on yields calculated through parametrizations of T_{ch} and μ_{B} . Table 4.1 summarizes the interpolation results. The contribution from $(\bar{\Lambda})\Lambda$ to (anti-)proton yields are labeled as δ . These δ values were subtracted from unity and multiplied by their respective (anti-)proton yields. The resulting (anti-)proton yields were then used for the entirety of this analysis portion. This procedure should be considered an upper limit for the feed-down contribution, since the experiment imposes loose primary vertex cut on the $(\bar{\Lambda})\Lambda$ decay daughter candidates. Nevertheless, these percentages are in general agreement with estimates in the aforementioned STAR papers.

Table 4.1: Interpolation results based on results in Ref. [156] for weak-decay feed-down contributions to (anti-)proton yields at STAR AuAu Collisions from $\sqrt{s_{\text{NN}}} = 11.5$ to 200 GeV.

	Proton	Anti-Proton
$\sqrt{s_{\text{NN}}} \text{ (GeV)}$	δ	δ
11.5	23.00%	48.00%
19.6	27.50%	44.00%
27.0	29.50%	41.50%
39.0	31.00%	40.00%
64.2	32.00%	38.50%
200	34.00%	36.50%

Every thermal fit was performed with T_{ch} (MeV) and V (fm^3) as free parameters, setting γ_{S} unity by assuming a full saturation of strangeness within V . This analysis focused on varying the particle species included in the fit. The particle species included in the temperature fits were $\pi K p$ (light), $\pi K p \Lambda \Xi \Omega K_{\text{S}}^0 \phi$ (all) and $K \Lambda \Xi \Omega K_{\text{S}}^0 \phi$ (strange), respectively. The inclusion of the kaons in the light fit was done in order to avoid too few degrees of freedom in the fit and has no effect on the extracted freeze-out temperature, since the kaon yield is rather insensitive to the temperature, as was also shown previously in Ref. [157].

Table 4.2: ‘The FIST’ Grand Canonical Ensemble Yield Fits via the PDG2016+ hadronic spectrum for the most central (0 - 10%) PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV at ALICE. The top panel shows the results for light ($\pi K p$), full ($\pi p \Lambda \Xi \Omega K_S^0 \phi d$), and strange ($K \Lambda \Xi \Omega K_S^0 \phi$) particle fits, while the bottom panel shows the results for the alternative light ($\pi p d$), full ($\pi p \Lambda \Xi \Omega K_S^0 \phi d$), and strange ($K \Lambda \Xi \Omega K_S^0 \phi$) particle fits, respectively, using the (anti-)deuteron yields from Ref. [158]. For all fits, $\mu_B = 0$, while T_{ch} and V were used as free parameters.

Excluding (anti-)deuteron				
Fit	μ_B (MeV)	T_{ch} (MeV)	$V(fm^3)$	χ^2/dof
$\pi K p$	0.0	143.2 ± 2.79	8031.7 ± 1263	5.65/4
$\pi K p \Lambda \Xi \Omega K_S^0 \phi$	0.0	149.6 ± 1.76	5764.4 ± 635.8	23.4/12
$K \Lambda \Xi \Omega K_S^0 \phi$	0.0	153.9 ± 2.30	4389.7 ± 640.8	10.5/8
Including (anti-)deuteron				
Fit	μ_B (MeV)	T_{ch} (MeV)	$V(fm^3)$	χ^2/dof
$\pi p d$	0.0	144.6 ± 2.39	7911.6 ± 1177	5.45/4
$\pi K p \Lambda \Xi \Omega K_S^0 \phi d$	0.0	150.1 ± 1.65	5613.6 ± 588.5	23.9/14
$K \Lambda \Xi \Omega K_S^0 \phi$	0.0	153.9 ± 2.30	4389.7 ± 640.8	10.5/8

As a proof of concept ad-hoc cross-check, an independent fit replacing the kaon yields with (anti-)deuteron yields in the light fit for the most central (0 - 10%) PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [158] was performed and no significant difference in the extracted freeze-out parameters and the quality of fit was found – shown in Table 4.2. Nevertheless, the present study did not implement the aforementioned $\pi p d$ light fit in the calculations since experimental (anti-)deuteron yields were not available for the entire center-of-mass energy range in question. For the two ALICE energies, μ_B was set to zero. For all fits to STAR AuAu data, μ_B was allowed to be a free parameter in the fits to also gauge its sensitivity to the flavor-specific fits.

4.2.2 System Size and Flavor Dependence of Chemical Freeze-out at ALICE

Unlike the previous analysis procedure, the System Size analysis approach does not employ the use of the SHM framework for a non-zero baryochemical potential. Instead, a focus is placed solely on the ALICE data available for all three collision systems (i.e. pp, pPb and PbPb).

The GCE treatment has been shown to inadequately reproduce experimental results in a centrality dependent manner across pp, pPb and PbPb collision systems measured by the ALICE Collaboration – particularly where $\langle dN_{\text{ch}}/d\eta \rangle \leq 20$ [159]. The aforementioned shortcoming of the GCE treatment is often attributed to the presence of non-equilibrium strangeness production in the smaller systems and may be partially remedied by employing a Strangeness Canonical Ensemble (SCE), in which strangeness is explicitly conserved within a correlation volume – labeled as V_C – using a single value of T_{ch} for all particle species. The question arises as to whether out of equilibrium strangeness in small systems provides an adequate description of the overall production of final state (multi)strange hadrons as measured at LHC energies. The absence of equilibrium, commonly represented by a non-negligible strangeness saturation parameter (γ_S) within the fireball volume of hadronic collisions, is known to diminish with increasing collision energy [160, 161].

In principle, strangeness equilibration scaling as an inverse function of collision energy can be equated with an increasing value of γ_S . At LHC energies, γ_S is expected to approach unity [159, 162], i.e. full saturation of strangeness is achieved. The dominating presence of strangeness enhancement at ALICE, even in small systems, is evident from an energy dependent comparison of final state anti-hadron to π^+ yields measured by ALICE and STAR [163], where production of strange baryons is fully saturated for collision energy values above $\sqrt{s_{\text{NN}}} = 62.4$ GeV (Figure 4.2).

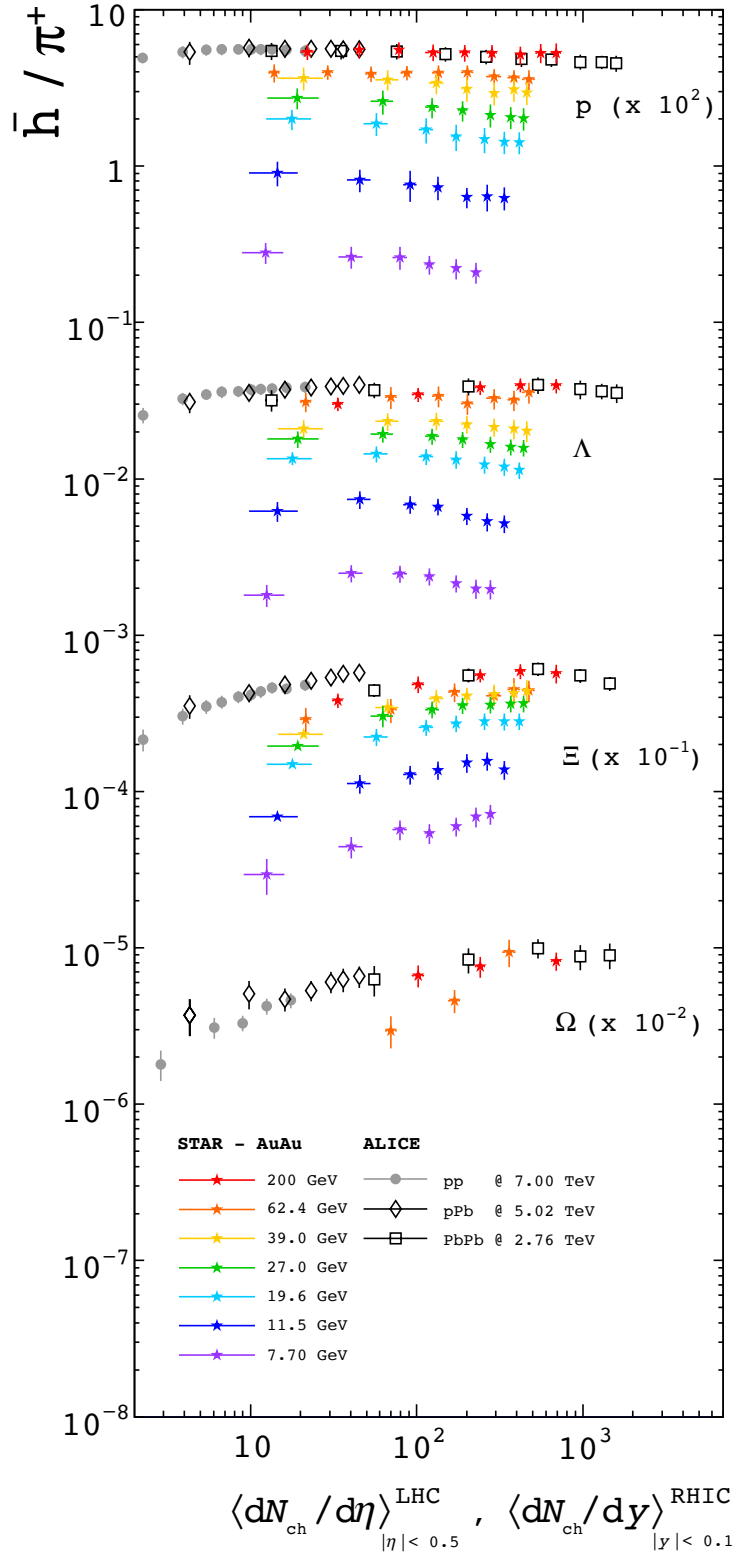


Figure 4.2: Anti-baryon to proton yield ratios as function of charged particle multiplicity density and collision energy at ALICE and STAR. Figure adapted from [163].

Nevertheless, in order to fully describe the experimental data, further considerations are made regarding the interplay between V_C , the fireball volume (V), the experimental rapidity window (Δy) and γ_S [164, 165].

This portion of the analysis extended the flavor-dependent two chemical freeze-out (2CFO) temperature approach described in the previous section to pp, pPb and PbPb collision systems measured by ALICE as a function of increasing $\langle dN_{\text{ch}}/d\eta \rangle$, rendering a natural explanation of strangeness enhancement from small to large systems at vanishing μ_B without further modifying Δy , V_C or γ_S .

In this analysis, an ideal non-interacting gas of hadrons and resonances within both the GCE and the SCE scenarios was employed for the sake of comparison. The analysis was itself two-fold: To gauge the sensitivity of the chemical freeze-out temperature at vanishing μ_B relative to the ensemble of choice, and to employ a 2CFO treatment onto the reigning ensemble.

Yield data for π^+ , π^- , K^+ , K^- , p , $\bar{p}$, Λ , $\bar{\Lambda}$, Ξ^- , $\bar{\Xi}^+$, Ω^- , $\bar{\Omega}^+$, K_S^0 , and ϕ for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV [36], pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [166–168] and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [146–149] across all available multiplicity classes were included in the analysis. The used multiplicity/centrality bins and corresponding values of charged particle multiplicity density for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV are defined in Table 4.3.

Table 4.3: Available centrality bins and corresponding values of charged particle multiplicity density for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. For the pp sample, the multiplicity classes are labelled in accordance to their generalized definitions in [36].

pp at 7.00 TeV		
Multiplicity Class	Centrality Bin	$\langle dN_{\text{ch}}/d\eta \rangle$
I-II	0-5%	17.47 ± 0.524
III-VI	5-28%	10.383 ± 0.313
VII-VIII	28-48%	6.057 ± 0.19
IX-X	48-100%	2.886 ± 0.135
pPb at 5.02 TeV		
	0-5%	45 ± 1
	5-10%	36.2 ± 0.8
	10-20%	30.5 ± 0.7
	20-40%	23.2 ± 0.5
	40-60%	16.1 ± 0.4
	60-80%	9.8 ± 0.2
	80-100%	4.3 ± 0.1
PbPb at 2.76 TeV		
	0-10%	1447.5 ± 54.5
	10-20%	966 ± 37
	20-40%	537.5 ± 19
	40-60%	205 ± 7.5
	60-80%	55.5 ± 3

In the GCE configuration, the thermal fits were performed with T_{ch} (MeV) and V (fm³) as free parameters, keeping $\mu_{\text{B}} = 0$ – as is the case at LHC energies – and setting γ_{S} and γ_{q} to unity in order to ensure a full saturation of strangeness and electric charge. Systematic HRG-based studies on the determination of the latter three parameters at top LHC energies can be found in Refs. [169,170]. The analysis focused on varying the particle species included in the fit in order to gauge the sensitivity of T_{ch} to each fit. The particle species included in the flavor-dependent temperature fits were $\pi K p$, $\pi K p \Lambda \Xi \Omega K_{\text{S}}^0 \phi$ and $K \Lambda \Xi \Omega K_{\text{S}}^0 \phi$, hereinafter referred to as “light”, “all” and “strange”, respectively. A comparison between the extracted T_{ch} values, for the light, strange and all fits, and their corresponding χ^2/dof measures as a function of their $\langle dN_{\text{ch}}/d\eta \rangle$ values across all three collision systems was performed. Again, since the kaon yields were shown to be insensitive to the freeze-out temperature, they were included in the light fit for the sake of having sufficient degrees of freedom. This analysis also did not implement the aforementioned $\pi p d$ light fit in the calculations since experimental (anti-)deuteron yields were not yet available for all multiplicity bins across all three collision systems at ALICE.

In the SCE configuration, the thermal fits were instead performed with T_{ch} (MeV) and V (fm³) as free parameters, keeping $\mu_{\text{B}} = 0$, $V_{\text{C}} = V$ and setting γ_{S} and γ_{q} to unity. Setting $V_{\text{C}} = V$ was done in order to guarantee a local conservation of strangeness within the calculated fireball volume per unit rapidity. The effect of letting $V_{\text{C}} \neq V$ on the calculated particle yields is shown in [164], however, this dissertation imposed the $V_{\text{C}} = V$ condition in order to simplify the thermal fit procedure without having to introduce an additional parameter to the fits.

The initial SCE analysis followed the same procedure than the early GCE trial, however, the following procedures were not performed for the GCE approach due to the deterioration of the quality of its fits across all three collision systems. A similar approach was presented in

[171,172], with a notable difference in the manner in which the pion yields were calculated in the thermal model. In this analysis, the pion, and all other, yields were based on experimental values measured over a single unit of rapidity around mid-rapidity.

For the SCE analysis, the $\pi K p \Lambda \Xi \Omega K_S^0 \phi$ yields were calculated by fixing the temperature to the flavor specific freeze-out temperatures at $\mu_B = 0$ to $T_L = 150$ MeV for πp and $T_S = 165$ MeV for $K \Lambda \Xi \Omega K_S^0 \phi$. The anti-hadron to π^+ ratios were calculated as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ and results were compared with the experimental data. The choice of the anti-hadron to π^+ ratio was made, as in [163], in order to use only particles produced during the evolution of the fireball. This will allow one to facilitate future centrality-dependent comparisons to lower collision energy measurements, where $\mu_B \neq 0$. In order to explicitly show the temperature dependence of the fireball volumes, the volumes were calculated across all three systems as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ with temperatures fixed to the aforementioned flavor specific temperatures, as well as $T = 158$ MeV for a non-flavor-dependent temperature.

5 RESULTS AND DISCUSSION

5.1 Energy and Flavor Dependence of Chemical Freeze-out at ALICE and STAR

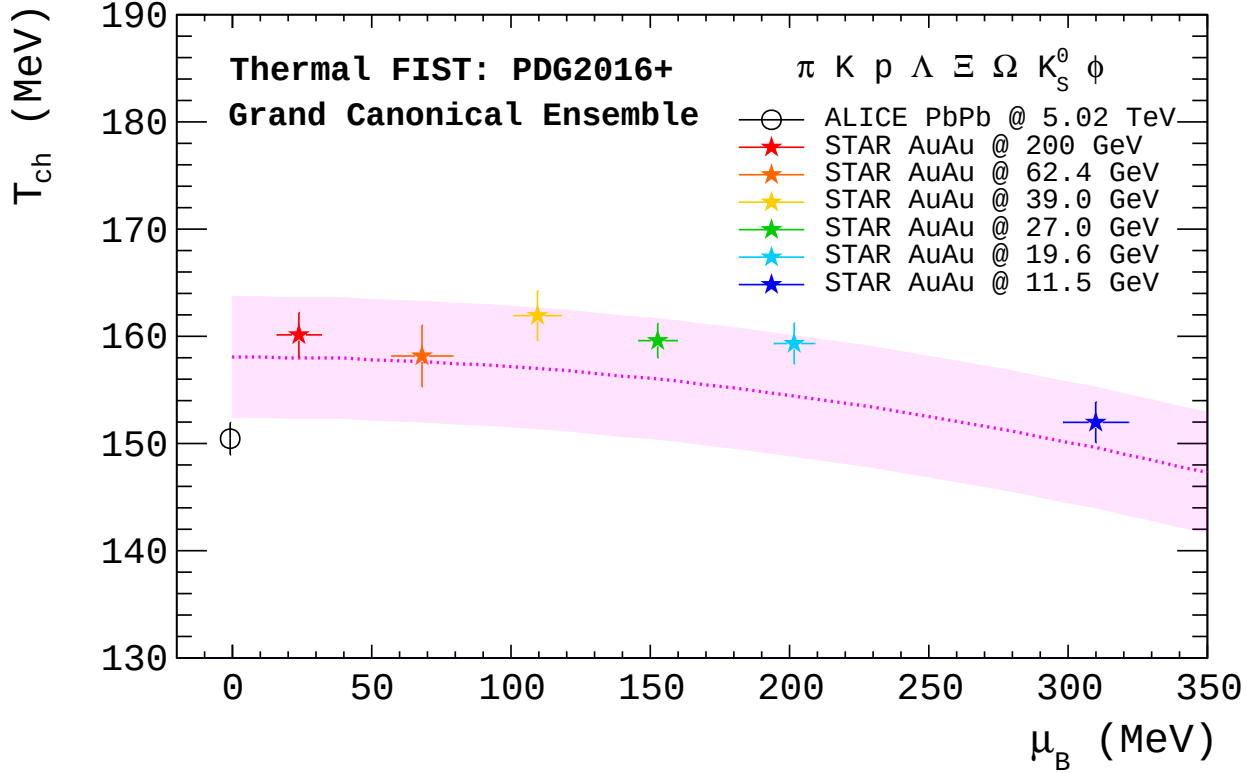


Figure 5.1: Full Grand Canonical Ensemble fits to STAR and ALICE data measured at collision energies ranging from $\sqrt{s_{\text{NN}}} = 11.5$ GeV to 5.02 TeV (0 - 10%) using The FIST with the PDG2016+ hadronic spectrum. Magenta bands shows a spline fit to the points. Figure from [136].

The freeze-out parameters, T_{ch} and V , were extracted for the full $(\pi K p \Lambda \Xi \Omega K_S^0 \phi)$, light $(\pi K p)$, and strange $(K \Lambda \Xi \Omega K_S^0 \phi)$ particle thermal fits. Figure 5.1 shows the extracted T_{ch} values from the full fits as a function of μ_B . The error bars for all extracted freeze-out parameter pairs were determined by the χ^2/rm minimization for each individual fit at each collision energy and are representative of their respective experimental yield errors. The magenta line represents the spline fit function with three nodes (PbPb at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and AuAu

at $\sqrt{s_{\text{NN}}} = 27.0$ and 19.6 GeV) to the freeze-out parameters. The magenta band represents the residual to the spline fit and is determined by the mean square error of the spline function.

Figure 5.2 shows T_{ch} from the light and strange fit as a function of μ_{B} ; both fits are compared to Lattice QCD calculations in Ref. [173]. Each flavor-dependent spline fit and error band was determined in the same manner as in Figure 5.1. The width of the lattice curve is based on the uncertainty of the pseudocritical temperature of the chiral susceptibility [173].

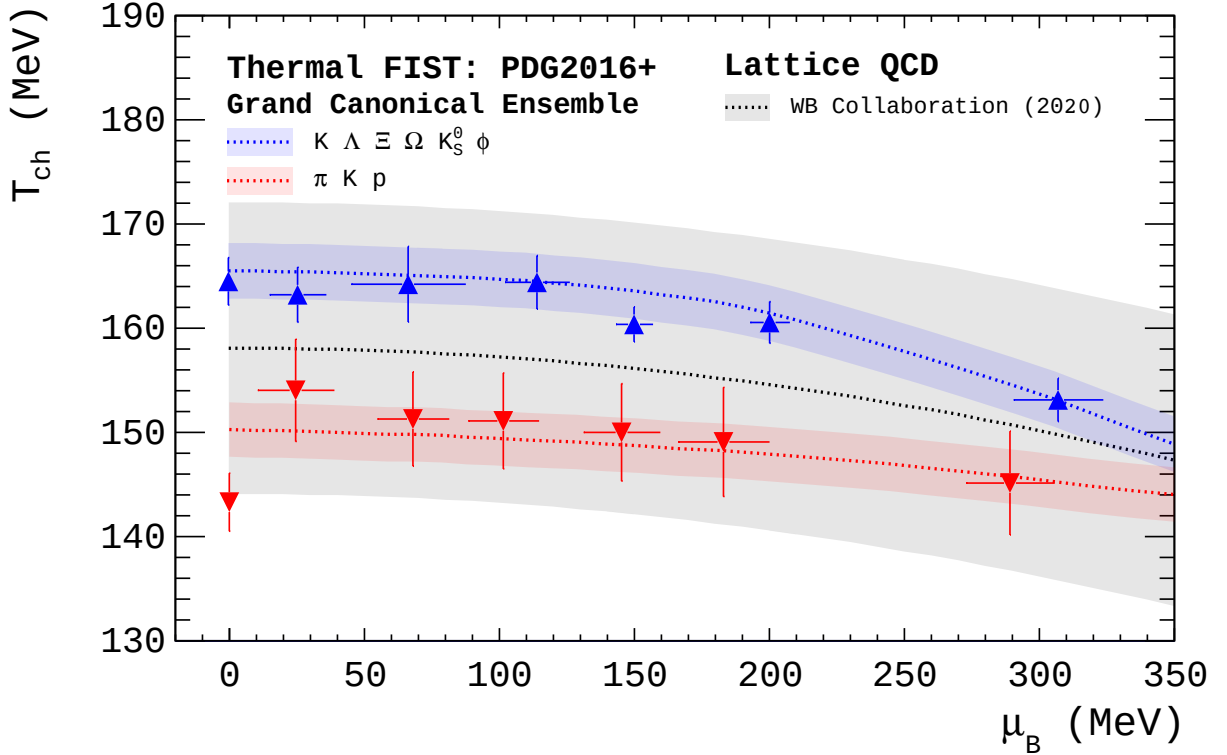


Figure 5.2: Strange (blue points) and light (red points) Grand Canonical Ensemble fits to STAR and ALICE data measured at collision energies ranging from $\sqrt{s_{\text{NN}}} = 11.5$ GeV to 5.02 TeV (0 - 10%) via ‘The FIST’ using the PDG2016+ hadronic spectrum. Figure from [136].

Detailed fit results for each energy including V and χ^2/dof , are shown in Tables 5.1 and 5.2. In general, the separation into light and strange particles improves the quality of the fits by at least a factor two at all energies.

Table 5.1: ‘The FIST’ Grand Canonical Ensemble Yield Fits via the PDG2016+ hadronic spectrum for collision energies ranging from $\sqrt{s_{\text{NN}}} = 11.5$ GeV to 5.02 TeV showing the full ($\pi K p \Lambda \Xi \Omega K_S^0 \phi$) particle fits. For all fits, μ_B , T_{ch} , and V were used as free parameters.

$\pi K p \Lambda \Xi \Omega K_S^0 \phi$				
$\sqrt{s_{\text{NN}}}$ (GeV)	μ_B (MeV)	T_{ch} (MeV)	$V(\text{fm}^3)$	χ^2/dof
5020	0.0	150.4 ± 1.50	6238.8 ± 538.8	78.9/12
2760	0.0	149.6 ± 1.76	5764.4 ± 635.8	23.4/12
200	25.0 ± 7.94	160.0 ± 2.07	1596.3 ± 198.3	23.3/10
62.4	68.4 ± 11.0	158.1 ± 2.87	1151.7 ± 178.2	41.0/11
39	110.3 ± 8.51	161.8 ± 2.32	751.6 ± 103.0	15.5/11
27	154.2 ± 6.92	159.4 ± 1.60	741.8 ± 70.4	12.8/11
19.6	202.8 ± 7.29	159.3 ± 1.91	643.2 ± 72.6	15.2/11
11.5	310.7 ± 11.7	151.8 ± 1.89	649.4 ± 78.8	15.9/11

Table 5.2: ‘The FIST’ Grand Canonical Ensemble Yield Fits via the PDG2016+ hadronic spectrum for collision energies ranging from $\sqrt{s_{\text{NN}}} = 11.5$ GeV to 5.02 TeV. The top and bottom sections show the light ($\pi K p$) and strange ($K\Lambda\Xi\Omega K_S^0\phi$) particle fits, respectively. For all fits, μ_B , T_{ch} , and V were used as free parameters.

$\pi K p$				
$\sqrt{s_{\text{NN}}} \text{ (GeV)}$	$\mu_B \text{ (MeV)}$	$T_{\text{ch}} \text{ (MeV)}$	$V(fm^3)$	χ^2/dof
5020	0.0	142.4 ± 1.70	9371.6 ± 902.1	14.6/4
2760	0.0	143.2 ± 2.79	8031.7 ± 1263.0	5.65/4
200	24.5 ± 14.0	153.9 ± 4.91	2210.0 ± 553.4	3.10/3
62.4	67.8 ± 13.1	151.2 ± 4.53	1721.7 ± 394.5	7.79/3
39	101.4 ± 12.9	151.0 ± 4.61	1368.8 ± 338.9	2.60/3
27	145.3 ± 14.0	149.9 ± 4.68	1333.3 ± 336.8	2.79/3
19.6	182.9 ± 16.7	149.0 ± 5.24	1186.1 ± 341.5	7.66/3
11.5	289.4 ± 16.2	145.0 ± 4.98	1074.8 ± 300.1	2.67/3
$K\Lambda\Xi\Omega K_S^0\phi$				
$\sqrt{s_{\text{NN}}} \text{ (GeV)}$	$\mu_B \text{ (MeV)}$	$T_{\text{ch}} \text{ (MeV)}$	$V(fm^3)$	χ^2/dof
5020	0.0	164.4 ± 2.28	3086.8 ± 366.7	1.58/8
2760	0.0	153.9 ± 2.30	4389.7 ± 640.8	10.5/8
200	25.7 ± 10.3	163.2 ± 2.64	1287.5 ± 208.8	17.3/6
62.4	66.7 ± 21.0	164.2 ± 3.63	784.5 ± 152.2	21.4/7
39	116.0 ± 11.7	164.4 ± 2.57	643.6 ± 97.7	6.73/7
27	156.0 ± 7.88	160.4 ± 1.68	695.0 ± 69.3	3.55/7
19.6	206.5 ± 8.30	160.7 ± 2.00	585.2 ± 69.3	3.87/7
11.5	321.0 ± 14.7	153.4 ± 2.08	574.8 ± 77.2	5.36/7

In Figure 5.2, the light and strange fits consistently fall within the lattice QCD crossover width as determined by detailed studies of the temperature dependence of the chiral order parameters [173]. The flavor-dependent fits agree with the calculated freeze-out temperatures from net-proton, net-charge and net-kaon fluctuations up to $\mu_B \simeq 150$ MeV [174,175], shown in Figure 5.3, which show similar differences in temperatures at freeze-out between the light and strange quark sector.

Both light and strange T_{ch} values in Figure 5.2 remain constant with increasing μ_B until $\mu_B \simeq 100$ MeV, where the strange and light fit begin to approach. The two fits converge within errors at $\mu_B \simeq 300$ MeV; therefore, it is proposed that a separate treatment of strange and light particles might not be meaningful at $\mu_B \geq 300$ MeV. The convergence of the two flavor-dependent temperatures is expected in the vicinity of a critical point in the QCD phase diagram.

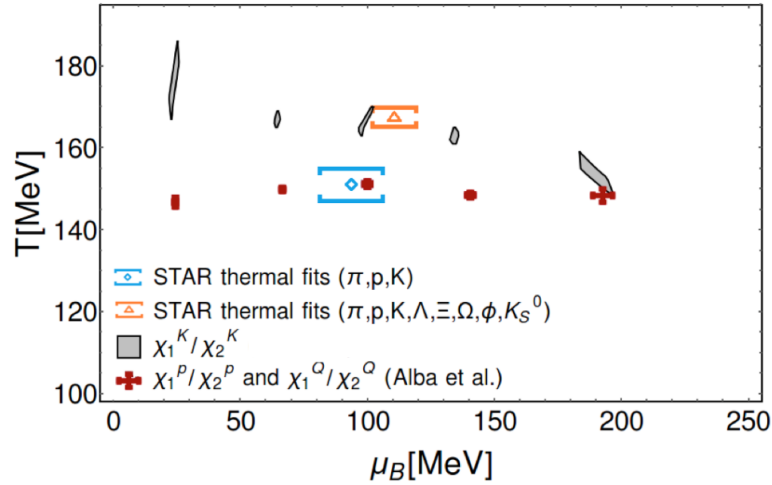


Figure 5.3: Extracted freeze-out parameters from the RHIC-Beam Energy Scan at $\sqrt{s_{\text{NN}}} = 200, 62.4, 39.0, 27.0$ and 19.6 GeV, shown from left to right. The points in red depict the values obtained from a combined fit of χ_1^Q/χ_2^Q and χ_1^P/χ_2^P , while the gray bands are obtained from fitting χ_1^K/χ_2^K . The freeze-out parameters obtained by the STAR Collaboration at $\sqrt{s_{\text{NN}}} = 39.0$ GeV from thermal fits to $\pi K p \Lambda \Xi \Omega K_S^0 \phi$ and to $\pi K p$ are shown as an orange triangle and a blue diamond, respectively, each with their corresponding errors (shown as closed brackets). Figure from [175].

The greatest difference between the two temperatures is seen at the highest energy, namely the ALICE top energy. To address the impact of a flavor dependent freeze-out on the particle abundances, yields were calculated for the full particle set ($\pi K p \Lambda \Xi \Omega K_S^0 \phi$) for ALICE PbPb 5.02 TeV with a one chemical freeze-out (1CFO) approach and were compared with yields calculated for light ($\pi K p$) and strange ($K \Lambda \Xi \Omega K_S^0 \phi$) particles separately with a two chemical freeze-out (2CFO) approach. The temperature(s) and volume(s) were fixed to the T_{ch} and V values shown in Table 5.2 for $\sqrt{s_{\text{NN}}} = 5020$ GeV. In the 1CFO approach, yields were calculated using a temperature of 150.4 MeV. In the 2CFO approach, the light and strange particle yield calculations were done with temperatures fixed to 142.4 MeV and 164.4 MeV, respectively. It is noted that the 1CFO temperature differs from the value quoted by ALICE, which is based on the Heidelberg-GSI model, by about 4 MeV, most likely due to the difference in the hadronic input spectrum [105].

Figure 5.4 shows a comparison between the 1CFO and 2CFO approaches for data sets at two exemplary energies, namely the preliminary 5.02 TeV central PbPb data from ALICE and the 39.0 GeV central AuAu data from STAR. The deviations of each yield calculation from the experimental value are shown at the bottom of each plot. It can be seen that the 2CFO approach provides an excellent and much improved description of the experimental data, rendering yields within one standard deviation of the experimental measurements for most particle species. The 2CFO treatment all but eliminates the tension between light and strange baryons, the so-called proton anomaly, seen in the 1CFO approach. It should also be noted that alternative methods to treat interactions in the SHM via the S-matrix approach [176] impact in particular the proton yields and improve the performance of the 1CFO method in the Heidelberg-GSI fits [38].

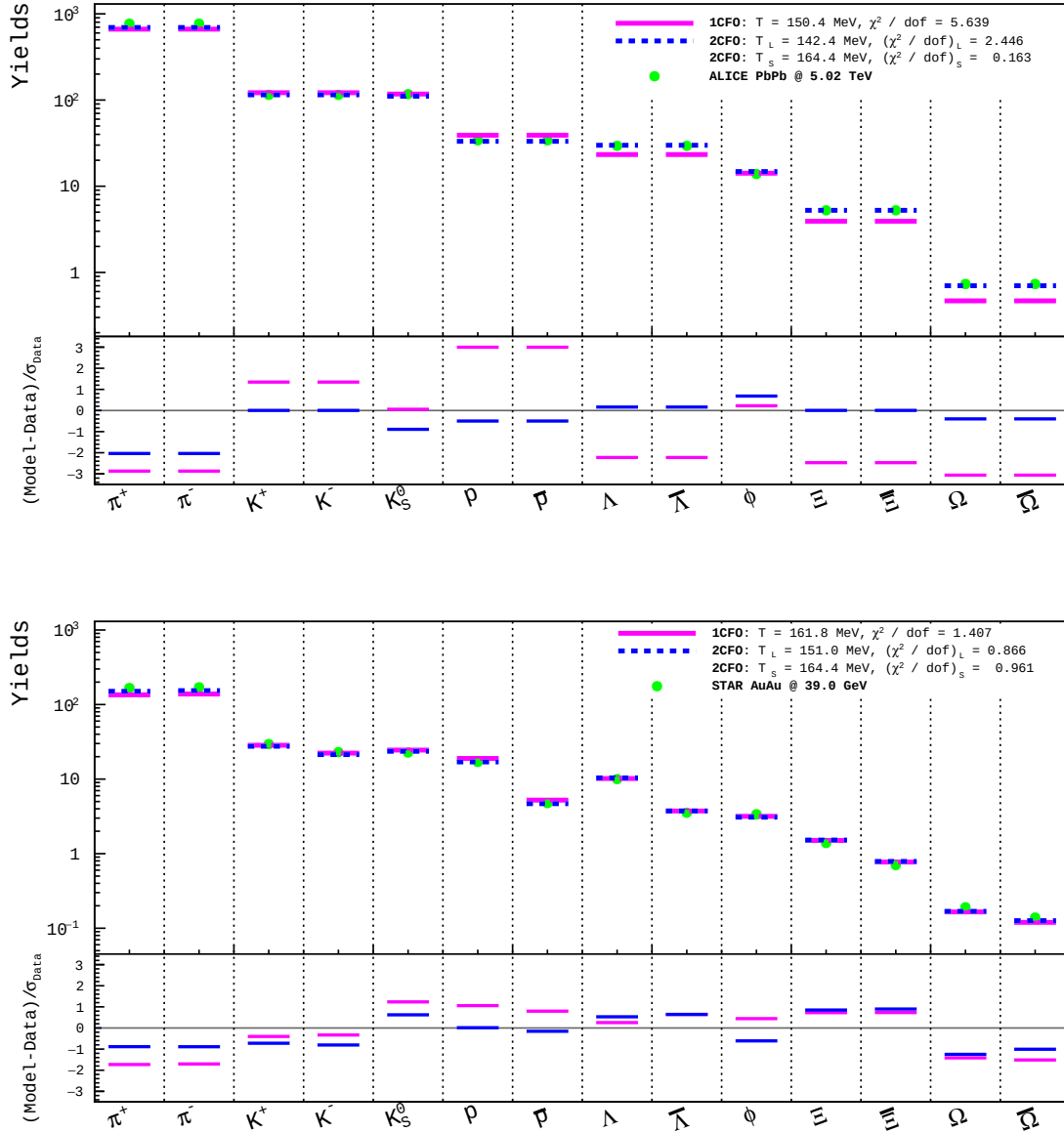


Figure 5.4: Grand Canonical Ensemble fits to ALICE PbPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV (0 - 10%) (top) and STAR AuAu collisions at $\sqrt{s_{NN}} = 39.0$ GeV (0 - 10%) (bottom), respectively, via The FIST using the PDG2016+ hadronic spectrum. Single temperature (1CFO) yield calculations are shown in magenta. Two temperature (2CFO) yield calculations are shown in dashed blue lines. The experimental data is shown in green. Figure from [136].

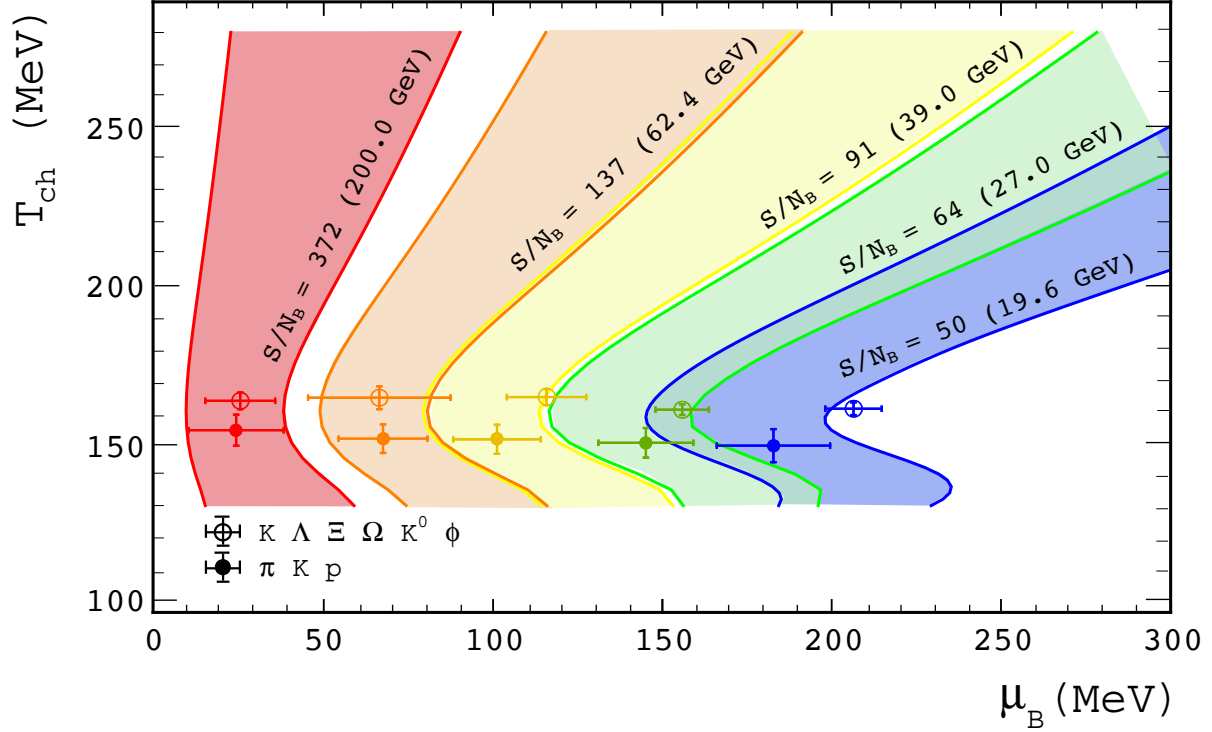


Figure 5.5: Isentropic trajectories (fixed S/N_B) in the T - μ_B plane based on a lattice QCD Equation of State expanded to finite density [177]. The closed circles show the *light* hadron freeze-out parameters, while the open points show the *strange* hadron freeze-out parameters in the 2CFO approach. The S/N_B ratios correspond to the RHIC energies 200, 62.4, 39, 27, 19.6 GeV, respectively.

As a cross-check with lattice QCD predictions, it was determined whether the extracted 2CFO parameters lie on the isentropic trajectories in the T - μ_B plane, which were calculated using a lattice QCD equation of state [177]. The freeze-out parameters are expected to lie within said trajectories in the absence of dissipative effects given that the medium generated by a heavy ion collision continues its expansion at a fixed entropy (S) value and a constant net baryon number (N_B). The validity of this approach to finite densities has been proven out to $\mu_B/T = 2$. Therefore, in Figure 5.5, the data is only shown for collision energies down to $\sqrt{s_{\text{NN}}} = 19.6$ GeV. The ALICE points are also excluded since at zero baryo-chemical potential S/N_B diverges. For all five collision energies, the light and strange freeze-out parameters lie well within the calculated trajectories. The uncertainties on the isentropic

trajectories are based on folding the errors of the light hadron freeze-out parameters for T_{ch} and μ_{B} into the calculation. A trend towards a sharper turn in the isentropic curve seems to develop with increasing density. It is likely that in order to capture said turn in future lattice QCD calculations beyond $\mu_{\text{B}}/T = 2$, it will be necessary to extend the Taylor expansion of the equation of state at finite density to terms of higher order than the c_6 coefficient [177].

5.2 System Size and Flavor Dependence of Chemical Freeze-out at ALICE

The freeze-out temperatures T_{ch} were extracted via ‘The FIST’ for the light, all, and strange particle thermal fits across increasing $\langle dN_{\text{ch}}/d\eta \rangle$ values for pp, pPb, and PbPb collision systems at ALICE in the GCE and SCE configurations using experimental particle yields.

The top panel of Figure 5.6 shows the extracted freeze-out temperatures T_{ch} as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ at $|\eta| < 0.5$ for the three different fits within the GCE treatment. The pp, pPb, PbPb points are shown as closed circles, open diamonds and closed squares, respectively. The red, green, and blue points pertain to the extracted T_{ch} values for the light, all, and strange fits, respectively. The colored bands were determined from the propagation of the error of the fits onto the freeze-out parameters. The bottom panel of Figure 5.6 shows the extracted freeze-out temperatures T_{ch} as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ at $|\eta| < 0.5$ for the three different fits within the SCE treatment, following the same labeling convention as in the top panel of the figure. Qualitatively contrasting both top and bottom panels, a deterioration of the flavor-dependent 2CFO approach is observed occurring in the pp and pPb systems when using the GCE scenario, starting at values of $\langle dN_{\text{ch}}/d\eta \rangle < 50$. This is not the case in the SCE treatment, where a flavor-dependent temperature separation is consistently present as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ across all three systems. It also is worth noting that in the grand canonical limit, i.e. for values of $\langle dN_{\text{ch}}/d\eta \rangle > 50$, both the GCE and SCE treatments render

almost identical T_{ch} values. Nevertheless, in order to qualitatively differentiate between the two treatments, an in-depth comparison of the fit quality of the aforementioned T_{ch} values as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ was performed, which is separately shown in Figure 5.7.

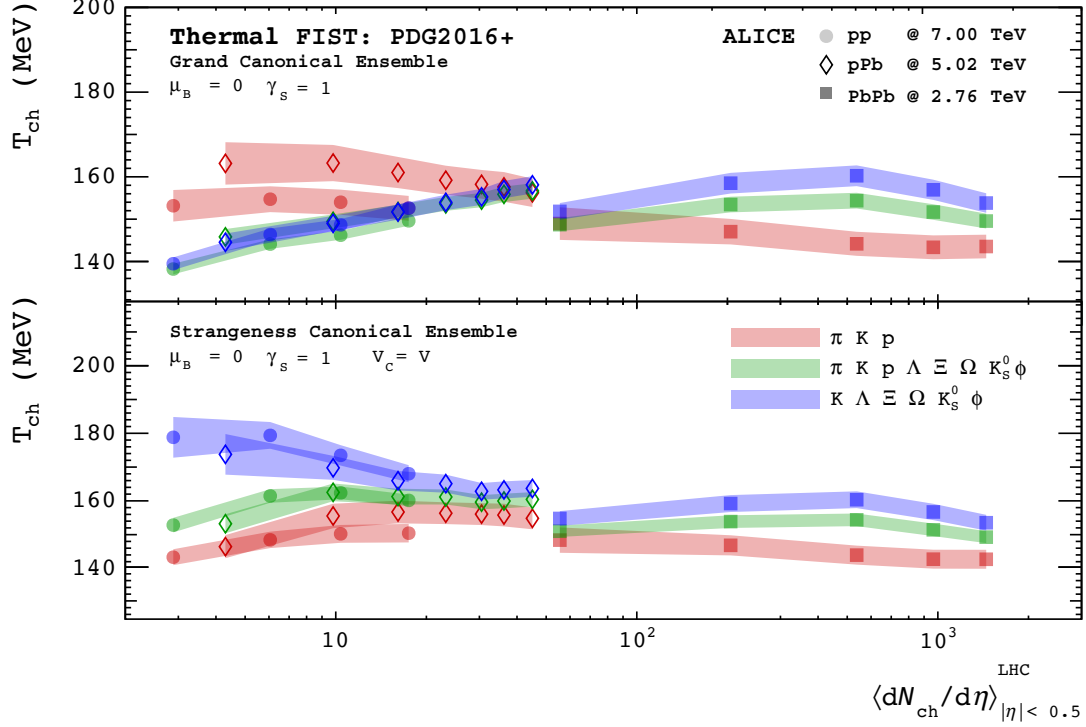


Figure 5.6: Top Panel: Flavor-dependent ‘The FIST’ Grand Canonical Ensemble Fits to Yields using the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV, respectively, shown as as closed circles, open diamonds and closed squares as a function of $\langle dN_{\text{ch}}/d\eta \rangle$. For all fits, $\mu_B = 0$ and $\gamma_S = 1$, while T_{ch} , and V were used as free parameters. Red, green, and blue points represent the light, all, and strange fits, respectively. Bottom Panel: Flavor-dependent ‘The FIST’ Strangeness Canonical Ensemble Fits to Yields using the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV as a function of $\langle dN_{\text{ch}}/d\eta \rangle$. For all Strangeness Canonical Ensemble fits, $\mu_B = 0$, $\gamma_S = 1$ and $V_C = V$, while T_{ch} , and V were used as free parameters. The bottom panel follows the same labeling convention used in the top panel. Figure from [137].

The top panel of Figure 5.7 depicts the corresponding χ^2/dof values for the GCE and SCE all fits of the top and bottom panels of Figure 5.6, shown in cyan and magenta, respectively. The pp, pPb, PbPb points are shown as closed circles, open diamonds and closed squares.

In the large system (PbPb), a consistent quality of fit for all multiplicity classes in the GCE and SCE scenarios is present, with almost identical values for both approaches. However, in the GCE case, a deteriorating quality of fits in the small systems (pp and pPb) can be seen, with $\chi^2/\text{dof} > 5$ for nearly all fits to pp data. This renders an inconsistent description of temperatures across multiplicity within the GCE treatment of smaller systems. Instead, using the SCE approach, a consistent quality of fit is attained, specifically, $\chi^2/\text{dof} \leq 5$, across all small and large systems with corresponding values $\langle dN_{\text{ch}}/d\eta \rangle \geq 10$.

The bottom panel of Figure 5.7 shows the $\Delta(\chi^2/\text{dof})$ values corresponding to the SCE fits performed in the light and strange fits from the bottom panel of Figure 5.6 – shown in orange and purple, respectively. As in the top panel, the pp, pPb, PbPb points are respectively shown as closed circles, open diamonds and closed squares. The $\Delta(\chi^2/\text{dof})$ values for the orange points in the bottom panel of Figure 5.7 were calculated as the difference between the corresponding χ^2/dof values of the SCE all (green) fits and the SCE light (red) fits from the bottom panel in Figure 5.6. Similarly, the $\Delta(\chi^2/\text{dof})$ values for the purple points in the bottom panel of Figure 5.7 were calculated as the difference between the corresponding χ^2/dof values of the SCE all (green) fits and the SCE strange (blue) fits from the bottom panel in Figure 5.6. This procedure was done in order to explicitly discriminate between notable differences in the quality of the fits when employing the 2CFO approach within the SCE, where any improvements of the fit quality would render $\Delta(\chi^2/\text{dof})$ values approximately equal to the magenta χ^2/dof values shown in the top panel. Conversely, any $\Delta(\chi^2/\text{dof})$ values below zero reflect a worsened fit quality. For all three systems, across all values of $\langle dN_{\text{ch}}/d\eta \rangle$, the calculated $\Delta(\chi^2/\text{dof})$ values for both the light (orange) and strange (purple) points generally follow the same trend as the all (cyan) points in the top panel. Without loss of generality, this suggests an overall improvement of the fits when employing the flavor-dependent 2CFO approach in the SCE treatment.

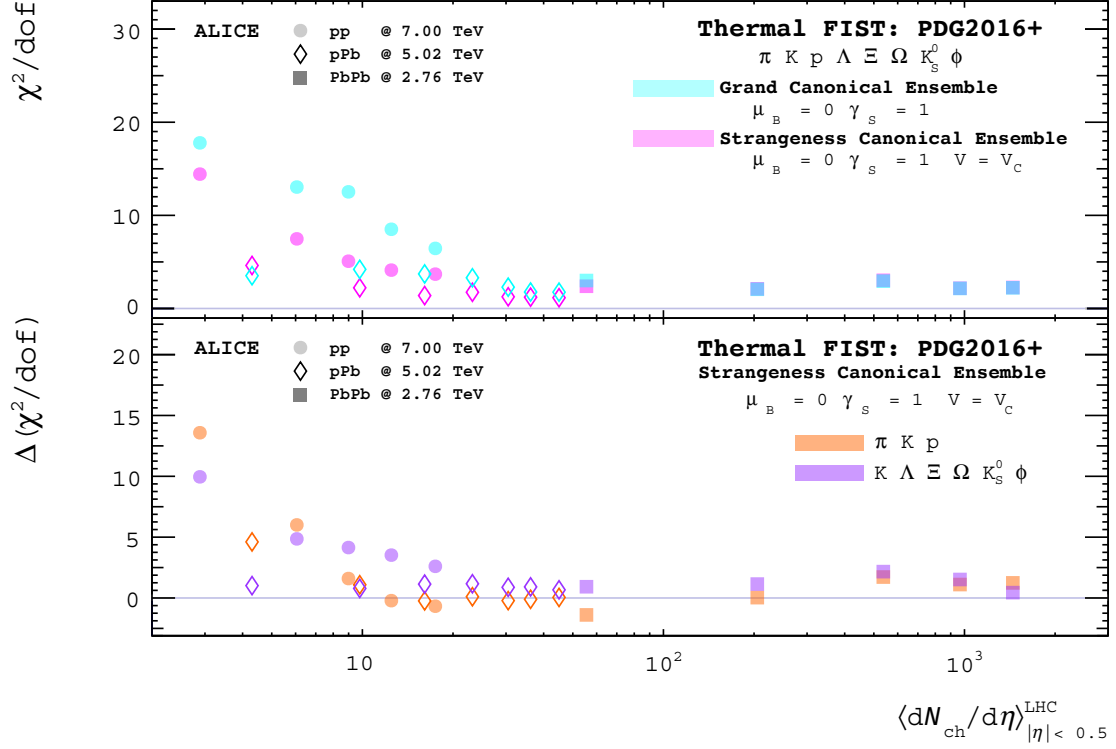


Figure 5.7: Top Panel: Flavor-dependent ‘The FIST’ χ^2/dof values of the Grand Canonical Ensemble (cyan) and Strangeness Canonical Ensemble (magenta) fits used in the all fits of the top and bottom panels of Figure 5.6, respectively, using the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV as a function of $\langle dN_{\text{ch}}/d\eta \rangle$. Bottom Panel: Flavor-dependent ‘The FIST’ $\Delta(\chi^2/\text{dof})$ values for the Strangeness Canonical Ensemble fits used in the light (orange) and strange (purple) fits of the bottom panel of Figure 5.6 via the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV as a function of $\langle dN_{\text{ch}}/d\eta \rangle$. For both the top and bottom panels, the corresponding χ^2/dof and $\Delta(\chi^2/\text{dof})$ values for pp, pPb and PbPb collision systems are respectively shown as closed circles, open diamonds and closed squares. Figure from [137].

The presence of negative $\Delta(\chi^2/\text{dof})$ values can only be seen for a few of the light fits and can be attributed to a decreasing number of degrees of freedom, which causes an overall increase in the calculated χ^2/dof . Moreover, the final χ^2/dof values for the flavor specific fits can be obtained by subtracting the values in the lower panel from the magenta points in the upper panel of Figure 5.7. One can thus observe a consistent quality of fit – specifically,

$\chi^2/\text{dof} \leq 5$ – for each of the three types of fits across the small and large systems with corresponding values of $\langle dN_{\text{ch}}/d\eta \rangle \geq 10$. It is noted that the quality of fit is consistent for both the light and strange fits for values of $\langle dN_{\text{ch}}/d\eta \rangle < 5$. Generally, the χ^2/dof values for the light and strange fits are lower than those for the all fit.

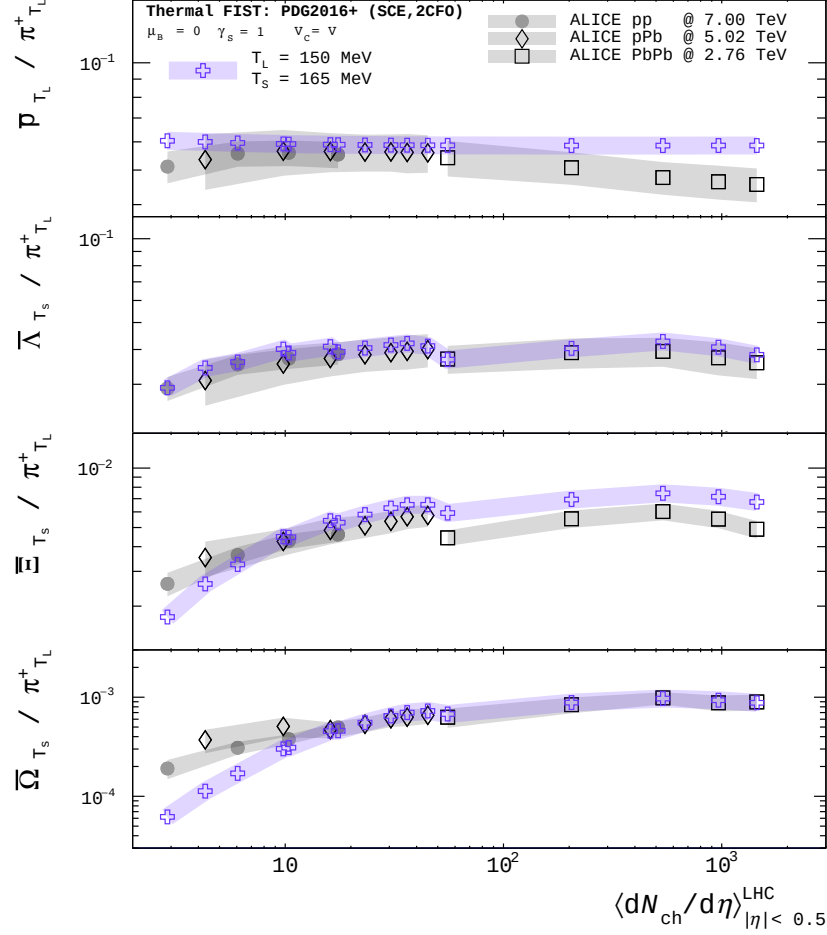


Figure 5.8: Two Chemical Freeze-out Temperature Strangeness Canonical Ensemble ‘The FIST’ Thermal Model anti-hadron to π^+ yield ratio calculations via the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV, and PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV as a function of $\langle dN_{\text{ch}}/d\eta \rangle$. From top to bottom: $\bar{p}/\pi^+$, $\bar{\Lambda}/\pi^+$, $\bar{\Xi}/\pi^+$, and $\bar{\Omega}/\pi^+$, respectively. ALICE experimental points for pp, pPb and both PbPb collision systems are respectively shown as grey closed circles, black open diamonds and black closed squares. Purple open crosses depict to the calculated Two Chemical Freeze-out Temperature anti-hadron to π^+ ratios at vanishing baryo-chemical potential using the Strangeness Canonical Ensemble framework based on the flavor specific temperatures extracted in [136]. For all calculations, $\mu_B = 0$, $\gamma_s = 1$ and $V_C = V$. Figure from [137].

Figure 5.8 shows the anti-hadron to π^+ ratio as a function of $\langle dN_{\text{ch}}/d\eta \rangle$, starting from the top panel down to the bottom: $\bar{p}/\pi^+$, $\bar{\Lambda}/\pi^+$, $\bar{\Xi}/\pi^+$ and $\bar{\Omega}/\pi^+$, respectively. The experimental points for ALICE pp, pPb and both PbPb collision systems are shown as grey closed circles, black open diamonds and black closed squares, respectively. The purple open crosses indicate the calculated 2CFO anti-hadron to π^+ ratios within the SCE framework in ‘The FIST’ using the flavor specific temperatures $T_L = 150$ MeV for π^+ and $\bar{p}$, and $T_S = 165$ MeV for $\bar{\Lambda}\bar{\Xi}$ and $\bar{\Omega}$. The corresponding model calculations of the ϕ to π^+ ratio are shown separately in Figure 5.10, keeping the same purple cross color convention, for the traditional treatment of net-strangeness of the ϕ meson of $S = 0$. For this analysis, the chemical temperatures extracted in [136] were chosen and used for central (0 - 10%) PbPb collisions at $\mu_B = 0$ across all centralities and system sizes.

The obtained results show an excellent agreement with the experimental yield ratios measured by ALICE across all three systems and are consistent with those shown by [171, 172]. Early calculations [178] requiring a factor accounting for non-equilibrium strangeness production, i.e. $\gamma_S < 1$, can be excluded based on this procedure and results from Ref. [172]. At these high collision energies, strangeness seems saturated even in the smallest systems. The main differences to Ref. [172] emerge from the fact that the presented results do not require different rapidity windows for pions and strange particles and no additional normalization factor was used to reproduce the experimental results. In this analysis, the accurate representation of final state particle yields relies solely on the use of flavor-dependent temperatures and fireball volumes – the relevant volumes as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ as determined by the model are shown in Figure 5.9.

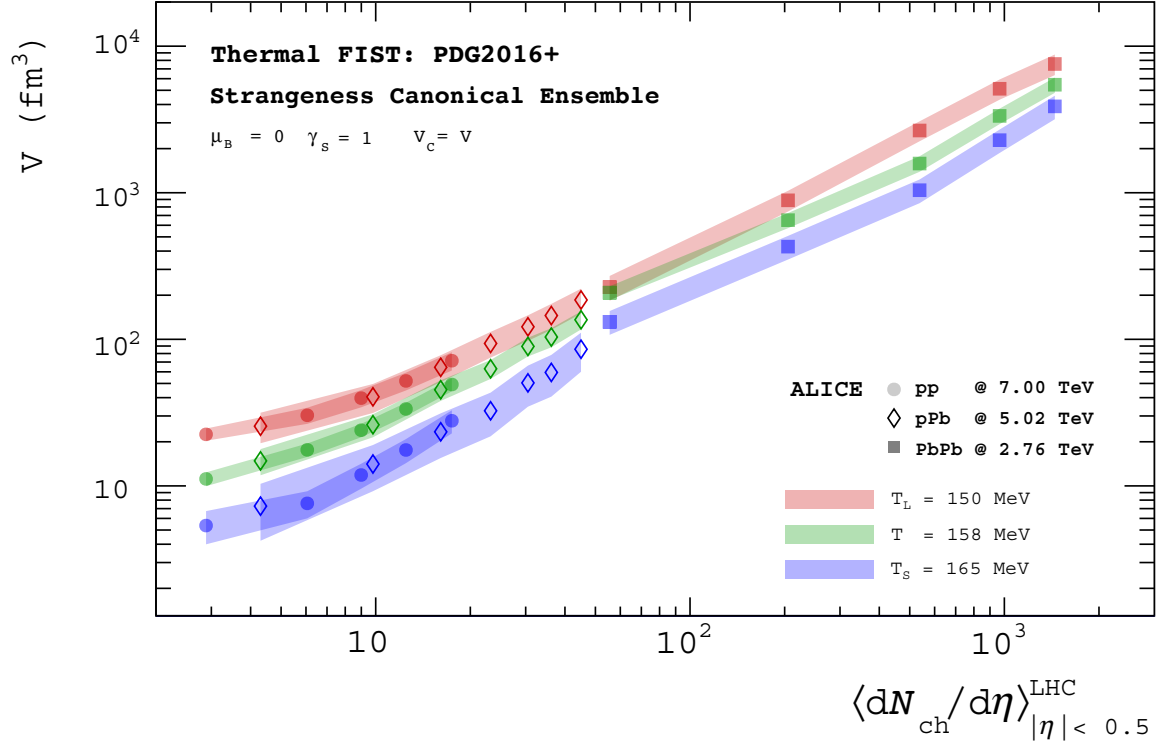


Figure 5.9: ‘The FIST’ Strangeness Canonical Ensemble Thermal Model temperature-dependent volume per unit rapidity calculations via the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and PbPb collisions at $\sqrt{s_{NN}} = 2.76$ TeV, respectively shown as as closed circles, open diamonds and closed squares as a function of $\langle dN_{ch}/d\eta \rangle$. The red, green, and blue points represent the volume calculations based on the flavor specific temperatures extracted in [136] at 150 MeV, 158 MeV and 165 MeV, respectively. For all calculations, $\mu_B = 0$, $\gamma_s = 1$ and $V_C = V$. Figure from [137].

Lastly, the treatment of net-strangeness of the ϕ meson was varied in the model between $S = 0$, $S = 1$ and $S = 2$ to gauge the sensitivity of the calculated yields and to elucidate the question of strangeness enhancement in the case of the vector meson. Figure 5.10 shows the ϕ to π^+ ratio as a function of $\langle dN_{ch}/d\eta \rangle$, following the same labeling convention for experimental points and 2CFO thermal model calculations as Figure 5.8, using an ad hoc variation of the total strangeness (S) value within the model for the ϕ meson. To reiterate, the π^+ and ϕ values were calculated at $T_L = 150$ MeV and $T_S = 165$ MeV, respectively. For the comparison to the data, the total strangeness of the ϕ meson was arbitrarily fixed to

$S = 0$ (purple open crosses), $S = 1$ (cyan open crosses) and $S = 2$ (magenta open crosses), respectively.

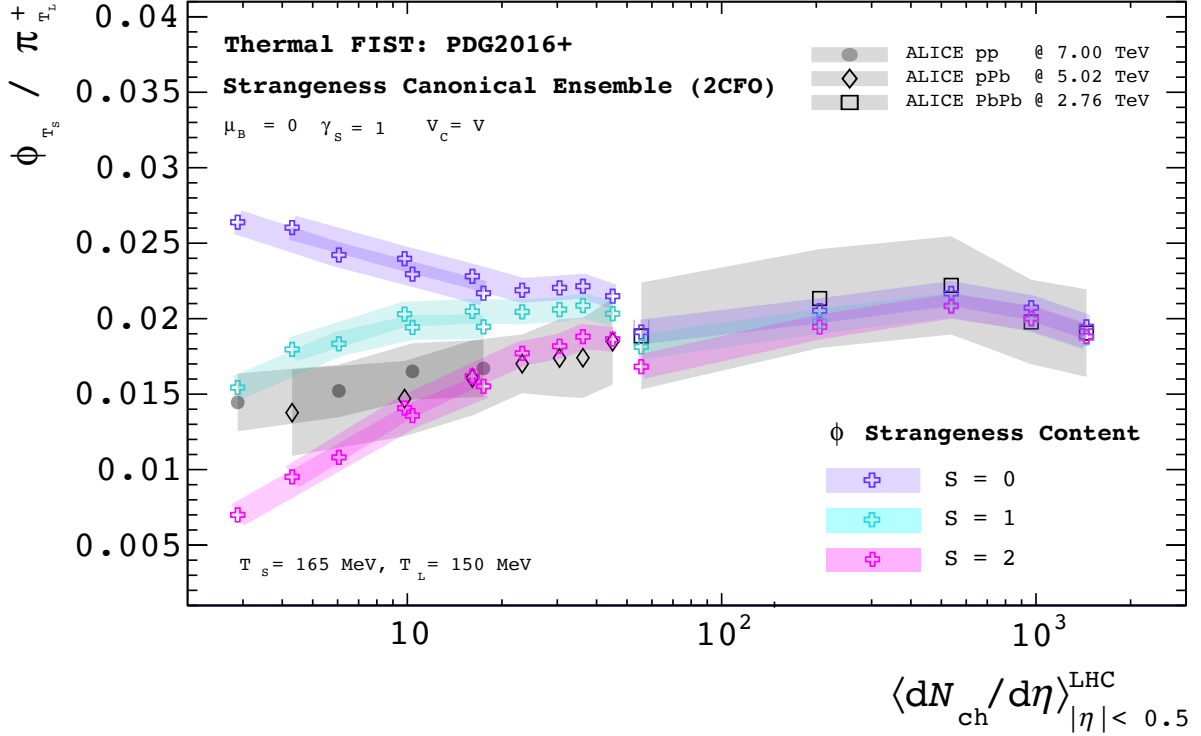


Figure 5.10: Two Chemical Freeze-out Temperature Strangeness Canonical Ensemble ‘The FIST’ Thermal Model ϕ to π^+ yield ratio calculations via the PDG2016+ hadronic spectrum for ALICE pp collisions at $\sqrt{s} = 7.00$ TeV, pPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and PbPb collisions at $\sqrt{s_{NN}} = 2.76$ TeV as a function of $\langle dN_{ch}/d\eta \rangle$. ALICE experimental points for pp, pPb, PbPb collisions are respectively shown in grey as closed circles, open diamonds, and closed squares. Purple, cyan and magenta open crosses depict the calculated SCE 2CFO ϕ to π^+ ratios in the SCE framework based on the flavor specific freeze-out temperatures at vanishing baryo-chemical potential shown in [136] fixing the total strangeness values for the ϕ meson to $S = 0$, $S = 1$ and $S = 2$, respectively. For all calculations, $\mu_B = 0$, $\gamma_S = 1$ and $V_C = V$. Figure from [137].

In the case of $S = 0$, the calculated yield ratios are well described by the 2CFO approach only in the large systems, for $\langle dN_{ch}/d\eta \rangle > 50$. The model vastly overestimates the value of the ratio for both the pp and pPb systems. In the case $S = 1$, an improvement is seen in the smaller systems, with the calculated values almost falling within the errors of

the experimental data for all multiplicity bins. Lastly, for the case of $S = 2$, the model underestimates the experimental yield ratios for values of $\langle dN_{\text{ch}}/d\eta \rangle < 10$, but does otherwise quite well. The obtained results suggest that for ϕ production in small systems the ϕ should not be considered a $S = 0$ particle, because simple flavor conservation and recombination arguments require more than a single string to fragment to form an $s\bar{s}$ state. Therefore, ϕ yields can be more accurately calculated, within the HRG framework, by assuming a strangeness content between 1 and 2.

6 CONCLUSION AND OUTLOOK

Calculations of the chemical freeze-out temperature (T_{ch}) based on particle yields from STAR and ALICE measured at collision energies ranging from $\sqrt{s_{\text{NN}}} = 11.5$ GeV to 5.02 TeV were presented. Based on a splined fit to the extracted thermal fit parameters at all energies, a light flavor freeze-out temperature $T_{\text{L}} = 150.2 \pm 2.6$ MeV and a strange flavor freeze-out temperature $T_{\text{S}} = 165.1 \pm 2.7$ MeV were calculated at vanishing μ_{B} , employing the GCE approach within the framework of the ‘The FIST’ HRG model package. Evidence for flavor-dependent chemical freeze-out temperatures in the crossover region of the QCD phase diagram was shown. The temperatures start to converge above $\mu_{\text{B}} \simeq 300$ MeV. The flavor-dependent two temperature approach via ‘The FIST’ was employed in order to successfully model and reproduce experimental yields at top ALICE energies. Thus, at the highest energies at RHIC and the LHC a separation of the hadronization temperature of light and strange particles seems likely. Furthermore, the acquired results suggest the existence of a critical point in the QCD phase diagram above $\mu_{\text{B}} \simeq 300$ MeV and a temperature below $T_{\text{ch}} \simeq 150$ MeV.

Additionally, determinations of freeze-out temperatures T_{ch} for the light, full, and strange particle thermal fits across increasing $\langle dN_{\text{ch}}/d\eta \rangle$ values for pp, pPb, and PbPb collision systems at ALICE in the GCE and SCE configurations from experimental particle yields via ‘The FIST’ using the PDG2016+ hadronic spectrum were employed. Moreover, thermal model anti-hadron to π^+ yield ratio calculations in these same collision systems as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ were also shown, with particular attention given to the treatment of the total strangeness content of the ϕ meson. In the scope of the Strangeness Canonical Ensemble within the framework of the Thermal FIST HRG model package, an excellent description of experimental yield ratios across all three systems was presented, measured by the ALICE Collaboration as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ at LHC energies, when employing flavor-dependent

chemical freeze-out temperatures under the assumption of fully saturated strangeness. In conclusion, the flavor-dependent T_{ch} separation established in heavy ion collisions, seems to prevail also at low $\langle dN_{ch}/d\eta \rangle$ values corresponding to the pp and pPb systems. This sustained separation may also be seen as an indication of QGP formation in small systems.

By providing evidence for the existence of a strange flavor hierarchy at chemical freeze-out, the presented results provide a narrow foothold into the examination of a generalized flavor hierarchy present in the cross-over region of the QCD phase diagram. An expansion of the presented methodology into the heavy flavor sector, particularly charm, is indeed a necessary and logical progression for providing further evidence (or lack thereof) for the existence of such a hierarchy.

Given that charm quarks are generated in the early stages of relativistic heavy ion collision, they propagate through the produced medium and the entire evolution of the collision system as it expands and chemically freezes out – the number of produced charm quarks is assumed to stay constant until hadronization is reached. In a heavy ion collision, the produced medium (i.e. QGP) provides a plethora of deconfined lighter (e.g. up, down and strange) quarks with which charm quarks can recombine and generate final state charmed hadrons; these are in turn detected by the experimental apparatus. Hence, one can generate a picture in which deconfined charm quarks within the QGP thermalize and subsequently freeze out at the phase transition of the QGP. This thermal-like behavior displayed by charmed hadrons allows for their production to be modeled in the context of statistical hadronization and allows existing thermal models to be expanded into the charm sector; albeit not based on the initial charm yield, which is modeled through the addition of a fugacity factor, but rather on the thermal freeze-out behavior of final state charmed hadrons. In principle, this makes it possible to broaden the aforementioned sequential strangeness freeze-out treatment by induction into the charm sector to produce a generalized model for sequential flavor

freeze-out.

As it currently stands, statistical hadronization models are capable of performing thermal fits to the measured yields of produced charm hadrons in relativistic heavy ion collisions. Open-source models, like ‘The FIST’ and THERMUS [179, 180] are more than capable of performing such calculations and possess the capacity of employing a charm analog to the SCE approach (i.e. a Charm Canonical Ensemble). The hadronic spectrum of THERMUS code itself has been recently updated with a large proportion of charm and bottom hadrons and hadronic resonances in order to undertake such a task. Coupled with the fact that the upgrade to the ALICE detector will provide a statistically robust sample size of measured charm and multi-charm hadrons in the upcoming high-luminosity runs at the LHC, investigations of a freeze-out scheme beyond two flavor-dependent temperatures can most certainly be conjectured.

Assuming the presented flavor-dependent effect on the extracted chemical freeze-out temperatures can be attributed to the mass of the strange quark itself, being close to that of the hadronization temperature, one could argue that such an effect on the freeze-out temperature should be more prominent in the charm sector due to the higher charm quark mass. In this fashion, the flavor dependence of the chemical freeze-out temperature can be interpreted as being generated by the individual masses of the quarks composing final state hadrons produced in heavy ion and elementary particle collisions.

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