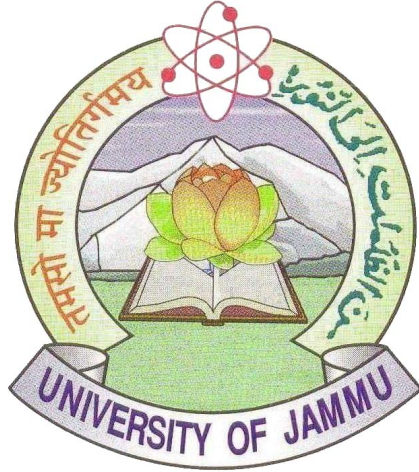

K_S^0 , Λ and $\bar{\Lambda}$ production in p–Pb collision system
with the ALICE detector at LHC energies



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SUBMITTED TO THE UNIVERSITY OF JAMMU

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by

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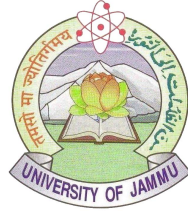
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My Spiritual Master

POSTGRADUATE DEPARTMENT OF PHYSICS,
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CERTIFICATE

It is certified that **Ms. Meenakshi Sharma**, has carried out the work presented in this thesis entitled “ **K_S^0 , Λ and $\bar{\Lambda}$ production in p–Pb collision system with the ALICE detector at LHC energies**” under my supervision and the work is worthy of consideration for the award of Ph.D. degree in Physics. It is further certified that:

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I, Meenakshi Sharma, D/O Sh. Subash Chander, hereby declare that the work reported in this thesis entitled “ **K_S^0 , Λ and $\bar{\Lambda}$ production in p–Pb collision system with the ALICE detector at LHC energies**” is entirely done by me under the supervision of Prof. Anju Bhasin, Department of Physics, University of Jammu. The work undertaken by me in this thesis is original and has not been copied from other sources without due acknowledgement. I also ascertain that no part of this thesis is presented elsewhere for the award of any degree or diploma of any University or Institution. Moreover this thesis is having 5% degree of plagiarism as per the Ouriginal analysis report.

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Meenakshi Sharma

ABSTRACT

The goal of heavy-ion collisions at ultra-relativistic energies is to study the properties of strongly interacting systems, which are described by the theory of Quantum Chromodynamics (QCD). It was predicted by QCD that a new state of matter, the so-called the Quark Gluon Plasma (QGP) can be created in such collision systems, if the temperature and energy density exceed a certain threshold, obtained by increasing the kinetic energy of the colliding beams. In this new state of matter, the constituent partons, namely quarks and gluons, are freed from nucleons in which they are normally confined by the strong force. Over the past two decades a panoply of measurements made in heavy ion collisions at the two highest energy Colliders : Relativistic Heavy Ion Collider (RHIC) and the Large Hadron Collider (LHC), combined with theoretical advances from calculations made in variety of frameworks, have led to a broad and deep knowledge of the properties of hot QCD matter. One of these discoveries is the formation of the Quark Gluon Plasma which is created through high energy heavy-ion collisions.

These “Little Bangs” create exploding little droplets of the hottest matter seen anywhere in the universe since it was a few microseconds old. Through these high energy heavy-ion collisions at these accelerators, we have increasingly quantitative descriptions of the phenomena manifest in these collisions and some of the key material properties of matter created in these collisions. In particular, we have determined that the quark-gluon plasma (QGP) created in these collisions is a strongly coupled liquid as a near-ideal Fermi Liquid with the lowest value of specific viscosity ever measured (almost no frictional resistance or viscosity). Experimental evidence also points towards collective motion of particles during QGP expansion and it was observed that the mean free path of partons (quarks and gluons) is comparable to inter particle spacing. The picture of quark-gluon plasma resembles a jigsaw puzzle, with many

pieces provided by the different experiments. The data from any one experiment alone is insufficient to provide a complete picture but the results from different experiments taken together are consistent. Despite the fact that a quark-gluon plasma explanation appears to be the most logical explanation, it is crucial to research this newly created matter at both higher and lower temperatures in order to properly characterise its properties and firmly establish the quark-gluon plasma interpretation. With this aim, the LHC facility at CERN has a design goal of 5.5 A TeV for Pb–Pb collisions allowing for higher energy densities and temperatures, thereby giving an opportunity to study the new state of matter and furthermore to understand how the universe has been evolving to what it is now.

This thesis work was carried out in the context of one of the experiments at LHC namely the ALICE (A Large Ion Collider Experiment). The main physics goal of the experiment is the creation and the investigation of the properties of the strongly-interacting matter in the conditions of high energy density and high temperatures, expected to characterize the medium formed in central heavy-ion collisions at these energies. Under these conditions, quark confinement into colorless hadrons should fade and a deconfined Quark-Gluon Plasma (QGP) should be formed. ALICE, with the current available energies and in the next future, will open a door in a whole and completely unexplored new regime for the physics of the strong interactions. The detectors of ALICE at the LHC, are especially designed to study the physics of the QGP produced in heavy-ion collisions. The extensive Particle identification (PID) capabilities of the ALICE detector allow for the study of a wide set of observables related to particle production mechanisms and sensitive to the properties and the evolution of the QGP.

The presence of a de-confined phase manifests itself with typical signatures that can be quantified by studying particle production. The relative production of strange hadrons with respect to non-strange hadrons in heavy-ion collisions was historically

considered as one of the first predicted signatures of signatures of QGP state and it was indeed observed in many experiments at SPS, RHIC and LHC. In ALICE experiment at LHC, an enhancement of strangeness production is observed in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV compared to pp collisions analysing the Run 1 data (2009–2013), thereby contributed to evidence that the QGP state had been formed. Small collision systems, such as proton-proton and proton–lead, aim to provide the reference data for collisions of heavy nuclei. However, inspection of high multiplicity pp and p–Pb interactions revealed surprising features. The latest results in proton-proton (pp) and proton-lead (p–Pb) collisions have revealed an increasing trend in the yield ratio of strange hadrons to pions with the charged-particle multiplicity in the event, showing a smooth evolution across different collision systems and energies. At the LHC, the observation of this new effect has raised a question: If we put enough energy in pp collision – can it be considered also as a small droplet of QGP? Answer to this question is a very hot topic in heavy ion physics on both sides – theory and experiment and it is still not fully understood. For the strange particle production enhancement the results at LHC suggest that there is a continuing rise in strangeness enhancement toward the high multiplicity pp collisions with saturation in heavy ion collisions. The “breaking region” where high multiplicity pp collisions and low multiplicity nucleus-nucleus collisions meet is still a region of big interest since there is lack of statistics from both data sets (pp and PbPb) and therefore room for various interpretations. To tackle this problem results from p-Pb collisions is important. Thus, main aim of this thesis is to better understand the production mechanisms for strange particles and hence the unexpected observations like strangeness enhancement phenomenon in small collision systems.

In this thesis, the results of the production of neutral strange hadrons (K_S^0 and Λ) in p–Pb collisions at a centre of mass energy of 8.16 TeV from Run 2 (2015–2018) will be presented. The transverse momentum (p_T) distributions and mean transverse

momentum, $\langle p_T \rangle$, are studied as a function of event multiplicity. An increase with multiplicity in $\langle p_T \rangle$ is observed for both the particles, indicating the presence of radial flow. The yields of the hadrons are also measured and normalised to pion yields in order to study the strangeness enhancement effect, originally predicted to indicate the presence of a QGP. A comparison to other collision energies and systems, pp, Pb–Pb and Xe–Xe, indicates that the ratio of the strange particles yield-to-pion yield follow a continuously increasing trend from low multiplicity pp to high-multiplicity Pb–Pb collisions, independent of the initial collision energy and system size. Such comparative studies will allow the energy dependence of this phenomenon to be investigated as well as a direct comparison between p–Pb and Pb–Pb results at high multiplicities. The enhancement of strangeness is clearly visible as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ as particles with larger strangeness content are relatively more enhanced (e.g. $\Omega + \bar{\Omega}$ with respect to K_S^0). This effect is purely due to the chemical composition of the hadrons under consideration as mesons and baryons with equal strangeness content show the same increase. Standard event generators such as PYTHIA, DIPSY and EPOS-LHC show discrepancies with the measured data. Modification to the usual mechanisms of particle production (such as color re-connection, color ropes, string shoving) were proposed to explain such behaviour. These mechanism show better agreement with the experimental data. The baryon-to-meson yield ratio, Λ/K_S^0 are investigated and compared for different centrality bins. An enhancement of this ratio is seen at intermediate p_T . This enhancement is, to our current knowledge, explained by effects of radial flow combined with processes like recombination during the hadronization of the created medium – the Quark Gluon Plasma.

In addition, the study of the nuclear modification factor as a function of transverse momentum (p_T) can give insight into their production mechanism. In particular, the large transverse momentum range covered by the identification techniques adopted in

the ALICE experiment provides the possibility of studying the competition between the hard mechanism (fragmentation) and soft mechanism (coalescence) in the different p_T regions.

This thesis is divided into seven chapters:

Chapter 1: The general physics context, with an introduction to QCD and the connection with the idea of the QGP, obtained at high density and temperature, is given in this first chapter. In addition a general description of the fundamental characteristics of heavy-ion collision and the time evolution of the created system are presented. Since the results which are obtained in intermediate collision system (p-Pb) will be presented so Cold Nuclear Matter Effects are also discussed in this chapter.

Chapter 2: This chapter is focused on the theoretical description of the strangeness production mechanisms (within the thermodynamical description) and the original idea of strangeness enhancement. Furthermore, results on the nuclear modification factor and the strangeness enhancement obtained in lower energy experiments are shown.

Chapter 3: Some of the recent ALICE results in small systems which are showing QGP-like features will be presented in this chapter and new novel tools to understand such results will be discussed.

Chapter 4: In this chapter the detection capabilities of the ALICE apparatus are given. Moreover, the LHC Run3 upgrades and its physics motivation will also be discussed

Chapter 5: Here, the identification technique, based on the topological reconstruction of the weak decay of the strange baryons and strange mesons, is described. The details of the several corrections which are applied to obtain the

final results are also added. The details of the total and uncorrelated systematic uncertainties obtained in this analysis for p_T differential and p_T integrated yields are also added in this chapter.

Chapter 6: This chapter includes the major results of the current analysis. The transverse momentum spectra, the results on the strangeness enhancement, baryon over meson ratio and the nuclear modification factors for the neutral strange hadrons at the LHC energy are presented and discussed. For better understanding the results, the comparison has been made with the results obtained at different energy and different collision systems.

Chapter 7: In this last chapter, the summary and interpretation of the results is included. The future perspectives of such studies will also be covered in this chapter.

Contents

List of Figures	xii
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List of Tables	xxvi
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1 Introduction	1
1.1 Nuclear Physics at High Energies	1
1.2 The Standard Model and Quantum Chromo Dynamics (QCD)	1
1.2.1 Confinement	3
1.2.2 Asymptotic Freedom	5
1.2.3 Debye Screening	6
1.2.4 MIT Bag Model	7
1.2.5 Lattice QCD	10
1.2.6 Phase Diagram of QCD	12
1.2.7 Limitations of the Standard Model	14
1.3 Time Evolution of Heavy-ion Collision	15
1.4 Quark Gluon Plasma (QGP)	18
1.4.1 Characteristics of QGP	20
1.4.1.1 QGP: A perfect liquid	20
1.4.1.2 Chiral Symmetry and Constituent Quark Masses . .	21
1.4.1.3 Deconfinement and Chiral Symmetry Restoration . .	22

CONTENTS

1.5	Signatures of Quark Gluon Plasma (QGP)	22
1.5.1	Jet quenching	23
1.5.2	Collective Flow	25
1.5.2.1	Radial Flow	25
1.5.2.2	Anisotropic flow	29
1.5.3	Direct photons and their puzzle	32
1.5.4	Quarkonium suppression (Debye screening effect)	33
1.5.5	Electromagnetic field/chirality	34
1.5.6	Vorticity and polarization	35
1.6	Modifications in the Cold Nuclear Matter	36
1.6.1	Cronin Effect	36
1.6.2	Shadowing, Anti-shadowing and EMC Effect	37
1.6.3	Color-Glass Condensate	38
1.7	Thesis Motivation	40
2	Strangeness Production in QGP	41
2.1	Probing QGP with Strangeness	41
2.2	Theoretical models	46
2.2.1	Hydrodynamic models	46
2.2.2	Microscopic models	48
2.2.3	Statistical models	49
2.2.3.1	Chemical Equilibrium ($\gamma_S = \gamma_q = 1$)	50
2.2.3.2	Chemical Semi-Equilibrium ($\gamma_q = 1, \gamma_S \neq 1$)	52
2.2.3.3	Chemical Non-equilibrium ($\gamma_q \neq 1, \gamma_S \neq 1$)	52
2.3	Production Mechanisms	54
2.3.1	Nuclear modification factors	56

3	LHC Results in Small Systems: Hints of QGP?	59
3.1	pp or p–Pb collisions: more than a reference for heavy ions	59
3.2	Experimental results in small systems	60
3.2.1	Collective effects	60
3.2.1.1	Two particle angular correlations:	61
3.2.1.2	Multi-particle correlations	64
3.2.1.3	Identified hadron spectra	66
3.2.2	Multiplicity dependent enhancement of strange and multi-strange hadron production:	69
3.2.2.1	Canonical suppression	70
3.2.3	Re-scattering effects:	76
3.3	Sources of the QGP-like effects in MC	76
4	The ALICE Experiment	79
4.1	The LHC Project	79
4.1.1	Proton Acceleration	82
4.1.2	Lead Ions Acceleration	83
4.1.3	Lead Ion Collisions	83
4.1.4	LHC Experiments	84
4.2	The ALICE Project	85
4.2.1	ALICE Coordinates	86
4.2.2	ALICE Subdetectors	88
4.2.2.1	ITS	89
4.2.2.2	TPC	91
4.2.2.3	V0 Detectors	95
4.2.3	Event Reconstruction	96
4.3	LHC Run 3 Upgrades	97

CONTENTS

4.3.1	Motivation for ITS Upgrade	99
4.3.2	Motivation for TPC Upgrade	100
4.3.2.1	Principal Operation of GEMs	101
4.3.3	Physics Motivations for Run 3	103
5	Strange hadrons identification at ALICE	105
5.1	Introduction to strange hadrons	105
5.1.1	K^0 meson and its anti-particle	105
5.1.2	K_S^0 meson and Λ ($\bar{\Lambda}$) baryon	108
5.1.3	V^0 identification	108
5.1.3.1	Topological Reconstruction	109
5.1.3.2	Invariant Mass Calculation	110
5.2	Data Sample and Event selections	112
5.2.1	Event Multiplicity Selections	112
5.3	Kinematic Cuts: Topological and Non-topological	113
5.3.0.1	Optimization of cuts	116
5.4	Signal Extraction	116
5.4.1	Statistical uncertainty	117
5.5	Signals Corrections	118
5.5.1	Feed-down Correction	118
5.5.2	Efficiency x Acceptance Correction	123
5.5.2.1	Dependence on Multiplicity	125
5.6	Systematic Uncertainties	127
5.6.1	Smoothing of systematic error	130
5.6.2	Multiplicity Dependent Systematics	133
5.6.2.1	Uncorrelated Systematics	133

6	Results and discussions	141
6.1	Transverse Momentum Spectra	141
6.1.1	p_T -integrated yields and average p_T	144
6.2	Ratio to non-strange particle yields: Strangeness Enhancement	147
6.3	Baryon Over Meson Ratio	149
6.3.1	Redistribution and enhanced baryon production	151
6.3.2	Baryon over meson ratio as a function of multiplicity	151
6.4	Nuclear Modification Factor	152
6.5	Ratio of Resonance particles to Ground-state particles	155
7	Conclusions & future perspectives	159
7.0.1	Future perspectives	160
A	Additional Plots	163
B	Lévy-Tsallis function	178
C	Geometry and centrality in heavy-ion collisions	179
C.0.1	Geometry	179
C.0.2	Centrality	181
	Bibliography	185

List of Figures

1.1	Standard model of elementary particles. Figure taken from [3]	3
1.2	Schematic view of the two possible scenarios for hadronic matter: 1) ordinary matter condition and 2) high density matter in a deconfined state. Figure taken from [4]	4
1.3	The conventional string breaking: the QCD string spanned between the static quark Q and the static antiquark $\bar{Q}$ breaks due to light $q\bar{q}$ pair creation. Figure taken from [8]	5
1.4	Calculations and measurements of the α_s coupling constant as a function of the energy scale Q . Figure taken from [10]	6
1.5	Pressure of the Hadron Gas and Quark-Gluon Plasma - Schematic of the pressure-temperature relationship for the Hadronic Gas and Quark-Gluon Plasma phase, using the MIT bag model. Figure taken from [18]	9
1.6	Lattice Predictions of Energy Density- Energy density of QCD medium as calculated with Lattice QCD. Also shown is the prediction from the ideal gas model, ε_{SB} . The critical temperature depends slightly on the number of light quark flavours considered. Figure taken from [20] . .	11
1.7	The QCD phase diagram as a function of baryon chemical potential and temperature. Credit: S Mukherjee/BNL	13

1.8	Evolution of the system created in the collision in a Minkowski-like plane[29]. Two scenarios are given: one with (left) and one without (right) QGP formation. The two freeze-out temperatures T_{Ch} (chemical freeze-out) and T_{Kin} (Kinetic freeze-out) are reported together with the critical temperature T_c	17
1.9	Sketch of the transition from normal hadronic matter to the QGP. At high enough temperature, quark pairs are created. When the quark density is large enough, the long distance strong potential-that confines the quarks in hadron-is screened by the other quarks. The quarks are deconfined in a “coloured” medium, the QGP. Figure taken from [34]	19
1.10	Pictorial view of the spontaneous breaking of the chiral symmetry. The quark, represented by the sphere, tends to occupy the minimum energy state: in the (a) configuration, it correspond to a symmetric state; in the (b) configuration, representing the confined quark in a nucleon, the parton is forced to break the symmetry and occupies an asymmetric state. Figure taken from [41].	21
1.11	Event display of dijet events in p+p (left) and Pb+Pb (right) collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV from the CMS experiment [43]	23
1.12	Nuclear modification factor R_{AA} for various particles as a function of the transverse momentum in Au+Au collisions at $\sqrt{s_{\text{NN}}} = 200$ GeV from the PHENIX experiment. Figure was taken from [44] and the references corresponding to each data are shown in the figure.	24
1.13	Schematic representation of radial flow in heavy-ion collisions (Figure adapted from [49]).	26

LIST OF FIGURES

- 1.14 Calculation of the m_T spectra for different particle species with non zero flow [46]. The effect on the spectral shape is more evident for heavier particles as the flow velocity increases. The contribution of resonance decays to the spectrum of pions is reported as $\pi^+(\text{all})$, this inclusions have important effects on the spectral shape, especially for low m_T 27
- 1.15 m_T spectra of π^- , K^- and p^- as measured at $\sqrt{s_{\text{NN}}} = 200$ GeV in central Au–Au collisions (left) and pp collisions (right) by the STAR Collaboration [47, 48]. 28
- 1.16 $\langle p_T \rangle$ for $\pi^\pm$, $K^\pm$ and $p\bar{p}$ as measured in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV and Au–Au collisions at $\sqrt{s_{\text{NN}}} = 200$ GeV [48]. A clear increase in the $\langle p_T \rangle$ is visible in central collisions. 28
- 1.17 Non-central heavy-ion collision of two nuclei in position space (left) as well as in momentum space (right). On the left, the spatial asymmetry of the participant region (colored in orange) is shown. Due to thermal pressure against the vacuum, pressure gradients within the participant region build up. The spatial anisotropy causes the pressure gradient along the x-axis to be larger than along the y-axis. The resulting momentum anisotropy (referred to as collective flow) that builds up in the reaction plane is shown in the right picture. (Figure adapted from [50].) 29

1.18	(Upper panel) v_2 as a function of the transverse momentum for $\pi^\pm$, $K^\pm$, K_S^0 , $p(\bar{p})$, ϕ , Λ ($\bar{\Lambda}$), Ω^- ($\bar{\Omega}^+$) in Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV from the STAR and PHENIX experiments [54–56]. Solid and dashed curves show the prediction of the hydrodynamic model [57]. (Lower panel) v_2 as a function of the transverse momentum for $\pi^\pm$, $K^\pm$, K_S^0 , $p(\bar{p})$, ϕ , Λ ($\bar{\Lambda}$) in semi-central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV from the ALICE experiment [58].	30
1.19	Direct photon yield (a), elliptic (b), and triangular (c) flow as a function of the transverse momentum [62] comparing to hydrodynamic model calculations with different assumptions. An effective temperature T_{eff} extracted from the inverse slope of the yield is shown in the figure. See [62] and references therein for details	32
1.20	R_{AA} of quarkonium states as a function of the binding energy in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV measured by the CMS experiment, which is consistent with the sequential suppression pattern [68–70].	34
1.21	Global polarization as a function of the collision energy for Λ and $\bar{\Lambda}$ hyperons as well as for Ξ and Ω hyperons [85, 86, 92–94]. Solid line and shaded bands denote calculations from viscous hydrodynamic model (vHLLE) with UrQMD initial condition [87] as well as a multi-phase transport (AMPT) model [88] respectively. Decay parameters α and γ of each particles [95] used in the measurements are shown in the figure. Note that the old results are corrected for recent update of the decay parameter.	36
1.22	Ratio of the nuclear structure functions for a fixed Q^2 as a function of the momentum fraction x . Plot taken from [101]	37

LIST OF FIGURES

1.23	Illustration of the parton density of protons at different energy resolutions.	38
2.1	Leading order Feynman diagrams for strangeness production.	43
2.2	(a,b) Enhancements in the rapidity range $ y < 0.5$ as a function of the mean number of participants N_{part} , showing LHC (ALICE, full symbols), RHIC and SPS (open symbols) data. Boxes on the dashed line at unity indicate statistical and systematic uncertainties on the pp or p-Be reference. Error bars on the data points represent the corresponding uncertainties for all the heavy-ion measurements and those for p-Pb at the SPS. (c) Hyperon-to-pion ratios as a function of N_{part} , for A-A and pp collisions at LHC and RHIC energies. The lines mark the thermal model predictions from [116] (full line) and [117] (dashed line).	45
2.3	Λ/K_S^0 ratios as a function of p_T measured by the ALICE Collaboration in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV, compared with Λ/K_S^0 and $\bar{\Lambda}/K_S^0$ ratios measured in Au-Au collisions at $\sqrt{s_{\text{NN}}} = 200$ GeV by STAR Collaboration. The dot-dashed line show the corresponding ratios from EPOS model. [141]	48
2.4	Experimental hadron yields and model calculations for the parameters of the best fit at the centre-of-mass energies of 7.6 GeV by the NA49 Collaboration (left panel) and 200 GeV by the STAR Collaboration (right panel) [146].	50
2.5	Integrated yields at mid-rapidity dN/dy in central (0–10%) Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV with results from a thermal fit. [147] .	51

2.6	The lower panel shows the non-equilibrium SHM fit to experimental data measured by the ALICE experiment in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV for 0–20% centrality. The upper panel shows the ratio of model values to experimental data for the three SHM variants and presents the key parameter values for: chemical non-equilibrium (solid squares), chemical semi-equilibrium (solid circles) and chemical equilibrium (solid triangles). [154]	52
2.7	Centrality dependence of the relative enhancement normalized to $\langle N_{\text{part}} \rangle$, for Ξ^- (left) and $\Omega^- + \bar{\Omega}^+$ (right), at different collision energies, obtained in the statistical approach with canonical formalism [156].	53
2.8	Illustration of the possible competition between fragmentation and coalescence at mid p_{T} . The depicted processes corresponds to: a) high p_{T} quark fragmentation into a meson, b) and c) creation respectively of a meson and baryon by recombination. [160]	55
2.9	Transverse momentum spectrum of charged hadrons $(h^- + h^+)/2$ at mid-rapidity for central Au–Au collisions produced at RHIC, measured by PHENIX (in red, 0–10% central events) and by STAR (in black, 0–5% central events). Prediction from fragmentation model (dotted black line) and recombination model (dashed blue line) and the sum of both (solid red line) are showed. [161]	56

LIST OF FIGURES

2.10	(a) Comparison of R_{AA} for unidentified particles in central Pb–Pb collisions at LHC to measurements at $\sqrt{s_{NN}} = 200$ GeV by the PHENIX and STAR experiments at RHIC. For both ALICE and PHENIX, the statistical and systematics errors are separated and represented by the lines and the boxes, respectively. For STAR the two errors are combined and represented by boxes. The vertical bars around R_{AA} equal to unity indicate the p_T independent scaling errors on R_{AA} . [164] (b) The nuclear modification factor of charged particles as a function of p_T in minimum bias (NSD) p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, compared to measurements in collisions at $\sqrt{s_{NN}} = 2.76$ TeV by the ALICE Collaboration. The errors are represented as in plot (a) for ALICE measurements. [165]	58
3.1	Schematic sketch of pp, p–Pb and Pb–Pb collisions	59
3.2	$\Delta\eta$ – $\Delta\phi$ angular correlation as measured (a) and (b) in pp collisions at $\sqrt{s} = 7$ TeV [186] for different reconstructed charged particle multiplicities (c) in p–Pb collisions at $\sqrt{s} = 5.02$ TeV [187] for the 0–20 % multiplicity class (the correlation of 60–100 % was subtracted) and (d) in most central Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [188]. One can clearly see the long range ridge at $\Delta\phi \sim 0$ and extending along $\Delta\eta$ for the events where large charged particle densities are produced	62
3.3	$\Delta\eta$ – $\Delta\phi$ angular correlation as measured in p–Pb collisions at CMS [190], ALICE [191], ATLAS [192] and LHCb [193]. One can clearly see the long range ridge at $\Delta\phi \sim 0$ and extending along $\Delta\eta$ for the events where large charged particle densities are produced	63

3.4	Second-order azimuthal anisotropy Fourier harmonics v_2 measured by CMS in pp, p-Pb and Pb-Pb collisions based on correlations among multiple charged particles [194, 195].	64
3.5	Four-particle cumulants $c_2\{4\}$ measured by ATLAS in p-Pb collisions (left) and in pp collisions (right) using different multiplicity selection methods [196].	65
3.6	Four-particle cumulants $c_2\{4\}$ measured by ATLAS in p-Pb collisions (left) and in pp collisions (right) with the “standard”, “two-subevent” and the “three-subevent” methods [197].	66
3.7	Ratios of p_T spectra of $\Lambda/2K_S^0$ measured by CMS in pp (left), p-Pb (middle), and Pb-Pb collisions (right) [177].	67
3.8	Ratios $p/\pi = (p + p)/(\pi^+ + \pi^-)$ (top plots) and Λ/K_S^0 (bottom plots) as a function of p_T in two multiplicity bins measured in the rapidity interval $0 < y_{CMS} < 0.5$. The ratios are compared to results in Pb-Pb collisions measured at midrapidity, shown in the right panels. The empty boxes show the total systematic uncertainty; the shaded boxes indicate the contribution uncorrelated across multiplicity bins (not estimated in Pb-Pb) [178].	68
3.9	Multiplicity dependence of the T_{kin} and β_T Blast-Wave [199] parameters fitted to pion, kaon and proton p_T spectra measured by ALICE in pp, p-Pb and Pb-Pb collisions [200]	69
3.10	p_T integrated yield ratios of strange and multi-strange hadrons to $\pi^+ + \pi^-$ as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta < 0.5}$ measured in the rapidity interval $ y < 0.5$ as measured in pp, p-Pb and Pb-Pb collisions [222]	70

LIST OF FIGURES

3.11 (a) $p/(\pi^+ + \pi^-)$ and $(\Lambda + \bar{\Lambda})/(2K_S^0)$ yield ratios as a function of $\langle dN_{\text{ch}}/d\eta \rangle$ as measured in pp collisions at $\sqrt{s} = 7$ TeV [228]. Baryon to meson ratios of hadrons with same strangeness show no significant increase. (b) Double ratios $h/(\pi^+ + \pi^-)$ over $h/(\pi^+ + \pi^-)$ in MB collisions[228].	71
3.12 Ratios of several particle species measured by ALICE as a function of the midrapidity pion yields for pp, p-Pb and Pb-Pb collisions compared to the THERMUS strangeness canonical suppression model prediction [202].	72
3.13 Ratios of p_T integrated particle yields K^{*0}/K , ϕ/π , ϕ/K and Ξ^-/ϕ in pp collisions at $\sqrt{s} = 13$ TeV as functions of $\langle dN_{\text{ch}}/d\eta \rangle_{ \eta < 0.5}$ [210, 211]. These measurements are compared with data from pp collisions at $\sqrt{s} = 7$ TeV [212], p-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [213, 214], and Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [215, 216]. The widths of the boxes for pp collisions at $\sqrt{s} = 13$ TeV and p-Pb collisions do not represent the uncertainties of $\langle dN_{\text{ch}}/d\eta \rangle_{ \eta < 0.5}$. The measurements for pp collisions at $\sqrt{s} = 13$ TeV are also compared to results from common event generators [217–220] and a Canonical Statistical Model calculation [212].	74
3.14 Resonance yield ratio to long-lived hadrons for pp, p-Pb, Xe-Xe, and Pb-Pb collisions, with comparisons to EPOS3 predictions and STAR data as a function of charged particle multiplicity ($\langle dN_{\text{ch}}/d\eta \rangle$). The bars and the lines represent the statistical and systematic errors, respectively[243].	75

4.1	The LHC is the last ring (dark blue line) in a complex chain of particle accelerators. The smaller machines are used in a chain to help boost the particles to their final energies and provide beams to a whole set of smaller experiments, which also aim to uncover the mysteries of the Universe [262]	80
4.2	Schematic view of ALICE Experiment	86
4.3	ALICE coordinates system and relevant variables. Adapted from [282]	87
4.4	Schematic view of ITS [289]	89
4.5	Schematic view of TPC [284]	92
4.6	(a) Specific energy loss measured by the ITS as a function of p. (b) Specific energyloss measured by the TPC as a function of rigidity for positive tracks. The curves represent parameterizations of the expected mean energy loss	93
4.7	(a) Distribution of measured V0A amplitude with NBD-Glauber model fit shown in red. The multiplicity classes are indicated by the vertical lines [292].(b) Correlation of weighted average time of flight of particles that are detected in the V0A and V0C detectors [293].	95
4.8	Event display for Pb–Pb collision recorded by ALICE during one of the first runs at centre of mass energy of 5.02 TeV per nucleon.	97
4.9	Layout of the new ITS detector.	100
4.10	(a) Close-up of GEM construction. (b)CTPC GEM stackup	102
5.1	Simulated Λ decay in ALICE	109
5.2	Kinematic Cuts - Schematic indicating some kinematic cut variables.	110
5.3	Λ Candidate’s invariant mass distribution	111
5.4	(a)Feed-down matrix for Λ , (b) Feed-down matrix for $\bar{\Lambda}$	119
5.5	(a) Feed-down spectra for Λ , (b) Feed-down spectra for $\bar{\Lambda}$	120

LIST OF FIGURES

5.6	(a) Feed-down ratio for Λ , (b) Feed-down ratio for $\bar{\Lambda}$	121
5.7	Ratio of feed-down spectra obtained using DC and MC Ratio method for Λ (a) and $\bar{\Lambda}$ (b)	124
5.8	Systematics obtained due to charged Ξ spectra used for feed-down correction.	125
5.9	(Top panel)Efficiency for K_S^0 in different multiplicity bins as a function of p_T . (Bottom panel) comparison with centrality integrated bin. . . .	126
5.10	(Top panel)Efficiency for Λ in different multiplicity bins as a function of p_T . (Bottom panel) comparison with centrality integrated bin. . . .	126
5.11	(Top panel)Efficiency for $\bar{\Lambda}$ in different multiplicity bins as a function of p_T . (Bottom panel) comparison with centrality integrated bin. . . .	127
5.12	Significance and signal losses, with respect to loosest cut considered, in low, mid and high p_T bins for the variable cosine of pointing angle.	129
5.13	Significance and signal losses, with respect to loosest cut considered, in low, mid and high p_T bins for the variable properlifetime.	130
5.14	Signal deviations seen in corrected K_S^0 spectra due to cut selection changes for the variable Properlifetime (a) and V0Daughters (b). Error bars are Roger-Barlow errors representing an acceptable deviation due to statistical fluctuations.	131
5.15	Total systematics obtained in 0-100% multiplicity class for K_S^0 (a) and Λ (b)	132
5.16	Total systematic uncertainties obtained for K_S^0 (a) and Λ (b) in three different multiplicity classes (0-10%, 10-30% and 30-100%) with roughly 1/3 of the MB statistics in each.	134
5.17	The multiplicity-uncorrelated uncertainties in comparison to the total (MB) systematic uncertatinty for K_S^0 (a) and Λ (b)	136

6.1	(Top panel): p_T spectra of K_S^0 in different V0A multiplicity classes. (Bottom panel): Ratio of the p_T spectra in different multiplicity classes to the minimum bias p_T spectrum.	142
6.2	(Top panel): p_T spectra of Λ (a) and $\bar{\Lambda}$ (b) in different V0A multiplicity classes. (Bottom panel): Ratio of the p_T spectra in different multiplicity classes to the minimum bias p_T spectrum.	143
6.3	p_T integrated yield per unit rapidity, dN/dy as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta <0.5}$ for of K_S^0 (a) and $(\Lambda + \bar{\Lambda})$ (b)	145
6.4	Average transverse momentum, $\langle p_T \rangle$ as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta <0.5}$ for of K_S^0 (a) and $(\Lambda + \bar{\Lambda})$ (b)	146
6.5	(a) $K_S^0/(\pi^+ + \pi^-)$ and (b) $(\Lambda + \bar{\Lambda})/(\pi^+ + \pi^-)$ yields as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta <0.5}$	148
6.6	(a) Λ/K_S^0 as a function of p_T in 0–5% and 60–80% multiplicity classes and two different energies, (b) Λ/K_S^0 as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta <0.5}$	150
6.7	Fit to the dependence on $\sqrt{s_{NN}}$ of the pp yields for K_S^0 for a p_T bin. The fitted function is used to interpolate the value at 8.16 TeV (red point on the curve).	152
6.8	Fit to the dependence on $\sqrt{s_{NN}}$ of the pp yields for $\Lambda + \bar{\Lambda}$ for a p_T bin. The fitted function is used to interpolate the value at 8.16 TeV (red point on the curve).	153
6.9	Reference pp spectra at 13 TeV (red marker) and 7 TeV (green marker) and interpolated spectra at 8.16 TeV (blue marker) for K_S^0	154
6.10	p–Pb spectra at 8.16 TeV (red marker) and re-binned interpolated pp spectra (blue marker) for K_S^0	154
6.11	Nuclear Modification Factor, R_{pPb} for K_S^0 and Λ as a function of p_T	156

LIST OF FIGURES

6.12	p_T integrated particle yield ratios for K^{*0} (892)/ K_S^0 and Λ (1520)/ Λ a function of $\langle dN_{\text{ch}}/d\eta \rangle_{ \eta <0.5}$. Comparison of Λ (1520)/ Λ is also made with data points from p-Pb 5.02 TeV	158
7.1	$(\Omega^- + \bar{\Omega}^+)/(\pi^+ + \pi^-)$ yields as a function of $\langle dN_{\text{ch}}/d\eta \rangle_{ \eta <0.5}$	162
A.1	invariant mass fits for K_S^0 in 0-1 % multiplicity bin in few p_T bins. . .	163
A.2	invariant mass fits for K_S^0 in 80-100 % multiplicity bin in few p_T bins	164
A.3	invariant mass fits for Λ in 0-1 % multiplicity bin in few p_T bins. . . .	165
A.4	invariant mass fits for Λ in 80-100 % multiplicity bin in few p_T bins .	166
A.5	invariant mass fits for $\bar{\Lambda}$ in 0-1 % multiplicity bin in few p_T bins. . . .	167
A.6	invariant mass fits for $\bar{\Lambda}$ in 80-100 % multiplicity bin in few p_T bins .	168
A.7	Significance and signal losses for K_S^0 , with respect to loosest cut considered, in low, mid and high p_T bins for the variables V0Radius and DCA NegToPV.	169
A.8	Significance and signal losses for K_S^0 , with respect to loosest cut considered, in low, mid and high p_T bins for the variables DCA PosToPV and TPC Crossed Rows.	170
A.9	Significance and signal losses for K_S^0 , with respect to loosest cut considered, in low, mid and high p_T bins for the variable DCA V0Daughters.	171
A.10	Signal deviations seen in corrected K_S^0 spectra due to cut selection changes. Error bars are Roger-Barlow errors representing an acceptable deviation due to statistical fluctuations.	172
A.11	Total Systematics obtained in 0-100% multiplicity class for $\bar{\Lambda}$	173
A.12	Total systematic uncertainties obtained in three different multiplicity classes (0-10%, 10-30% and 30-100%) with roughly 1/3 of the MB statistics each for $\bar{\Lambda}$	173

A.13	The multiplicity-uncorrelated uncertainties in comparison to the total (MB) systematic uncertainty for $\bar{\Lambda}$	173
A.14	(a) K_S^0 , (b) Λ and (c) $\bar{\Lambda}$ p_T spectra fitted with different functions to compute systematics due to extrapolation.	174
A.15	(a) K_S^0 , (b) Λ and (c) $\bar{\Lambda}$ raw p_T spectra in different multiplicity classes	175
A.16	p_T integrated yield per unit rapidity, dN/dy as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta <0.5}$ for of K_S^0 and $(\Lambda + \bar{\Lambda})$ including data points from Pb–Pb collisions . .	176
A.17	Average transverse momentum, $\langle p_T \rangle$ as a function of $\langle dN_{ch}/d\eta \rangle_{ \eta <0.5}$ for of K_S^0 and $(\Lambda + \bar{\Lambda})$ including data points from Pb–Pb collisions . .	177
C.1	Top Left: Two heavy ions before the collision with the impact pa- rameter b . Top Right: The spectators remain unaffected while in the participant zone, particle production takes place - UrQMD calcula- tion [335]. Bottom: An illustration of the often used nomenclature. Three collisions in the transverse plane are shown (from left to right): peripheral, semi-central, and central	180
C.2	Distribution of the sum of amplitudes in the V0A scintillators for p–Pb collisions at $\sqrt{s_{NN}}$ 8.16 TeV. The distribution is fitted with the NBD- Glauber fit shown as a line. The inset shows a zoom of the most peripheral region.	182

List of Tables

1.1	Dataset collected at RHIC	16
1.2	Dataset collected at LHC.	16
4.1	PID techniques provided by the ALICE barrel and single arm detector subsystems.	94
4.2	Decay topology PID methods implemented by the ALICE barrel de- tector subsystems.	94
5.1	Selections applied to K_S^0 , Λ and $\bar{\Lambda}$	138
5.2	Values of the selection cuts applied for systematics estimation of K_S^0	139
5.3	Values of the selection cuts applied for systematics estimation of Λ ($\bar{\Lambda}$)	140

Chapter 1

This chapter is an introduction to the concepts of the Quantum Chromo Dynamics (QCD) and Quark-Gluon Plasma (QGP) and how it can be created in high energy heavy-ion collisions. In this chapter we will discuss how the QGP offers exceptional chances to test quantum chromodynamics under extreme conditions of temperature and density. Many signatures supporting the formation of the QGP have also been discussed.

Chapter 1

Introduction

1.1 Nuclear Physics at High Energies

Since the discovery of the radioactive processes, physicists have had the chance to look into the underlying principles that govern how matter behaves. The curiosity of the scientists of the late 19th and early 20th centuries propelled their quest to broaden the boundaries of human knowledge. During those times, the present notion of a nucleus was created. Since then, there have been increasingly more theoretical and experimental advancements, with a bigger influence on technological advances and on the course of history. The observation of “compelling evidence for the presence of a new state of matter in which quarks, instead of being linked up into more complicated particles such as protons and neutrons, are liberated to roam freely” was made public by CERN in February 2000 [1].

1.2 The Standard Model and Quantum Chromo Dynamics (QCD)

The Standard Model of particle physics describes two categories of fundamental particles: quarks and leptons, which can interact with each other via three funda-

CHAPTER 1. INTRODUCTION

mental forces: electromagnetic, weak, and strong. The forces are transmitted via force carriers. In particular, in the case of an electromagnetic interaction, photons are transferred between particles, bosons $W^\pm$ or Z^0 carry the weak force, and gluons carry the strong interaction. There are six quarks and six leptons (and their anti-particles). Quarks can be identified based on their flavour and there are six flavours: up, down, strange, charm, bottom and top. We distinguish the following leptons: electron (e^-) and electron neutrino (ν_e), muon (μ^-) and muon neutrino (ν_μ), and finally tau (τ^-) and tau neutrino (ν_τ). In Figure 1.1, all these fundamental particles are listed. This figure also includes one more particle, which is the Higgs boson, responsible for giving mass to the particles. The forces are distinguished based on their characteristic behaviour. Between particles with an electric charge there exists the well-known electromagnetic force. The weak force is responsible for the nuclear decays as it has the ability to change the flavour of quarks and leptons. Finally, only particles containing a colour charge, such as quarks and gluons, are subject to the strong force.

Quantum Chromodynamics (QCD) [2], which derives its name from the colour charges of particles impacted by the strong interaction, is the theory that explains the strong force. There are three charges in particular: **red**, **green**, and **blue**. Each quark can only have one colour charge, while a gluon has two charges. In contrast to photons, which do not have an electric charge and are therefore prohibited from interacting with one another, gluons can interact with themselves since they carry colour charge. The strength of the force can be quantified by a coupling constant α_{QCD} . At large distances (or low energies), the α_{QCD} increases, while at small distances (or large energies), the α_{QCD} approaches asymptotically to zero. This is in contrast to how the well-known electromagnetic force behaves, which is strongest at short distances and quite weak at long distances. Colour confinement [5] and asymptotic freedom [6] are

1.2. THE STANDARD MODEL AND QUANTUM CHROMO DYNAMICS (QCD)

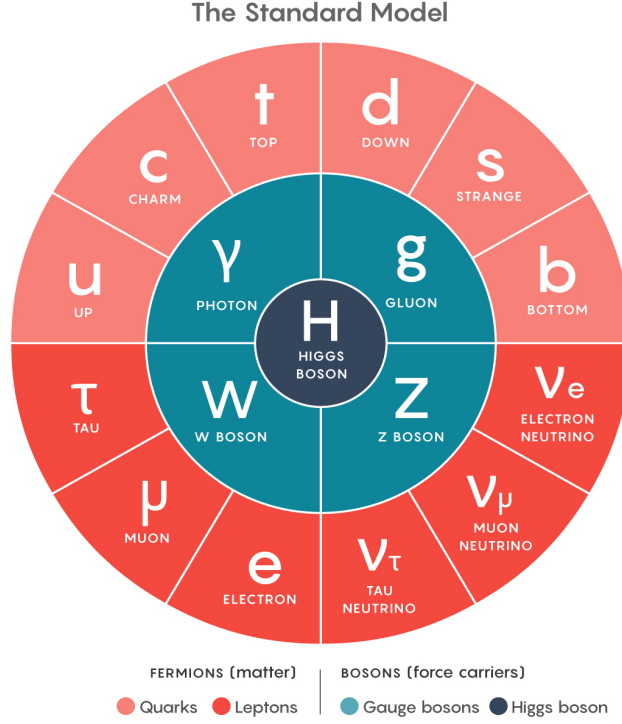


Figure 1.1: Standard model of elementary particles. Figure taken from [3]

the two primary characteristics of the strong interaction as a result of the properties of coupling constant properties in QCD.

By raising their energy, the electrons of atoms can be separated from their nucleus. This can be performed by increasing the temperature of a gas or by applying strong electromagnetic fields. In principle a similar effect can be achieved for hadronic matter by reaching the asymptotic freedom regime [7]. The only practical way to free the hadron constituents is to drastically raise the energy scale of the system through compression or heating, shown schematically in Figure 1.2.

1.2.1 Confinement

The potential in quantum electrodynamics is inversely related to the charge separation, r ,

$$V_{QED}(r) \propto -\frac{\alpha_{em}}{r}, \quad (1.1)$$

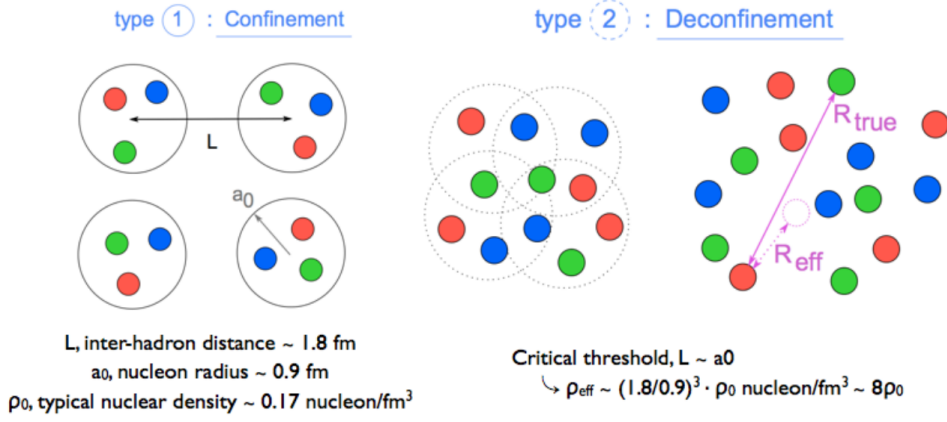


Figure 1.2: Schematic view of the two possible scenarios for hadronic matter: 1) ordinary matter condition and 2) high density matter in a deconfined state. Figure taken from [4]

where α_{em} is the electromagnetic coupling constant. However, in quantum chromodynamics, the behaviour is different with a strong potential. As two colour charges are pulled apart, the gluons self-interact and the energy in the field creates more gluons, which pull the field lines together as schematically depicted in Figure 1.3. This behaviour adds an additional ‘string-like’ term (proportional to the ‘string’ constant, k), to the potential giving the Cornell potential for QCD

$$V_{QCD}(r) \propto -\frac{4}{3} \frac{\alpha_s}{r} + kr, \quad (1.2)$$

where α_s is the strong coupling constant [5]. The second term dominates as the colour charges are further separated, and the equation predicts that the force will continue to grow infinitely. Instead of further separating the quarks, it actually becomes energetically advantageous to break the string and produce a new quark-antiquark pair from the gluon field as the energy level rises. This results in the confinement of quarks in hadrons, making it unable to examine single quarks experimentally.

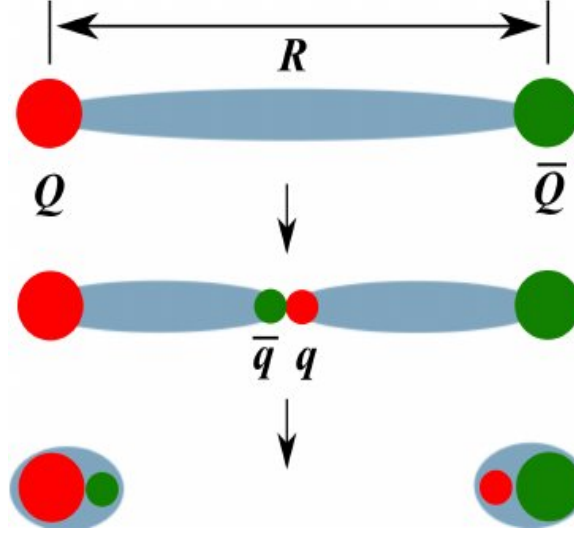


Figure 1.3: The conventional string breaking: the QCD string spanned between the static quark Q and the static antiquark $\bar{Q}$ breaks due to light $q\bar{q}$ pair creation. Figure taken from [8]

1.2.2 Asymptotic Freedom

Due to vacuum fluctuations permitted by the Heisenberg Uncertainty Principle, the coupling ‘constants’ in both QED and QCD actually change depending on the separation of bare charges. An electron in quantum electrodynamics (QED) is surrounded by electron-positron pairs which screen the charge of the electron. The screening is reduced and the electron’s apparent charge appears to be larger as another charge is brought closer to it. This effect is represented by the coupling constant, α_{em} increasing. Reduced distances signify strong momentum transfers in the interaction, so the strength of the coupling constant is usually described in terms of the four momentum, Q . For the QED coupling constant, this is described by

$$\alpha_{em}(Q) = \frac{\alpha}{1 - \frac{2\alpha}{3\pi} \ln(Q/m_e)}, \quad (1.3)$$

where $\alpha = \frac{e^2}{4\pi}$ and m_e is the mass of the electron[9]. In QCD, a similar effect is seen however, a colour charge will be surrounded by virtual gluon loops as well as quark-

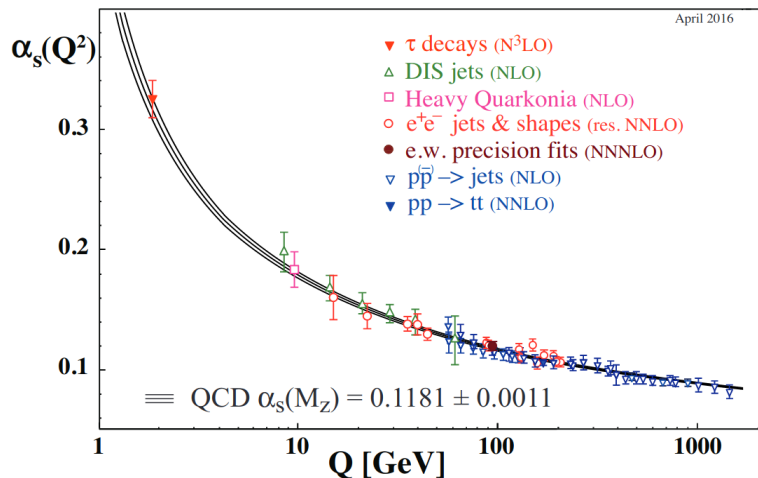


Figure 1.4: Calculations and measurements of the α_s coupling constant as a function of the energy scale Q . Figure taken from [10]

antiquark pairs. These gluon loops have an anti-screening effect that counteracts the screening of the quark-antiquark pairs. The strength of the coupling constant is described by

$$\alpha_s(Q^2) = \frac{\alpha(Q_0^2)}{1 + \frac{11N_f - 2n_f}{12\pi} \alpha(Q_0^2) \ln(Q^2/Q_0^2)}, \quad (1.4)$$

which gives the value of α at Q^2 relative to a value measured at Q_0^2 [6]. The factors N_f and n_f are given by the number of colour charges (three) and the number of quark flavours (six) respectively. With these values, the anti-screening effect of the gluons dominates and as the colour charges are brought closer together, or Q^2 increases, the anti-screening is reduced and the strength of the charge decreases as can be seen in Figure 1.4, which illustrates the running of the QCD coupling constant. This is known as asymptotic freedom and was discovered by David Politzer, David Gross and Frank Wilczek [11–13], earning them the Nobel prize in 2004.

1.2.3 Debye Screening

A plasma is a system with a high density of mobile charges. The polarised charges that separate them will screen out any forces that might exist between any

1.2. THE STANDARD MODEL AND QUANTUM CHROMO DYNAMICS (QCD)

two charges. It is frequently seen in electromagnetic plasmas, but can also be used to model the deconfinement of quarks at lower energy scales than asymptotic freedom. When the distances separating charges is reduced significantly, long range interactions are screened so that only short-range interactions need to be taken into account. This alters the QCD potential to

$$V_{QCD}(r) = -\frac{4}{3} \frac{\alpha_s}{r} e^{-r/\lambda_D}, \quad (1.5)$$

where λ_D represents the Debye screening length [14]. The Debye screening length shortens as the matter density rises as a result of heating or compression. When λ_D becomes comparable to hadronic radii at ≈ 1 fm, the forces between the quarks become smeared by nearby colour charges and the hadrons melt into a quark medium where quarks and gluons continue to interact vigorously but cannot be bonded into hadrons.

1.2.4 MIT Bag Model

The MIT bag model is a QCD-inspired model which can describe basic properties of hadrons and incorporates the confinement of quarks inside the hadron. It was developed by a group hosted at Massachusetts Institute of Technology in the early seventies [15, 16]. This model, is a basic hadron model from 1974 that still includes the fundamental physics [17]. According to this model, a bag is a region of space where the valence quarks of a specific hadron are thought to be localised. The quarks have no interaction potential within the bag; instead the confining force is represented by an external pressure B . The quarks cannot escape from the bag because of these boundary conditions. The Hamiltonian of the model can be written as

$$H_{bag} = H_{kinetic} + H_{spin-spin} + H_{volume} + H_0,$$

where the volume term is simply given by BV . Here B is the bag constant and V the

bag volume, which is in the easiest case only the volume of a sphere $V = \frac{4}{3}\pi R^3$. Since Lorentz invariance must be preserved, the bag constant B , which stands in for the external pressure, is a scalar. The prediction of the temperature where a phase change would be expected between hadronic matter and the quark-gluon plasma would be instructive in evaluating the feasibility of creating this matter. By comparing the model's predictions with the observed hadron mass spectra, the parameter B may be fixed and is found to be around $B^{1/4} \approx 200$ MeV [18].

Although this model may appear unduly straightforward, it has one significant benefit: because the quarks in the bag are characterised as non-interacting particles, one can roughly describe their behaviour as that of an ideal gas. Similar to this, a gas of hadrons can be thought of as non-interacting on the scales of the strong force. The pressure and energy density of an ideal massless gas are given by [19]

$$P = g \frac{\pi^2}{90} T^4, \quad \epsilon = 3g \frac{\pi^2}{90} T^4, \quad (1.6)$$

where P and ϵ are the pressure and energy density of a gas of temperature T , and g is the number of degrees of freedom. Although initially assumed to be three pions for the hadron gas, the number of degrees of freedom will change with temperature when alternative hadronic states or quark flavours become accessible. There are 8 bosonic gluons in a QGP, each with two spin states, and there are two types of fermionic quarks of interest, each with three colour states, two spin states, and one antiparticle. The total number of degrees of freedom for a QGP is thus equation 1.5

$$g_{QGP} = 2 \times 8 + \frac{7}{8}(2 \times 3 \times 2 \times 2) \sim 37, \quad (1.7)$$

where the factor of $\frac{7}{8}$ reflects the Fermi-Dirac statistics of the quarks[19].

The bag pressure must also be considered for the QGP; this reduces the overall

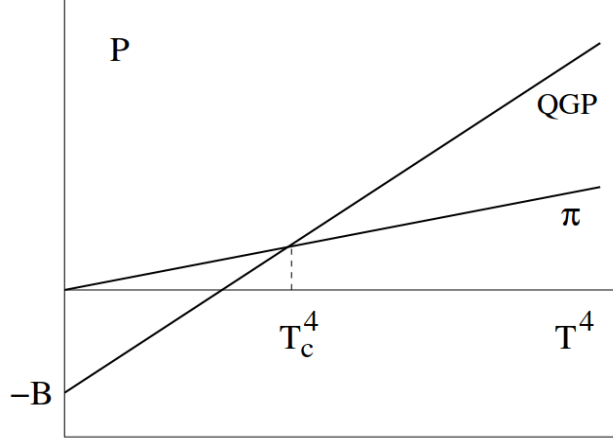


Figure 1.5: Pressure of the Hadron Gas and Quark-Gluon Plasma - Schematic of the pressure-temperature relationship for the Hadronic Gas and Quark-Gluon Plasma phase, using the MIT bag model. Figure taken from [18]

pressure and increases the energy density by B . In a mixed state, the state with the higher pressure will crowd out the other state, and so an approximation is reached for the transition temperature between the hadron gas (HG) and QGP phases by setting

$$P_{HG} = P_{QGP} \quad (1.8)$$

$$3 \frac{\pi^2}{90} T^4 = 37 \frac{\pi^2}{90} T^4 - B \quad (1.9)$$

Figure 1.5 provides a graphic representation of this. The critical temperature (T_c), which occurs at about 150 MeV, is determined by the temperature at which the two curves cross. Note that this temperature justifies our limiting of the degrees of freedom for the hadron gas to the three pions. It is likewise of the same order as the bare strange quark mass; including it marginally increases the degree of freedom of the QGP and decreases the critical temperature without significantly altering the physics. The research using Lattice QCD can help us understand its impact; these studies are discussed below.

1.2.5 Lattice QCD

In lattice QCD, the Lagrangian associated with the theory is discretised onto a finite space-time grid. Figure 1.6, which depicts the predicted energy density as a function of temperature based on these calculations, is drawn from the detailed description of this process in [20]. The exact details of the prediction vary depending on how many light quarks are taken into account, but for a plausible model consisting of two light quarks and a strange quark with a mass ~ 170 MeV, the critical temperature comes out to be around 170 MeV. To start, it should be said that this is fairly similar to the forecast made by the rudimentary model in the preceding section 1.2.4. At this temperature, there is a rapid increase in energy density that is consistent with a phase transition to a QGP state, which has a much higher degree of freedom. The energy density prediction from the ideal gas model, ε_{SB} is also indicated. When the temperature is high enough to effectively overwhelm the contact forces, the system should gravitate toward this value. This does not appear to be the case within an order of magnitude of the transition temperature, and would suggest that truly free QCD behaviour requires temperatures much higher than those achievable in modern experiments. Any QGP developed in the lab would have lingering interactions, which would allow for behaviour like hydrodynamics. The diagram shows the equivalent temperatures predicted by this theory as well as the energy densities measured at various accelerators. These represent fixed target experiments at the Super Proton Synchrotron (SPS) at CERN, and higher energy collider experiments at RHIC. A prediction is also shown for the LHC. The Bjorken formula [21], which connects the energy density at thermalization to the transverse energy density in the detector, is used to estimate the energy density.

Overall, the lattice QCD calculations suggest, firstly, that a phase transition at around $T_c \sim 170$ MeV is to be expected, and that collisions of heavy ions at RHIC,

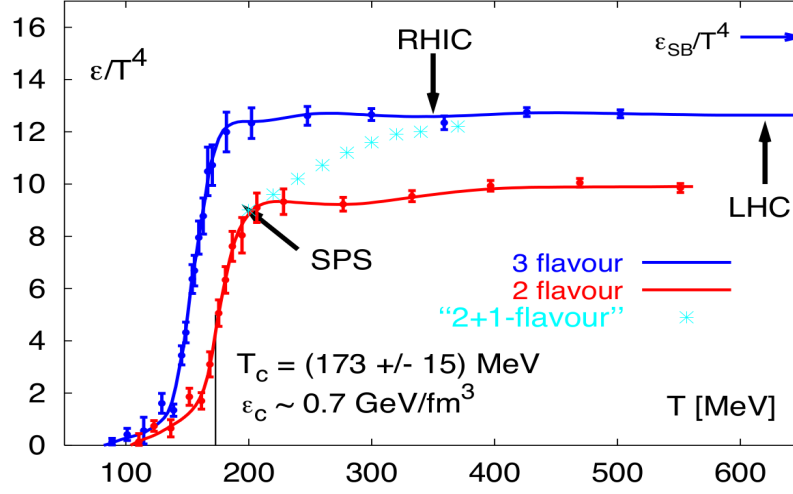


Figure 1.6: Lattice Predictions of Energy Density- Energy density of QCD medium as calculated with Lattice QCD. Also shown is the prediction from the ideal gas model, ϵ_{SB} . The critical temperature depends slightly on the number of light quark flavours considered. Figure taken from [20]

and particularly at the LHC, should reach temperatures well over this transition point

Another important Lattice QCD prediction is that the mass of quarks in a QGP is effectively lower than in the QCD vacuum [22]. Since low energy QCD is non-perturbative, it is impossible to approximate the gluon clouds surrounding any actual quark using perturbation theory. The theory is instead renormalized, where fundamental parameters are expressed in terms of experimentally measured properties. This leads to effective quark masses in the vacuum known as dynamical masses; approximately 300 MeV for up and down quarks, and around 500 MeV for strange quarks [23]. As the strong coupling drops, the effective quark masses also drop, and in a QGP they reach their bare masses; a few MeV for the up and down quarks and around 100 MeV for the strange quark [24]. This is important for the behaviour of QGP, as it may speed up the chemical equilibrium; the statistically most likely distribution of quark flavours in the medium. Particularly, strangeness equilibrium should be attained in a QGP more quickly than in a hadron gas given that the strange quark mass is now of the order of the QGP phase transition.

1.2.6 Phase Diagram of QCD

Lattice QCD computations were used in numerous investigations to shed light on the dynamics of the QCD phase. Only recently the technical and theoretical advancements allowed to better constrain the characteristics of this phenomenon. The phase change is anticipated to occur when the energy density surpasses the critical value ϵ_c , which is of the order of $1\text{GeV}/\text{fm}^3$ based on recent lattice QCD calculations [25].

The phase diagram in Figure 1.7 can be explained as follows:

1. The phase transition can happen at low temperatures but extreme nuclear matter densities ($\mu_B \rightarrow \infty$) are required. This scenario can be obtained by compressing cold nuclear matter and reaching density values higher 3-10 times ρ_0 ($\rho_0 \sim 0.15/\text{fm}^3$ is the baryon density at equilibrium) [26].
2. The phase transition at vanishing baryochemical potential ($\mu_B \rightarrow 0$) can only be accomplished by raising the system temperature above T_c . It is anticipated that this transition will be a “crossover.” The deconfined and confined phases would coexist close to the transition boundaries in such way that the entropy evolution as a function of the temperature ($\mathbb{S}(T)$) is continuous.
3. At intermediate μ_B and at intermediate temperatures the transition is of the first order, implying that a discontinuity in the entropy $\mathbb{S}(T)$ is located at temperature T_c . The first order transition needs a latent heat to maintain the temperature constant while the phase transition is happening.

The core of neutron stars is assumed to be able to achieve the circumstances for the first case (low T , large μ_B), where it is conceivable to greatly increase the matter density without rising the system temperature [27]. The other two scenarios are reproducible in a lab. As will be demonstrated later, high temperatures are obtained

1.2. THE STANDARD MODEL AND QUANTUM CHROMO DYNAMICS (QCD)

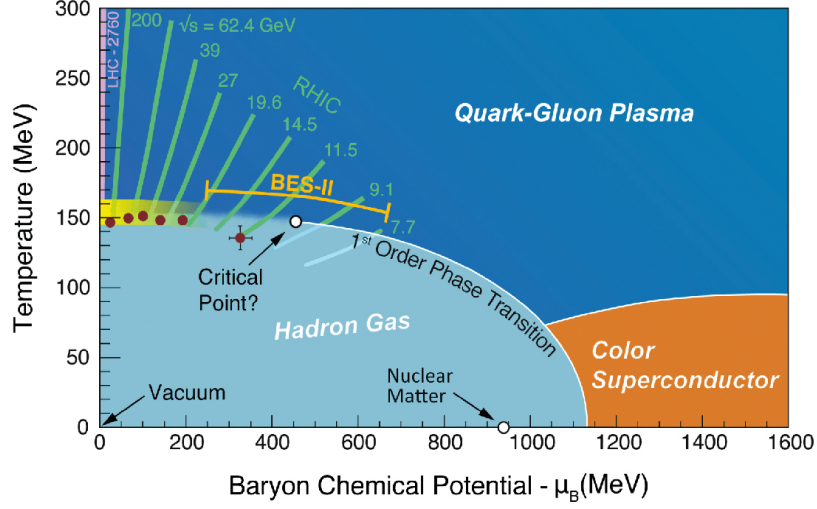


Figure 1.7: The QCD phase diagram as a function of baryon chemical potential and temperature. Credit: S Mukherjee/BNL

when nuclei collide at ultra-relativistic energies, causing the energy density to exceed ε_c . At BNL, CERN, GSI, and other facilities, these kinds of experiments have been and continue to be conducted in both collider and fix target mode. Only the type of nuclei employed and the collision energy itself affect the values of μ_B and T . The latter, in particular, identifies different regimes: starting from the 1 AGeV scale (at the SIS facility) colliding nuclei are stopped resulting in a moderated compression and heating of the nuclear matter. As the collision energy increases the temperature rises and the baryon chemical potential of the system created at rest lowers. This is because the total stopping of the colliding nuclei is less likely to occur at higher energies. The study of the QCD phase diagram has recently taken centre stage in the physics programmes of the greatest laboratories. This is currently accomplished by measuring collisions at the highest energies (LHC) or by examining the phase transition by altering the collision energy (Beam Energy Scan at RHIC). As more and more facilities devoted to nuclear physics are anticipated, this area of study will be crucial for the advancement of science.

So far the highest energy per collision was reached at the Large Hadron Collider

by using lead ions ($^{208}\text{Pb}^{82+}$) in 2015 [28] at $\sqrt{s}$ of 5.02 TeV per nucleon pair ($\sqrt{s_{\text{NN}}} = 5.02$ TeV) i.e. an overall $\sqrt{s}$ of ~ 2 PeV, which makes the LHC the first collider able to reach the PeV scale. At this energy scale the “fireball” created in the collision has no trace of the initial nuclei and is to good approximation constituted by a (net-)baryon-free hadronic matter, similar to the condition of our universe in its first moments.

1.2.7 Limitations of the Standard Model

SM is compact and many aspects of it have been experimentally verified to high accuracy. However, SM can’t explain some of the experimental observations which are listed below:

1. **Hierarchy problem:** The huge difference in the strength of fundamental forces is one aspect of the “hierarchy problem”. It also refers to the wide range in mass for the elementary particles. We know that electron is about 200 times lighter than the muon and 3500 times lighter than the tau. Same thing for the quarks: the top quark is 75000 times heavier than the up quark.

The hierarchy problem is also related to the Higgs boson mass. It is expected that quantum loop correction in the Higgs boson propagator contributes to the Higgs boson mass. Approximately it can be written as $m_H^2 \approx m_{bare}^2 + \Delta m_h^2$, where m_H is measured Higgs mass, m_{bare} is a negative free parameter, and Δm_h is loop correction term which is proportional to high mass scale Λ . If the SM is valid up to the Planck scale ($\Lambda_{pl} \sim 10^{19}$ GeV), this loop correction inflates the Higgs mass to a very large value. But it requires some unnatural fine-tuning of SM parameter to constraint the experimental observed Higgs mass ($m_H \sim 125$ GeV) around the electroweak scale ($\Lambda_{ev} \sim 100$ GeV). This is known as the hierarchy problem.

2. **Neutrino masses:** In the SM, neutrinos are massless. However, the exper-

imental observation of neutrino oscillations between their flavor eigen states proves that they possess a non-zero mass. In the SM, it is not completely understood how neutrinos gain mass.

3. **Dark matter:** Cosmological observations like galaxy rotation curves and galaxy cluster masses indicate that a significant part of the universe’s mass budget is not observable through any known interaction. Ascribing this unobservable mass to particles, such particles are called as dark matter particles. However, none of the SM particles is a good candidate for such dark matter.
4. **Gravity:** The SM describes well three fundamental forces: the electromagnetic force, the strong force, and the weak force at the quantum level. Furthermore, the coupling constants of these forces are expected to evolve to a unified value at the Planck scale. Finding an equivalent quantum theory of gravity that fits in this framework, however, remains the outstanding problem facing the unification of all forces.
5. **Baryon asymmetry :** We observed that the universe is dominated by matter. The SM does not explain the imbalance of matter and antimatter in the Universe.

1.3 Time Evolution of Heavy-ion Collision

The heavy nuclei—such as gold at the Relativistic Heavy-Ion Collider (RHIC) and lead at the Large Hadron Collider (LHC)—are accelerated to nearly the speed of light at these two greatest hadron colliders. The details of the different collision systems and energies at RHIC and LHC are listed in table 1.1 and 1.2. The nuclei are Lorentz contracted at such high velocities, forming narrow discs (commonly referred to as “pancakes”) with high densities and significant gluon occupation. The two nuclei are

CHAPTER 1. INTRODUCTION

then brought into collision. The nuclei collide and pass through one another, leaving a system behind that cools and expands.

Table 1.1: Dataset collected at RHIC

Collision System	Year	Total particle energy [GeV/nucleon]
Au+Au	2000, 2001, 2002, 2004, 2006 2008, 2010, 2011, 2014, 2016 2017, 2018, 2019,2020, 2021	27.9, 65.2, 100, 9.8, 31.2 4.6, 19.5, 3.85, 5.75, 13.5 7.3, 27.2, 9.8, 4.59, 8.65
polarised p+p	2002, 2003, 2004, 2005, 2006, 2008 2009, 2011, 2012,2013,2015,2017,2022	100.2, 204.9, 31.2, 249.9 254.9, 254.2
d+Au	2003,2008,2016,2021	(100.7+100), (31.3+31.1) (9.9+9.8), (19.6+19.4)
Cu+Cu	2005	100, 31.2, 11.2
U+U	2012	96.4
Cu+Au	2012	99.9+100
h+Au	2014	103.5+100
polarised p+p(RHICf)	2017	254.9
Zr+Zr	2018	100
Ru+Ru	2018	100
Au + Au (Fixed Target)	2018, 2019, 2020, 2021	3.85, 26.5, 7.3, 4.59 31.2, 19.5, 13.5, 9.8 5.75, 44.5, 70, 100
O+O	2021	100
polarised p+Au	2015	103.9+98.6
polarised p+Al	2015	103.9+98.7
polarised pp2pp	2009	100.2

Table 1.2: Dataset collected at LHC.

Collision System	Year	Centre of mass energy (TeV)
Pb-Pb	2010, 2011	2.76
	2015, 2018	5.02
Xe-Xe	2017	5.44
p-Pb	2013	5.02
	2016	5.02, 8.16
p-p	2009-2013	0.9, 2.76, 7, 8
	2015, 2017	5.02
	2015-2018	13

After the collision has occurred, the fireball evolves through a series of states as depicted in Figure 1.8:

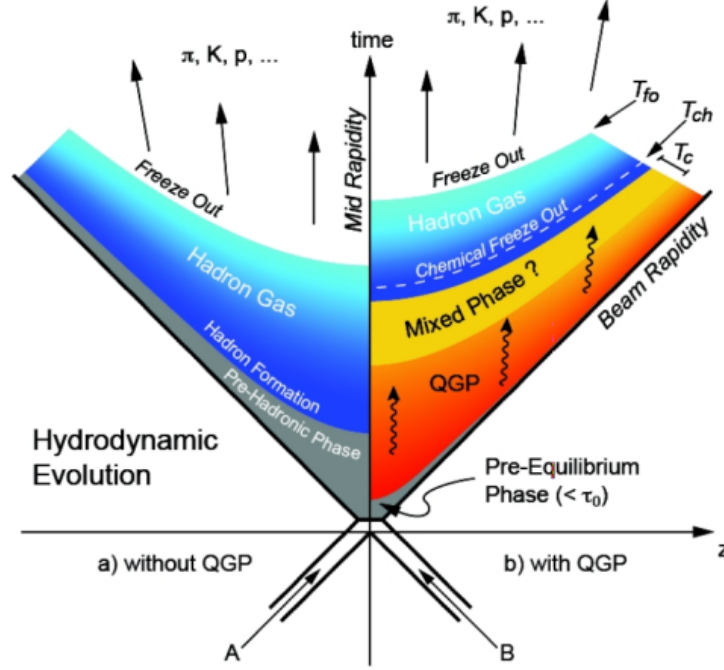


Figure 1.8: Evolution of the system created in the collision in a Minkowski-like plane[29]. Two scenarios are given: one with (left) and one without (right) QGP formation. The two freeze-out temperatures T_{Ch} (chemical freeze-out) and T_{Kin} (Kinetic freeze-out) are reported together with the critical temperature T_c .

1. **Pre-equilibrium ($t \lesssim 1 \text{ fm}/c$):** The two nuclei just collided. The partons of every participating nucleon interact producing a large amount of quarks and gluons. Hard scatterings at this stage result in partons with tremendous velocity that will subsequently coalesce into massive “hard” hadrons or fragment into jets. Additionally, prompt photons are created. The system is now formed by a dense inhomogeneous droplet of strongly interacting QGP matter.
2. **Thermalization ($t \sim 1\text{--}10 \text{ fm}/c$):** Due to the numerous interactions of each newly-created parton, the medium reaches thermal equilibrium. Additionally, the inelastic scatterings take place and alters the “chemical” composition of the medium by altering the flavor of partons. Now in equilibrium, the system develops an internal pressure that encounters no resistance from the void around it. Consequently, the system expands quickly with a drop in the temperature.

Hadrons begin to form when the temperature drops because the energy density of the system is unable to keep partons apart.

3. **Hadronization ($t \sim 20 \text{ fm/c}$):** As the energy density crosses ε_c ($T \equiv T_c$) hadronization takes over. The system loses its degrees of freedom as all partons are confined into hadrons. The interactions between hadrons keep the system in thermal equilibrium.
4. **Chemical freeze-out:** The temperature decreases down to $T \equiv T_{\text{Ch}}$, at this point all inelastic interactions among hadrons have totally ceased. The system is now formed by a gas of hadrons that interact only (pseudo-)elastically.
5. **Kinetic freeze-out:** As the temperature drops and the mean free path grows higher than the average distance between hadrons, the hadrons begin to decouple from the medium at this point. This results in a complete stop of all interactions, the temperature is $T \equiv T_{\text{Kin}}$. As a result, there is no longer any distribution of kinetic energy among the hadrons, and the momentum spectrum of each particle is fixed.

1.4 Quark Gluon Plasma (QGP)

The experimentally well-established confinement prevents quarks and gluons to be observed freely. Under typical circumstances, they are constantly bound in hadrons. But until a few microseconds after the big bang, they were released from their hadronic imprisonment. When the pseudo-critical temperature changed from higher to lower temperatures, quarks and gluons began to create hadrons. It is feasible to produce the so-called Quark-Gluon Plasma (QGP) through nuclear collisions at extremely high energy and high temperatures [30]. In general, at high enough temperatures and/or high enough densities, quarks and gluons are released from confinement, which is shown pictorially in Figure 1.9.

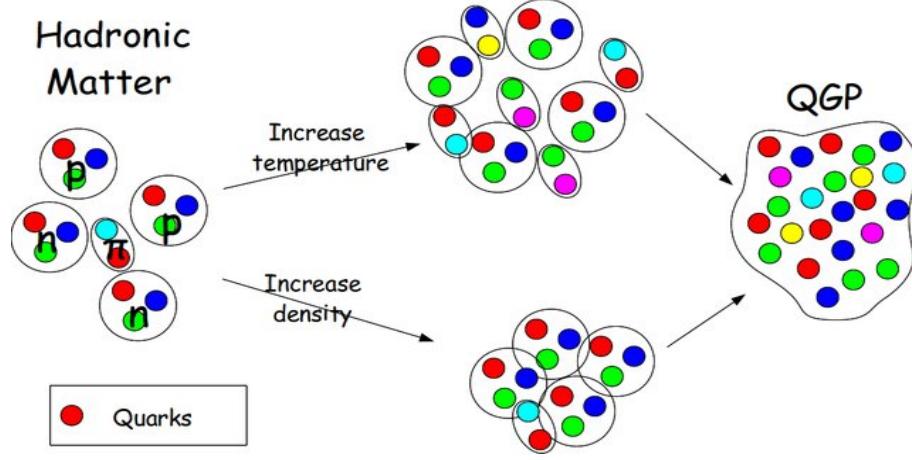


Figure 1.9: Sketch of the transition from normal hadronic matter to the QGP. At high enough temperature, quark pairs are created. When the quark density is large enough, the long distance strong potential-that confines the quarks in hadron-is screened by the other quarks. The quarks are deconfined in a “coloured” medium, the QGP. Figure taken from [34]

John Collins and Malcolm Perry theorised the existence of a new superdense state of matter made of deconfined quarks rather than hadrons that could exist in the centre of neutron stars and in the very early stages of the universe after the Big Bang because of the asymptotically free behaviour of quarks at small separation scales [31, 32]. The same year, Cabibbo and Parisi [33] proposed that only weakly interacting quark matter would exist in the asymptotically free regime.

One can get a rough estimate of a limiting scale in QCD from Heisenberg’s uncertainty principle ($\Delta p \Delta x \geq \hbar$) assuming the size of a hadron ($\approx 1\text{fm}$) as limiting length:

$$\Lambda_{QCD} = \frac{\hbar c}{1\text{fm}} = \frac{197\text{MeV fm}}{1\text{fm}} \approx 200\text{MeV} \quad (1.10)$$

In 1965, a fascinating discovery regarding the transition temperature was found. Based on the growing density of hadronic states with increasing energy, R. Hagedorn contended that there is a maximum temperature for hadronic matter [35]. He predicted an exponential behaviour for the density of states as a function of mass.

$$\rho(m) \propto (m_0^2 + m^2)^{-\frac{5}{4}} \exp(m/T_H), \quad (1.11)$$

with his calculation he observed a limiting temperature T_H of around 158 MeV. But since then the number of known resonances increased.

The Hagedorn temperature T_H was revisited lately by Cleymans and Worku [36]. They performed a fit to the hagedorn spectrum using all known resonance states from the PDG (Particle Data Group, the collaboration that gathers and summarises the field of particle physics and linked disciplines; this summary is issued every two years and is commonly known as PDG in the field) of 2008. They discovered a value for $T_c = 177.1$ MeV [36] using this. In other words, the deconfinement temperature calculated from lattice QCD and the limiting temperature of the hadron gas are of the same order.

1.4.1 Characteristics of QGP

1.4.1.1 QGP: A perfect liquid

The occurrence of a phase transition to a new state of matter had been hypothesised even before the formulation of the QCD [37] based on the mass spectrum of resonances created by hadronic collisions. This discovery has been linked to a phase change following the construction of the QCD as an asymptotically free theory [38]. The phase transition is anticipated to occur at a critical temperature $T_c = 175$ MeV, which corresponds to an energy density $\varepsilon \simeq 1$ GeV/fm³, according to QCD lattice predictions. The ultra-relativistic heavy-ion collisions generate the energy density circumstances more conducive to the formation of this phase of coloured matter known as the Quark Gluon Plasma, starting with the ordinary nuclear matter (Figure 1.7).

The QGP was predicted to behave as an ideal non-interacting gas in the region of the transition from hadronic to partonic matter, as suggested by extrapolations of the asymptotic high temperature behaviour of QCD; however, the observation

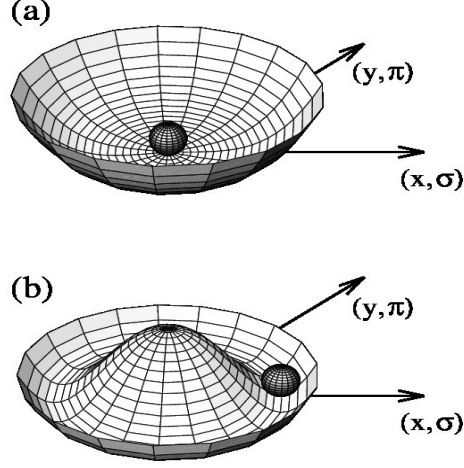


Figure 1.10: Pictorial view of the spontaneous breaking of the chiral symmetry. The quark, represented by the sphere, tends to occupy the minimum energy state: in the (a) configuration, it correspond to a symmetric state; in the (b) configuration, representing the confined quark in a nucleon, the parton is forced to break the symmetry and occupies an asymmetric state. Figure taken from [41].

made by RHIC seems to suggest a rapid thermalization of the dense matter produced in relativistic heavy-ion collisions [39]. It can be inferred from this that the hotter and denser medium produced by these collisions is still strongly interacting. Early non-perturbative studies have also demonstrated that the high temperature phase of QCD is not merely an ideal quark and gluon gas; even at temperatures $T \simeq (2-3)T_c$, significant departures from the ideal gas behaviour have been inferred [40].

1.4.1.2 Chiral Symmetry and Constituent Quark Masses

For light hadrons, only a small fraction of the hadrons masses are recovered when the masses of the quarks they are composed of are summed. This is supported by the spontaneous chiral symmetry breakdown mechanism in the QCD theory. The QCD Lagrangian becomes chirally symmetric, or invariant under separate flavour rotations of right- and left-handed quarks, if the mass term is ignored. For the up and down quarks and quite well for the strange quark, ignoring the mass term is a

decent approximation. This symmetry is broken by the non-zero vacuum expectation value of the scalar quark density operator (“chiral condensate”), which results in a dynamic mass of order 300 MeV for the up and down quarks and around 450 MeV for the strange quark. A pictorial view of the spontaneous chiral symmetry breaking mechanism can be seen in Figure 1.10.

1.4.1.3 Deconfinement and Chiral Symmetry Restoration

The partial chiral symmetry restoration that occurs during the transition from hadronic matter to the QGP has an impact on the masses of the quarks. The quarks “dress” themselves with gluons when they are enclosed in hadrons, resulting in an effective constituent quark mass of roughly 300 MeV (corresponding to 1/3 of the proton mass or 1/2 of the ρ -meson mass). The mass of the constituent quarks in the confined phase must be produced spontaneously through the confinement interaction because the fundamental bare quarks in the QCD Lagrangian are, on the other hand, nearly massless (spontaneous breaking of chiral symmetry). As a result, when deconfinement takes place, this extra mass is probably “lost” and the quarks return to their intrinsic bare mass [42]. A Lagrangian with massless fermions possesses chiral symmetry; which enables the decomposition of the quarks into independent left- and right-handed massless spin one-half states, which, for massive fermions, become mixed. Hence the mass shift transition in QCD is often referred to as chiral symmetry restoration. It is crucial to emphasise that it does not necessarily coincide with deconfinement: when the hadrons are dissolved into quark constituents, the free and now coloured quarks may still interact to form coloured bound states.

1.5 Signatures of Quark Gluon Plasma (QGP)

The QGP is not a stationary medium but subjected to dynamical evolution. The expansion of the system leads to a cooling followed by the final formation of hadrons, which are particles built from the available quarks and gluons. A detector eventually

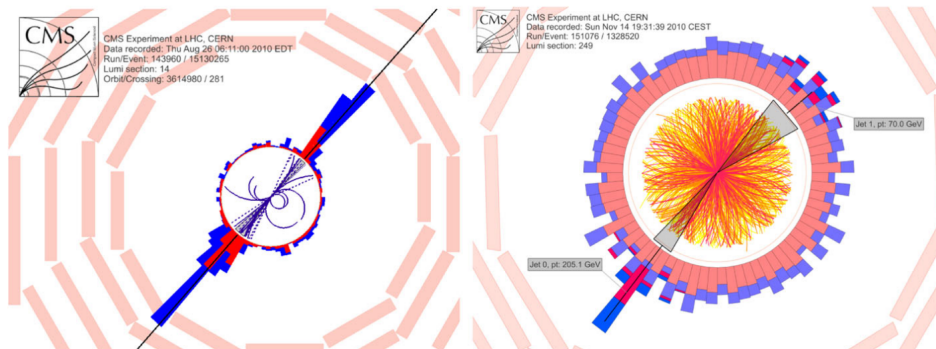


Figure 1.11: Event display of dijet events in p+p (left) and Pb+Pb (right) collisions at $\sqrt{s_{NN}} = 2.76$ TeV from the CMS experiment [43]

measures these newly formed particles, which are either made up of three quarks (baryons) or a quark and an anti-quark (mesons). In addition to the hadrons, leptons (i.e. electrons, photons) are produced. Since the whole collision evolution with a duration of 10^{-23} s is technically not possible to be followed, observables that show the properties of the medium and underlying processes over the course of the many evolution phases are essential. Some of these observables or signatures of QGP are listed below :

1.5.1 Jet quenching

In heavy-ion collisions, as already discussed the system life time is extremely short and therefore it is almost impossible to use external probes. Instead such a “tomography” can be performed by energetic and/or highly penetrating particles produced in initial parton-parton scatterings. The scattered partons have high transverse momentum and traverse the QGP losing their energy. Since the partons cannot exist on their own, they fragment into a spray of hadrons called a jet. Energy loss of partons in the medium can be understood by collisional and radiative processes similar to energy loss in QED.

Collisional energy loss is due to elastic scatterings between the initial primary parton and a parton from the medium and is expected to linearly scale with the

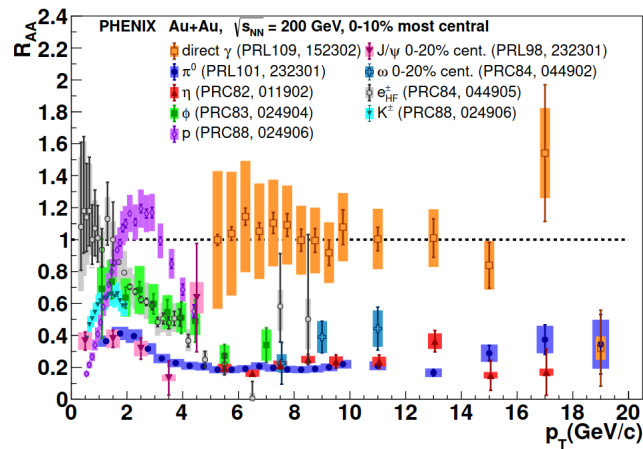


Figure 1.12: Nuclear modification factor R_{AA} for various particles as a function of the transverse momentum in Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV from the PHENIX experiment. Figure was taken from [44] and the references corresponding to each data are shown in the figure.

path length of the initial parton when traveling the uniform medium, while radiative energy loss takes place due to gluon radiation.

Figure 1.11 shows event displays of jets in proton-proton (p+p) and nucleus-nucleus (A+A) collisions. Two collinear (“back-to-back”) jets forced by momentum conservation are clearly observed in p+p collisions, while one of jets is significantly suppressed in A+A collisions: the phenomena called “jet quenching”. The jet quenching has been studied by measuring a so-called nuclear modification factor R_{AA} defined as the ratio of normalized single particle yields in p+p and A+A collisions.

Figure 1.12 shows R_{AA} of various particle species as a function of the transverse momentum p_T measured by the PHENIX experiment. A strong suppression of light hadron production ($R_{AA} < 1$) at high p_T was observed for the first time at RHIC, while R_{AA} of direct photons (photons produced at all stages through the system evolution except those from hadronic decays) is consistent with unity as expected since the photons do not interact via the strong force.

1.5.2 Collective Flow

The idea that the fireball may be imagined as a thermalized gas of highly interacting particles with quark and gluons as degrees of freedom is confirmed by the success of the Statistical Hadronization Model (SHM) in characterising particle generation in both pp and A–A collisions. The fireball expands as it cools down after being formed in the collision. The system is pushed from within by the fireball’s internal pressure, which causes it to expand. The particles velocities in the fireball are not independent, but rather are influenced by pressure gradients. As particles begin to exhibit correlated velocities that rely on their position, the fireball begins to behave like a fluid. As a direct consequence, particle velocities won’t be distributed in the end state according to random thermal motion, but rather will maintain a collective motion imprinted by their common source. This means that the momentum distribution of particles is affected by the presence of a thermalized medium, just as it would be in an increasing volume of gas. In contrast to “non-collective flow,” in which particles exhibit uncorrelated velocities, this characteristic is generally referred to as “collective flow.” The existence of a strongly interacting QGP that exhibits hydrodynamic development can be used to explain why collective flow occurs in heavy-ion collisions.

1.5.2.1 Radial Flow

The abundances of the different particle species are defined after the chemical freeze-out, but the particles continue to interact with one another through elastic scatterings. All interactions end and the primary particle momentum spectra become fixed only after the thermal freeze-out. The transverse-momentum spectrum at low p_T can be roughly characterised by a Maxwell-Boltzmann distribution if the system is in thermal equilibrium, where the slope parameter T_{fo} indicates the thermal freeze-out

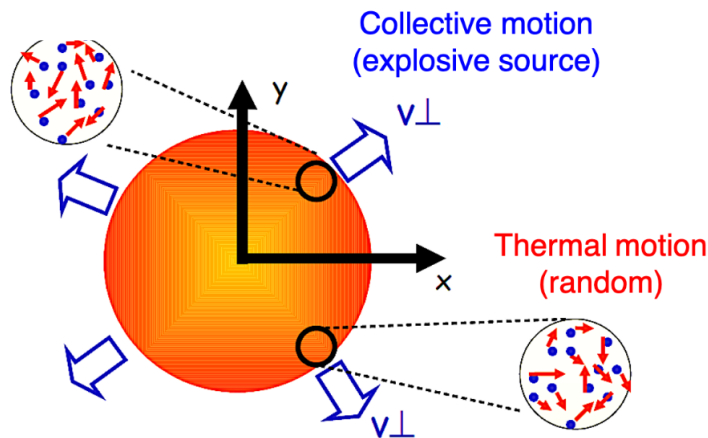


Figure 1.13: Schematic representation of radial flow in heavy-ion collisions (Figure adapted from [49]).

temperature:

$$\frac{1}{p_T} \frac{dN}{dp_T} \propto e^{-\sqrt{m^2 + p_T^2}/T_{fo}} \quad (1.12)$$

However, in A–A collisions a different slope for particles with different masses was observed [45]. This is the effect of a collective motion in the transverse plane, called radial flow, which is caused by the internal pressure generated by the QGP (shown in Figure 1.13) and is superimposed to the thermal random motion. Thus, the slope parameter (T_{slope}) is the sum of the thermal freeze-out temperature and the radial flow terms:

$$T_{slope} = T_{fo} + \frac{m}{2} \langle \beta_T \rangle^2 \quad (1.13)$$

where $\langle \beta_T \rangle$ is the collective transverse average velocity of the medium. The fact that the system is expanding, and as a result, the velocity is not uniform along the radial direction, adds further complexity.

At low p_T the formulation given in Equation 1.13 allows one to distinguish between the influence brought on by radial flow and that generated by thermal motion (T_{fo}). The radial flow appears to be connected to the particle mass, producing various slopes depending on the type of particle (breaking the m_T -scaling). In general, it is

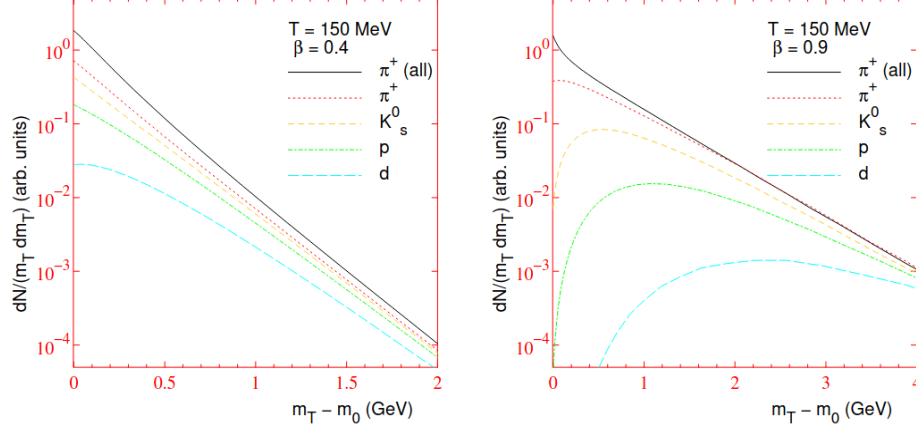


Figure 1.14: Calculation of the m_T spectra for different particle species with non zero flow [46]. The effect on the spectral shape is more evident for heavier particles as the flow velocity increases. The contribution of resonance decays to the spectrum of pions is reported as $\pi^+(\text{all})$, this inclusions have important effects on the spectral shape, especially for low m_T .

demonstrated that the spectra flatten at low momenta while experiencing a steeper decline at larger p_T . Naturally, these formulations have certain practical limits. For instance, pions, the lightest hadrons, quickly enter the relativistic case and are difficult to measure at very low momenta (especially in collider experiments). Additionally, resonance decays and Bose-Einstein statistics both have an impact on non-relativistic pions. The solid black line in Figure 1.14 shows the spectra with contribution of resonance decays included in the calculations.

Experimental measurements of the radial flow can be made using the momentum distribution or m_T of the identified particles. In Figure 1.15 [47, 48], identified π^- , K^- , and $\bar{p}$ spectra for central Au–Au collisions and pp collisions are displayed as an example. The presence of the radial flow in the case of heavy ions breaks the m_T -scaling, which is very obvious in the pp scenario where all three spectra have the same slope (m_T scaling at work). As anticipated, radial flow flattens the spectra in the low p_T area in a mass dependent manner, but as momenta increase, all spectra begin to adopt the same slope.

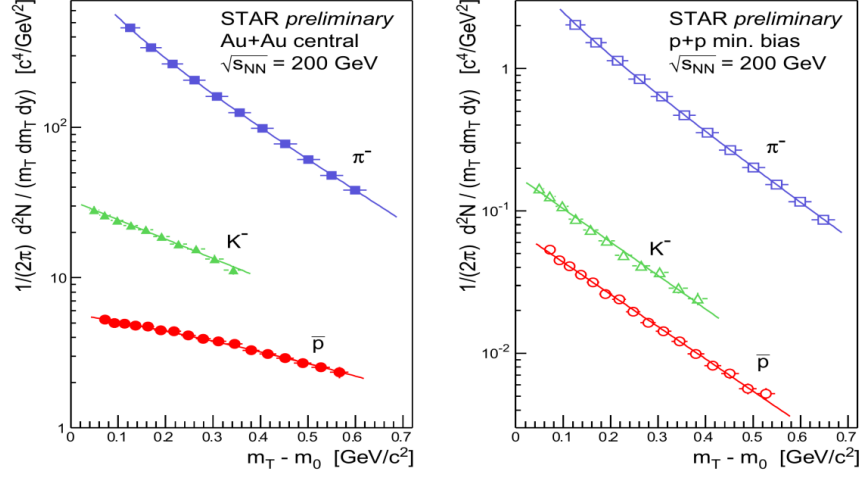


Figure 1.15: m_T spectra of π^- , K^- and p^- as measured at $\sqrt{s_{NN}} = 200$ GeV in central Au–Au collisions (left) and pp collisions (right) by the STAR Collaboration [47, 48].

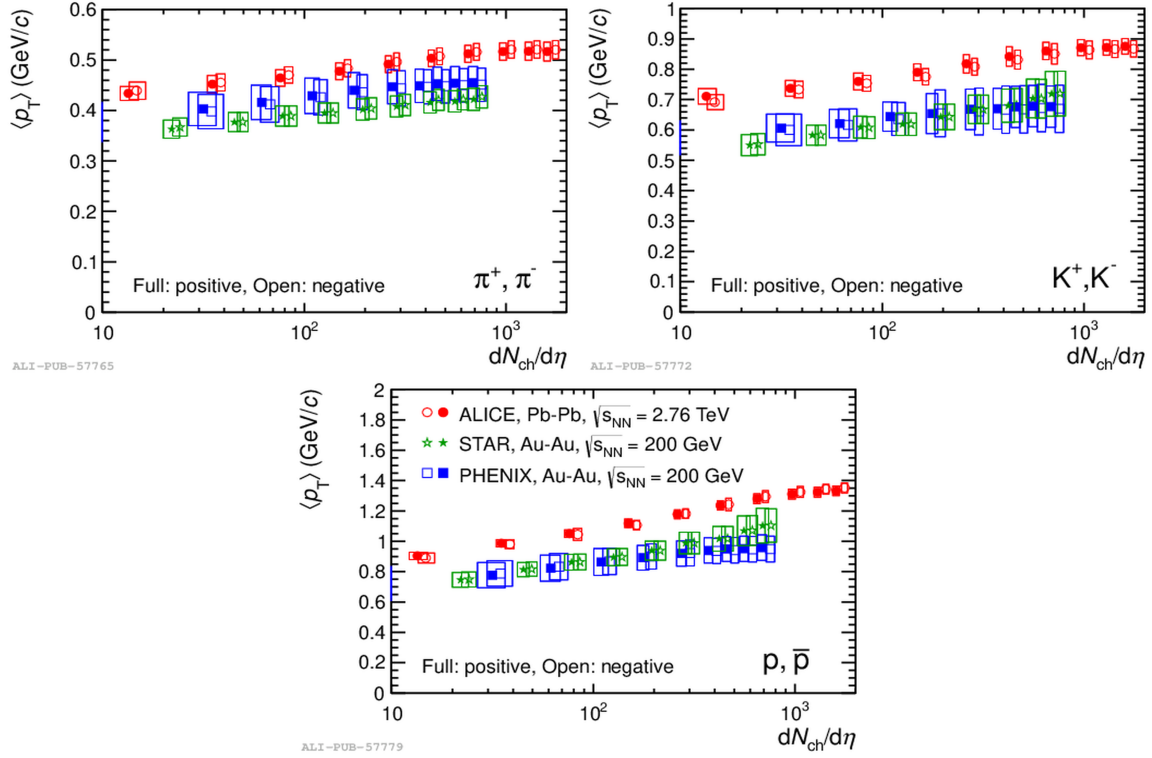


Figure 1.16: $\langle p_T \rangle$ for $\pi^\pm$, $K^\pm$ and $p\bar{p}$ as measured in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV and Au–Au collisions at $\sqrt{s_{NN}} = 200$ GeV [48]. A clear increase in the $\langle p_T \rangle$ is visible in central collisions.

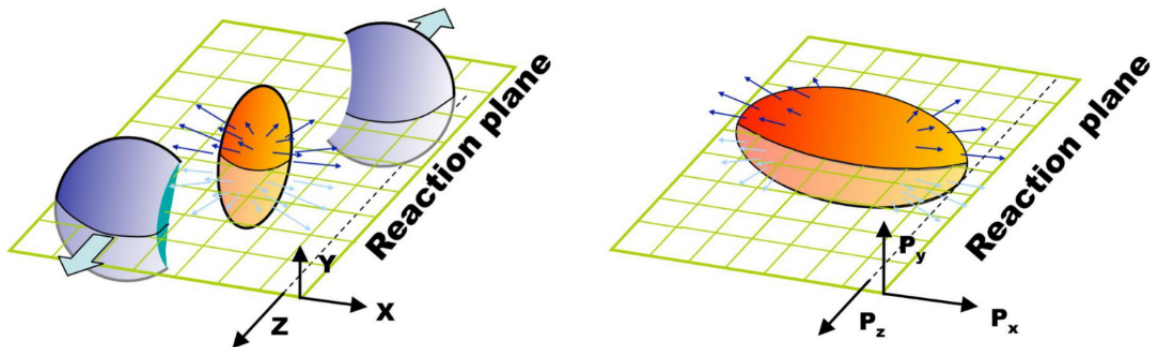


Figure 1.17: Non-central heavy-ion collision of two nuclei in position space (left) as well as in momentum space (right). On the left, the spatial asymmetry of the participant region (colored in orange) is shown. Due to thermal pressure against the vacuum, pressure gradients within the participant region build up. The spatial anisotropy causes the pressure gradient along the x-axis to be larger than along the y-axis. The resulting momentum anisotropy (referred to as collective flow) that builds up in the reaction plane is shown in the right picture. (Figure adapted from [50].)

The change in the spectral shape has a mass-dependent impact on the mean of the distribution as well, which causes the $\langle p_T \rangle$ to rise with particle mass. As a function of charged particle density at mid-rapidity for different collision energies, this is observed from the results at the RHIC and LHC energies reported in Figure 1.16 [48]. When flow is stronger, or when energy densities are higher, it is anticipated that the increase in the $\langle p_T \rangle$ will be more noticeable.

Since the energy density of the fireball and radial flow are both higher in central collisions, the increasing trend of $\langle p_T \rangle$ with $dN_{ch}/d\eta$ is expected there. For the same reason, the $\langle p_T \rangle$ is much greater when $\sqrt{s_{NN}}$ is increased from 200 GeV to 2.76 TeV. When comparing data in the most central collisions for different particle species, the mass dependency becomes apparent. The $\langle p_T \rangle$ is found to be constant between particle and anti-particle while assuming a hierarchical mass ordering.

1.5.2.2 Anisotropic flow

The overlapping region of two Lorentz contracted nuclei in non-central collisions is not circular but shaped like an almond as shown in Figure 1.17. The relationship between the mean free path λ of the particles and the size of the system R affects

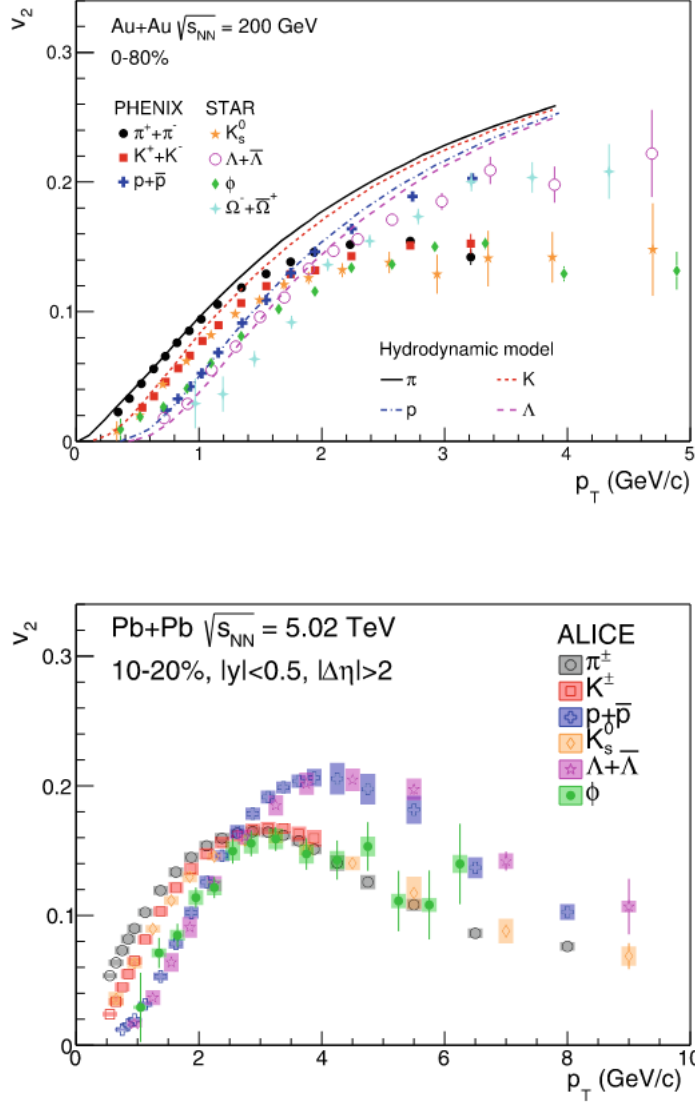


Figure 1.18: (Upper panel) v_2 as a function of the transverse momentum for $\pi^\pm$, $K^\pm$, K_S^0 , $p(\bar{p})$, ϕ , Λ ($\bar{\Lambda}$), Ω^- ($\bar{\Omega}^+$) in Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV from the STAR and PHENIX experiments [54–56]. Solid and dashed curves show the prediction of the hydrodynamic model [57]. (Lower panel) v_2 as a function of the transverse momentum for $\pi^\pm$, $K^\pm$, K_S^0 , $p(\bar{p})$, ϕ , Λ ($\bar{\Lambda}$) in semi-central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV from the ALICE experiment [58].

the emission pattern of the particles. The particle emission is isotropic in the transverse direction when λ is greater than R . However, when $\lambda \ll R$, a hydrodynamic explanation applies and the particle emission becomes anisotropic.

Hydrodynamic flow is derived by pressure gradients. The reaction plane is defined by the impact parameter $\vec{b}$ and the beam axis. In the overlapped region of two nuclei, the pressure gradient is steeper in the plane of the reaction plane and because of that more particles are produced in plane than out of plane. Consequently, the azimuthal distribution exhibits a distinctive $\cos(2\phi)$ modulation (elliptic flow), which is thought to be significant for the investigation of the hydrodynamic properties [51, 52]. Thus, the initial spatial anisotropy of the almond shape is converted to the momentum anisotropy called the elliptic flow. An important feature of the elliptic flow is that it is sensitive to the early stage of the collisions.

Experimentally, the azimuthal distribution is evaluated in terms of a Fourier expansion [53],

$$E \frac{d^3N}{d^3p} = \frac{d^2N}{2\pi p_T dp_T dy} \left(1 + \sum_{n=1}^{\infty} 2v_n \cos(n\phi) \right), \quad (1.14)$$

where ϕ is the azimuthal angle of produced particles with respect to the reaction plane. The second-order coefficient v_2 quantifies the strength of the elliptic flow.

The top panel of Figure 1.18 shows elliptic flow v_2 for pions, kaons, protons, ϕ , Λ , and Ω in mid-central Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV [54–56]. In $p_T < 2$ GeV/ c , v_2 increases with p_T and a clear mass dependence is observed, which is well described by the hydrodynamic model as shown with solid and dashed curves [57]. On the other hand, in higher p_T , there is a clear departure from the solid curves and two loci for mesons and baryons become visible. At the LHC, very similar behavior is observed: the bottom panel of Figure 1.18 shows v_2 for pions, kaons, protons, ϕ and Λ in mid-central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [58]. As seen at RHIC, a mass dependence is seen in the low p_T region while grouping of mesons

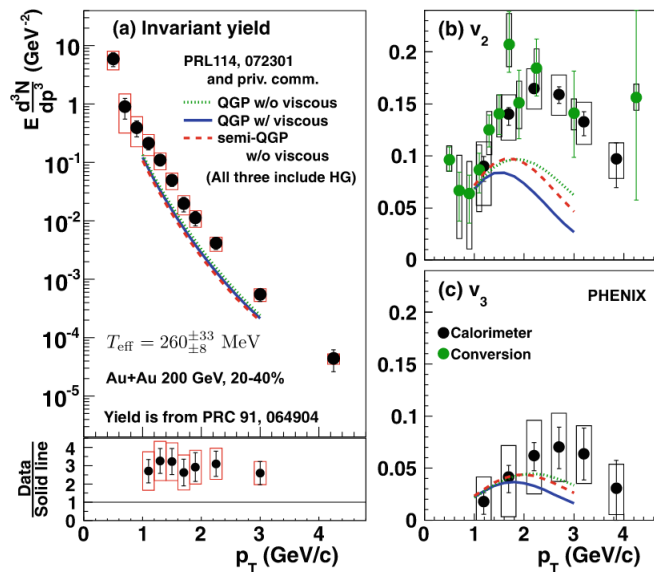


Figure 1.19: Direct photon yield (a), elliptic (b), and triangular (c) flow as a function of the transverse momentum [62] comparing to hydrodynamic model calculations with different assumptions. An effective temperature T_{eff} extracted from the inverse slope of the yield is shown in the figure. See [62] and references therein for details

and baryons is observed also in higher p_T , which suggests a different mechanism of particle production above $p_T \gtrsim 2$ GeV/ c compared to the lower p_T region, i.e., quark coalescence/recombination.

1.5.3 Direct photons and their puzzle

Due to their small cross section, photons are penetrating probes. They can be used to study the space-time evolution of the system since they are produced at all stages through the collision. Transverse momentum distributions of direct photons have been measured at both RHIC and the LHC [59–61], where the enhancement at low p_T (< 4 GeV/ c) in central A+A collisions relative to the scaled p+p data is described well by perturbative QCD. The excess indicates the photon production due to thermal radiation from the QGP. Figure 1.19a shows the direct photon yields after the subtraction of scaled p+p yield. An effective temperature extracted from the excess is found to be $T_{eff} = 260^{+33}_{-8}$ MeV for 20–40% Au+Au collisions at RHIC top energy and $T_{eff} = 297^{+12}_{-41}$ MeV for 0–20% Pb+Pb collisions at the LHC, both of

which are much hotter than the critical temperature, T_c .

Since photons are primarily emitted at early times (high temperature), one would anticipate that azimuthal anisotropy of direct photons would be small because the flow is developed later in time with the collective expansion of the system. However, experimental data at RHIC and the LHC show a sizable v_2 (even v_3) of direct photons comparable to the hadron v_2 [62–64]. Figure 1.19 shows the direct photon yield (a) and v_2 (b) and v_3 (c) measured at RHIC, comparing to hydrodynamic calculations. So far none of models can satisfactorily explain both the yield and v_2 simultaneously, which is known as “photon puzzle” (see [65] for recent review). There are still ongoing efforts from the experimental side to reduce the uncertainty as well as ones from the theoretical side with new ideas.

1.5.4 Quarkonium suppression (Debye screening effect)

Heavy quarkonium, such as charmonium or bottomonium, is expected to dissolve in the QGP when the potential between the quarks is screened by an abundance of colour charges of quark and gluon in the plasma, or the Debye screening effect. This suppression of heavy quarkonium has been proposed as a signature of the QGP [66]. Since the first measurement of the J/Ψ yields in Pb+Pb collisions at $\sqrt{s_{NN}} = 17$ GeV at SPS [67], heavy quarkonia have been thoroughly studied at RHIC and also at the LHC. Quantitative understanding of the yields at higher energies is found to be complicated since there are at least two competing effects: suppression due to the screening effects and the enhancement due to the recombination process.

Quarkonia having radius larger than the Debye length, the range of the interaction, are dissolved in the plasma. Therefore, the weaker bound quarkonium (larger radius) is expected to dissolve more completely compared to the stronger bound quarkonium (smaller radius). Thus, when suppressions of various quarkonia are compared, sequential suppression of their yields is expected according to their radii/binding en-

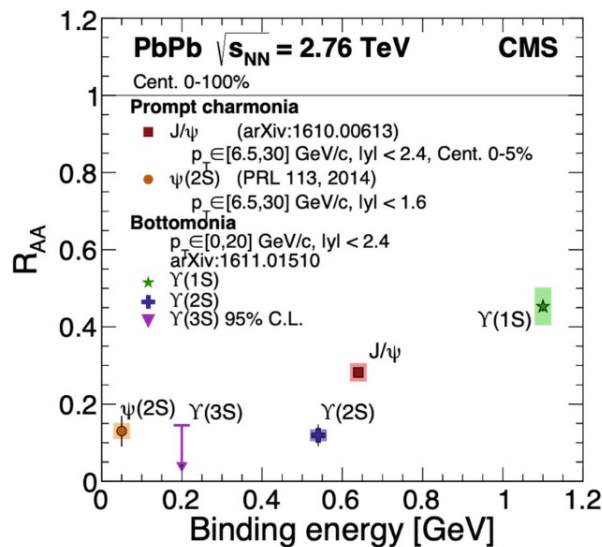


Figure 1.20: R_{AA} of quarkonium states as a function of the binding energy in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV measured by the CMS experiment, which is consistent with the sequential suppression pattern [68–70].

ergies. As shown in Figure 1.20, R_{AA} of charmonia, J/Ψ and $\Psi(2S)$, and bottomonia, $\Upsilon(1S)$ and $\Upsilon(2S)$, have been measured and the sequential suppression behavior has been clearly observed at the CMS experiment [68–70] which supports the assumption of QGP formation. For quantitative understanding, further theoretical and experimental studies are needed.

1.5.5 Electromagnetic field/chirality

In heavy-ion collisions, a short-lived strong magnetic field is created due to the moving electric charges carried by protons inside (fragmented) nuclei. The peak magnitude is anticipated to reach 10^{13-14} Tesla [71, 72], exceeding by a few orders than the magnetic field on the surface of neutron star called magnetars. The measurement of the strong field itself is of great interest. The lifetime of the field has large uncertainty and could be significantly extended depending on the electric conductivity of QGP due to Faraday’s law of induction [72–74]. To put it another way, one can examine the conductivity of QGP by studying the QGP response to the strong magnetic field generated by the charged spectator fragments. It is suggested that the effect of

the magnetic field appears in difference of directed flow (the first-order coefficient in Equation 1.14) between particles and antiparticles [75, 76], although the uncertainty is still too large to make any statement [77, 78].

Particularly for asymmetric collisions, such as Cu+Au, the electric field should be formed in the initial state in addition to the magnetic field. The effect may be seen in the charge difference of directed flow caused by the Coulomb force, which is dependent on the electric conductivity of QGP [79] and the temporal evolution of charge generation, such as quark and antiquark production [80]. Such a charge difference is demonstrated experimentally in hadron directed flow [81], showing the existence of the initial electric field.

1.5.6 Vorticity and polarization

Similar to the initial magnetic field, large orbital angular momentum is expected to be created in the initial state for non-central collisions. Due to spin-orbit coupling, a fraction of the orbital angular momentum would be transferred to the created matter, leading to global polarization of produced particles [82–84]. The term “global” refers to net spin alignment along the initial angular momentum direction which is perpendicular to the reaction plane and coincides with the direction of the initial magnetic field. At the STAR experiment, global polarization of Λ and $\bar{\Lambda}$ hyperons was measured and the extracted vorticity is found to be $\omega \sim 10^{22} s^{-1}$ [85, 86].

The matter created in the collisions is realized as the most vortical fluid ever observed. The polarization may also help to constrain the lifetime of the initial magnetic field since the polarization due to magnetic-spin coupling differs in the sign between particles and antiparticles. Theoretical models [87–91] can describe the energy dependence of the polarization quantitatively as shown in Figure 1.21 where calculations from viscous hydrodynamic model [87] and a multi-phase transport model [88] are compared. For both the models local thermal equilibrium is assumed and the

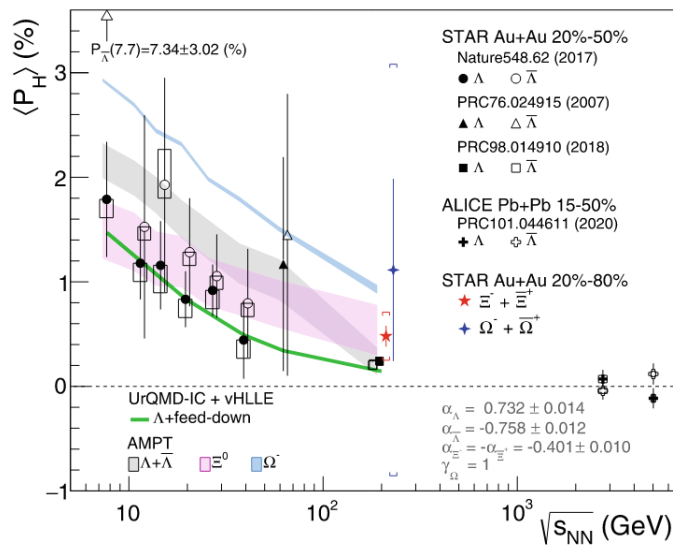


Figure 1.21: Global polarization as a function of the collision energy for Λ and $\bar{\Lambda}$ hyperons as well as for Ξ and Ω hyperons [85, 86, 92–94]. Solid line and shaded bands denote calculations from viscous hydrodynamic model (vHLL) with UrQMD initial condition [87] as well as a multi-phase transport (AMPT) model [88] respectively. Decay parameters α and γ of each particles [95] used in the measurements are shown in the figure. Note that the old results are corrected for recent update of the decay parameter.

polarization is calculated based on thermal vorticity at freeze-out [84, 96].

1.6 Modifications in the Cold Nuclear Matter

In order to give a definitive interpretation of the quantities observed in Pb–Pb collisions, the initial conditions of the Pb–Pb collisions have to be well understood. Those conditions can be explored in p–Pb collisions since the hot and dense medium is not expected to be created. Commonly, those conditions are called as cold nuclear matter effects. A brief explanation of these criteria will be provided in this section.

1.6.1 Cronin Effect

The results from the investigation of hadron production in pp collisions and p-nucleus (p-A) collisions in 1970 demonstrated that the invariant hadron cross sections in p-A relative to that in pp collisions not only scale with the quantity of binary collisions, but also exhibit a strong dependence on the atomic nuclei (A) for intermediate

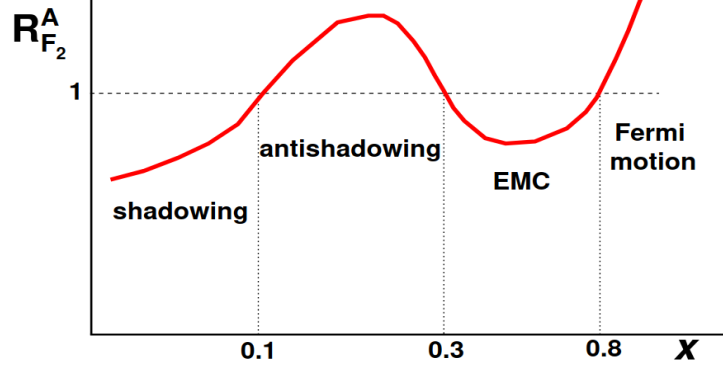


Figure 1.22: Ratio of the nuclear structure functions for a fixed Q^2 as a function of the momentum fraction x . Plot taken from [101]

transverse momenta [97, 98]. In this study, the invariant cross section dependence with A was parametrized as [97]:

$$\frac{Ed\sigma}{d^3p}(p_T, A) = \frac{Ed\sigma}{d^3p}(p_T, 1)A^\alpha \quad (1.15)$$

They discovered that $\alpha > 1$ for transverse momenta greater than 2 GeV/c. This indicates that an enhancement of the hadron production is expected in that region of p_T . This phenomenon is called Cronin effect and it is explained in terms of multiple scattering of the incoming partons when they propagate through the target nucleus [99].

1.6.2 Shadowing, Anti-shadowing and EMC Effect

It is well known that the nuclear structure functions of free nucleons ($F_2^{nucleon}$) differs from those of nucleons confined in a nucleus (F_2^A). The EMC experiment made this discovery when compared the nuclear structure function of nucleons bound in nuclei with those nucleons of deuterium. The deuterium was used in this study because it is the best approximation to the structure function of a free nucleon [100]. The behaviour of the nuclear distribution function can be observed in the $R_{F_2}^A$ which

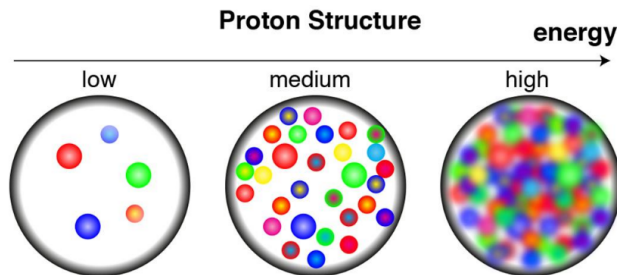


Figure 1.23: Illustration of the parton density of protons at different energy resolutions.

is defined as[101]:

$$R_{F_2}^A(x, Q^2) = F_2^A(x, Q^2) / A F_2^{nucleon}(x, Q^2) \quad (1.16)$$

where A is atomic number. The effects are normally classified in different regions of x [100, 101] as it is shown in Figure 1.22:

- $x > 0.8$ is called the fermion motion.
- $0.1 - 0.2 < x < .8$ the behaviour in this region is called “EMC” effect.
- $x \approx 0.1 - 0.2$ is called anti-shadowing.
- $x < 0.05 - 0.1$ this phenomenon is called shadowing.

1.6.3 Color-Glass Condensate

For many years, Deep Inelastic Scattering (DIS) studies have been used to study the structure of protons. The fact that the number of quarks and gluons inside a nucleon increases quickly as the high energy resolution is increased is one significant finding of the DIS experiments. Figure 1.23 provides a visual representation of this.

The enhancement of the parton density at high energy interactions is due to the Lorentz contraction of nucleons while traveling to velocities close to the speed of light [102]. The nucleons are composed by three valence quarks bound by gluons and the

valence quarks eventually fluctuate into new gluons and $q - \bar{q}$ pairs. In low energy interactions, only a small number of these extra gluons and quarks are visible because of their short life and because they do not affect the reaction [102]. In contrast, for very high energy collisions, the Lorentz contraction of nucleons not only modifies the geometry of the nucleon, but it also increases the lifetime of these fluctuations making them observable.

Another significant finding of the DIS experiments is that the partonic content of protons is dominated by gluons at tiny values of the Bjorken variable because the valence quarks are negligible at these scales and the sea quarks are suppressed by the weak coupling constant [102]. As a result, the gluon density rises at lower x values and higher Q^2 values. Because low energy gluons begin to recombine to generate more energetic ones, it is projected that the continual increase in gluon density at decreasing values of x will eventually reach saturation.

Strongly interacting systems with high parton densities $\alpha_s \ll 1$ are considered as non-perturbative in the QCD theory. This is due to the number of processes involving multiple partons which makes non-linear effects (which break down pQCD) more important. The scale in which the non-linear effects becomes important is known as saturation scale (Q_s). As a result, the process with a $p_T \ll Q_s$ can be affected by the gluon saturation regime. The Color-Glass Condensate (CGC) [102] is a theory that provides a description of the process taking place below the saturation scale. The saturation scale depends on the x and on the nucleon thickness as $Q_x \sim A^{1/3}$ [102]. Therefore, the gluon saturation effects are expected to be present in Pb–Pb and p–Pb collisions at the LHC energies and this can further suppress the hadron production at high p_T .

1.7 Thesis Motivation

The research work is carried out in the context of one of the experiments at LHC namely the ALICE (A Large Ion Collider Experiment) in p-Pb collision system at $\sqrt{s_{\text{NN}}} = 8.16$ TeV using Run 2 data. The main physics goal of the experiment is the creation and the investigation of the properties of the strongly-interacting matter in the conditions of high energy density and high temperatures, expected to characterize the medium formed in central heavy-ion collisions at these energies.

The motivation of the thesis is to explore the “breaking region” where high multiplicity pp collisions and low multiplicity nucleus-nucleus collisions meet. This is still a region of big interest since there is lack of statistics from both data sets (pp and PbPb) and therefore room for various interpretations. To tackle this problem results from p-Pb collisions is important. Thus, main aim of this thesis is to better understand the production mechanisms for strange particles and hence the unexpected observations like strangeness enhancement phenomenon in small collision systems.

The detailed work done in her physics analysis includes the transverse momentum (p_{T}) spectra measurements in different multiplicity classes, the ratio of the p_{T} integrated yield of strange hadrons to pions, baryon to meson ratio (Λ/K_{S}^0) as a function of p_{T} and multiplicity, nuclear modification factor (R_{pPb}) calculations and resonance to ground state yield ratio as a function of multiplicity.

Chapter 2

This chapter gives the theoretical description of strangeness production in QGP. The baryon-to-meson ratio (baryon anomaly) will also be explained by considering different production mechanisms.

Chapter 2

Strangeness Production in QGP

2.1 Probing QGP with Strangeness

In order to observe the variations in the production of strange particles in presence of the QGP, it is useful to define the strangeness content as the ratio (Wroblewski coefficient)

$$W_s \equiv \frac{2\langle s\bar{s} \rangle}{\langle u\bar{u} \rangle + \langle d\bar{d} \rangle} \quad (2.1)$$

where the quantities in brackets represent the number of quark-antiquark pairs [103].

The particle yields are investigated with the purpose of reconstructing the strange-to-light-quarks W_s ratio. The greater s quark mass in comparison to the u and d masses suppresses the generation of strange particles in ordinary nuclear matter [104]. The creation of strange baryons in low-energy nucleon-nucleon collisions is suppressed due to the absence of strange valence quarks in the interacting nucleons. As opposed to this, if deconfinement occurs in ultrarelativistic nucleus-nucleus collisions, the yields for strange-particle species should be higher than the corresponding yields at the same energy in nucleon-nucleon collisions: here the reactions leading to strangeness production are partonic and have lower thresholds and higher cross-sections, in particular if the strange mass is reduced by the partial restoration of the

chiral symmetry [105]. The threshold for the $s\bar{s}$ pairs should be lowered as a result, allowing their formation via gluon-gluon fusion in the QGP conditions [106]. In these conditions, the quarks should retain only their bare masses, reducing from ~ 500 MeV to ~ 150 MeV for the s quark. Moreover, in a system rich of baryons with respect to the antibaryons with high concentration of up and down quarks, the Pauli principle could block the production of $u\bar{u}$ and $d\bar{d}$ pairs: the energy of the lowest free level available for a light quark should grow enough to be comparable with the energy needed to create a $s\bar{s}$ pair.

The concepts of the strangeness enhancements can be summed up in more depth as follows: when color bonds are broken, the chemically equilibrated deconfined state has an unusually high abundance of strange quarks [107]. Perturbative QCD studies on the dynamical process of chemical equilibration has shown that only the gluon component in the QGP is able to produce strangeness rapidly, allowing formation of (nearly) chemically equilibrated dense phase of deconfined, hot, strangeness-rich, quark matter in relativistic nuclear collisions [108]. So strangeness enhancement is directly connected to gluon presence in QGP. The high density of strangeness formed in the reaction fireball favors formation of multi-strange hadrons which are produced rarely if only individual hadrons collide [109].

The appropriate observable is the “enhancement” E of strange particle Y [110]:

$$E_Y = \frac{\langle N_Y^{A-A} \rangle \langle N_{part}^{p-A} \rangle}{\langle N_{part}^{A-A} \rangle \langle N_Y^{p-A} \rangle} \quad (2.2)$$

where $\langle N_{part}^{p-A} \rangle$ and $\langle N_{part}^{A-A} \rangle$ are the numbers of participant in a p-A (A-A) collision and $\langle N_Y^{p-A} \rangle$ and $\langle N_Y^{A-A} \rangle$ are the number of strange particles produced in a p-A (A-A) collision. It is expected that $E_Y > 1$ and increasing with strangeness content of the

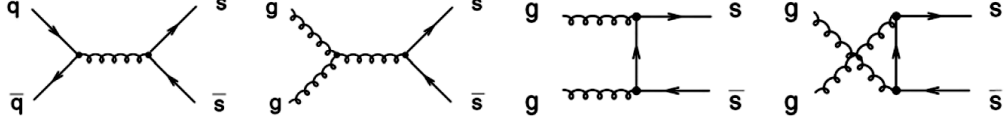


Figure 2.1: Leading order Feynman diagrams for strangeness production.

particle. The reason is that strangeness-producing processes in a QGP

$$qq \leftrightarrow ss \quad gg \leftrightarrow ss \quad (2.3)$$

should equilibrate faster than the corresponding processes in a hadron gas (HG)

$$\pi^+\pi^- \leftrightarrow KK \quad \pi N \leftrightarrow \Lambda K. \quad (2.4)$$

This can be seen relatively quickly by computing the momentum exchange $\langle Q \rangle$ needed for these processes ($\langle Q \rangle \sim 2m_s$ for Equation 2.3, $\langle Q \rangle \sim 2m_K$ for Equation 2.4). Hence, the strangeness abundance should reach chemical equilibrium much faster in a quark gluon plasma than in a hadron gas. Since the initial strangeness abundance in collisions is zero, the number of strange particles in a system of a certain lifetime with a phase transition should be parametrically higher than for a similar system where the transition did not occur. This is particularly true for multi-strange particles: in the hadron gas phase, these can only be produced sequentially [111, 112]

$$\pi\pi\pi N \leftrightarrow \pi\pi\Lambda K \leftrightarrow \pi\Xi KK \leftrightarrow \Omega KKK \quad (2.5)$$

We have mainly two lines of argumentation:

1. The production threshold for the production of strangeness via $s\bar{s}$ pairs is considerably lower than the one for hadrons. For this argument they compare the

energy needed to produce strange particles for a gas of quarks and gluons to those for a hadron gas. The associated $s\bar{s}$ pair production can proceed by the fusion of two gluons $g + g \rightarrow s + \bar{s}$ or two (massless) light quarks ($q = u, d$), $q + \bar{q} \rightarrow s + \bar{s}$, which involves only $Q_{QCD} = 2m_s \approx 200$ MeV. But in the processes connected to a hadron gas like $NN \leftrightarrow N\Lambda K$ the involved scale Q_{HG} is larger: $Q_{HG} = m_\Lambda + m_K - m_N \approx 670$ MeV. This should lead to larger production of strangeness in the QGP compared to a hadron gas.

2. The timescale for the equilibration of the production of strange particles in a QGP is significantly smaller than that for a hadron gas. This means that the number of particles produced in a QGP should be close to the equilibrium value and not be suppressed by any dynamical effects which would occur from hadronic interactions. In contrast to the light quark method, the generation via gluons has three possible lowest order Feynman diagrams. Figure 2.1 displays all lowest order diagrams. By calculating the rate of $s + \bar{s}$ production, Rafelski and Müller are able to demonstrate that the dominant production process is $g + g \rightarrow s + \bar{s}$. If QGP is formed, this would enable increased production of (multi-)strange hadrons in a heavy-ion collider. But in a completely hadronic scenario, this would not be the case.

Figure 2.2 shows the yield measured in AA collisions compared to the scaled pp (and pBe for the NA57 [113] results) versus the average number of participants [114, 115]. The figure also shows data from STAR [118] (an experiment at RHIC). The comparison of these different experimental results indicate that the observed enhancement decreases with the centre-of-mass energy of the collision. There is another observation from the results which shows the increase of the yield relative to the scaled reference, for particles containing more strangeness. This ordering of the enhancement E of the particles : $E(\Omega) > E(\Xi) > E(\Lambda)$ was predicted in [119]

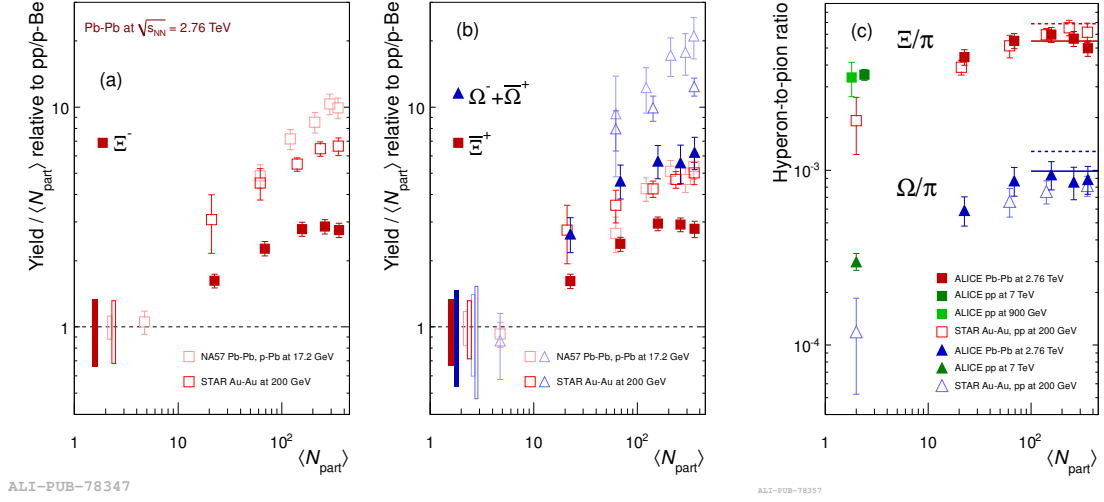


Figure 2.2: (a,b) Enhancements in the rapidity range $|y| < 0.5$ as a function of the mean number of participants N_{part} , showing LHC (ALICE, full symbols), RHIC and SPS (open symbols) data. Boxes on the dashed line at unity indicate statistical and systematic uncertainties on the pp or p-Be reference. Error bars on the data points represent the corresponding uncertainties for all the heavy-ion measurements and those for p-Pb at the SPS. (c) Hyperon-to-pion ratios as a function of N_{part} , for A-A and pp collisions at LHC and RHIC energies. The lines mark the thermal model predictions from [116] (full line) and [117] (dashed line).

According to today's understanding, the visible enhancement is not due to an enhancement in AA collisions but a suppression in pp [114]. Since this can be explained using a thermal model with a canonical formulation of strangeness conservation, this suppression is known as canonical suppression. The conservation of quantum numbers is carried out precisely on an event-by-event basis in the classical formulation of the thermal model.

While the grand canonical formulation, which is utilised in heavy-ion collisions, just averages the related chemical potential, such as the baryochemical potential or the chemical potential of strangeness, to apply the conservation of quantum numbers. The observed ordering related to the strangeness content [120] and the decline of the enhancement with increasing collision energy [121] are both very well described by the canonical formulation.

Strangeness itself continues to be an intriguing probe of the medium formed in

heavy-ion collisions. In QCD itself, the strange quark has a significant role. It can be created thermally inside the hot fireball because its mass is well separated from the light (u, d) and heavy (c, b) quarks, whereas b and c quarks are created in the initial hard collision. Strange quarks, however, also have certain influence on the thermodynamic characteristics of QCD. The mass m_s of the strange quark controls the properties of the chiral and the deconfinement transition [122].

2.2 Theoretical models

Due to the fundamentally diverse features of the underlying hadron-hadron and consequent parton-parton interactions at various incidence energy and kinematic intervals, there is currently no one single theoretical explanation for these interactions. The hard processes, or processes with significant four-momentum transfer, Q^2 , can be described using perturbative quantum chromodynamics (pQCD). However, because to the absence of the large Q^2 -scale, pQCD is formally incorrect for the description of the soft interactions. Therefore, phenomenological models are the most effective way to characterise low p_T particle generation.

For hadronic and nuclear collisions, a wide range of models have been created. They can be separated into microscopic (string, transport, cascade, etc.) and macro (statistical and hydrodynamical) models.

A general overview of these models will be provided in the sections that follow, with special emphasis on the models that will be utilised in comparison of experimental findings on strange baryons with theoretical forecasts.

2.2.1 Hydrodynamic models

After the pre-equilibrium phase, when thermodynamical equilibrium is quickly reached, the system expands and cools down (Figure 1.8). Using a hydrodynamical technique, the system's behaviour can be described. Landau first suggested this type of description in 1953 [123]. With the help of a measurement on heavy-ion collisions

at ultra-relativistic energies carried out at the RHIC, this initial hypothesis was confirmed. According to significant values of elliptic flow obtained at larger energies (in particular STAR, PHOBOS and PHENIX results) indicate that the matter produced in Au–Au collisions at RHIC does not behave as a weakly interacting gas (for which elliptic flow v_2 would be negligible), but instead as a strongly coupled liquid [124]. This conclusion is supported by the latest LHC results, which also qualitatively support hydrodynamic predictions of elliptic flow for Pb-Pb collisions at the LHC energy [125].

The initial conditions provide the input for the hydrodynamic evolution of the system $P(\epsilon, n_B)$ (EoS, with energy density ϵ , pressure P and baryon density n_B). Given the success of the hydrodynamical approach in the ideal limit ($\eta/s = 0$) in explaining the experimental data and the theoretical lower limit for the shear viscosity over entropy ratio ($\eta/s > 1/4\pi$) [126–128], the state produced during a relativistic heavy-ion collision is interpreted as a fluid that is almost “perfect”. However, there is some disagreement in the description of the data (particularly at large p_T), which encourages the use of dissipative hydrodynamics. The stringent assumption of local thermal equilibrium, which is required in the ideal limit, is also relaxed in such conditions to the assumption of near local thermal equilibrium, widening the scope of applicability of hydrodynamic description.

VISH2+1 [129] is an illustration of a pure viscous hydrodynamic model. On the other hand, the Krakow model [130] incorporates non-equilibrium corrections brought on by viscosity in the change from a hydrodynamic description to one involving the particles of the final state. The Cooper-Frye freeze-out picture is used in both models to describe the change from hydrodynamic fluid to free particles that enter the detector [131]. In this illustration, it is assumed that as the hydrodynamic expansion comes to an end, the momentum distribution of the final state particles is essentially the

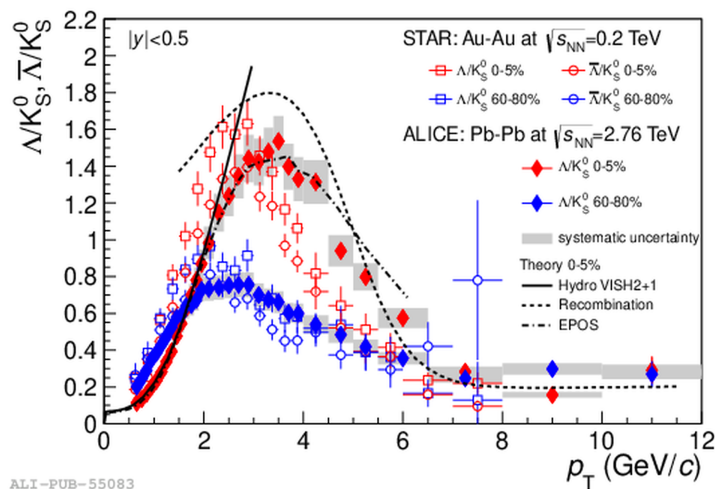


Figure 2.3: Λ/K_S^0 ratios as a function of p_T measured by the ALICE Collaboration in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV, compared with Λ/K_S^0 and $\bar{\Lambda}/K_S^0$ ratios measured in Au–Au collisions at $\sqrt{s_{NN}} = 200$ GeV by STAR Collaboration. The dot-dashed line show the corresponding ratios from EPOS model. [141]

same as the momentum distribution within the fluid, and that some fluid components instantly transform into free particles.

2.2.2 Microscopic models

Using a microscopic description of the dynamics of the constituents of the system based on transport theory is another technique to explain the transition from the hydrodynamic to the free particle phase or from the pre-equilibrium to the QGP phase. Such models are known as cascade models. Finding solutions to the equations of motion is the main objective of this transport model, which follows the individual trajectories and interactions of all particles (baryons and mesons in the case of hadronic transport models or quarks and gluons in the case of a parton cascade models). The interaction cross sections and the decay widths of the unstable particles are the two most important components in this description. The crucial thing to remember is that the particles are moving through a hot, dense, and possibly even not equilibrated medium, which could dramatically alter the prior values relative to those that are understood in the vacuum. This is an extremely challenging goal to

achieve, necessitating substantial model approximations.

Hadron String Dynamics (HSD) [135], Ultra-relativistic Quantum Molecular Dynamics (UrQMD) [132, 133], and Relativistic Quantum Molecular Dynamics (RQMD)[134] are a few examples of hadronic transport models used in heavy-ion physics to analyse particle production.

Similar to VISH2+1, which also includes a hadronic cascade (UrQMD) after the partonic hydrodynamic phase, HKM [136, 137] is an illustration of an ideal hydrodynamic model. On the other hand, the EPOS model (Energy conserving quantum mechanical multiple scattering approach, based on Partons, Off-shell remnants, and Splitting of parton ladders) [138–140], follows a phenomenological approach that is based on an initial parton model. It aims to be a comprehensive model and event generator, describing all p_T domains with the same dynamical picture. It specifically combines hydrodynamics and predicts how high p_T hadrons interact with the expanding fluid. Figure 2.3 illustrates how well this model has explained the baryon-to-meson increase seen by the STAR and ALICE Collaborations in the intermediate p_T region ($2 < p_T < 6$ GeV/ c).

2.2.3 Statistical models

The statistical hadronic models (SHM) use the statistical formalism to describe the system after hadronization. The goal is to identify the ensemble of observables that best describes the chemical composition in the hadronic phase using the fewest possible thermodynamic variables. The partition function is the fundamental component of these models. Using the Grand-Canonical (GC) formalism, the particle production from the large volume system produced by the A-A collisions can be explained. The mean hadron multiplicities $\langle N_i \rangle$ [142] can be defined in this case as:

$$\langle N_i \rangle = (2J_i + 1) \frac{V}{(2\pi)^3} \int d^3p \frac{1}{\gamma_S^{-S_i} \exp[(E_i - \mu_i)/T_{\text{ch}}] \pm 1} \quad (2.6)$$

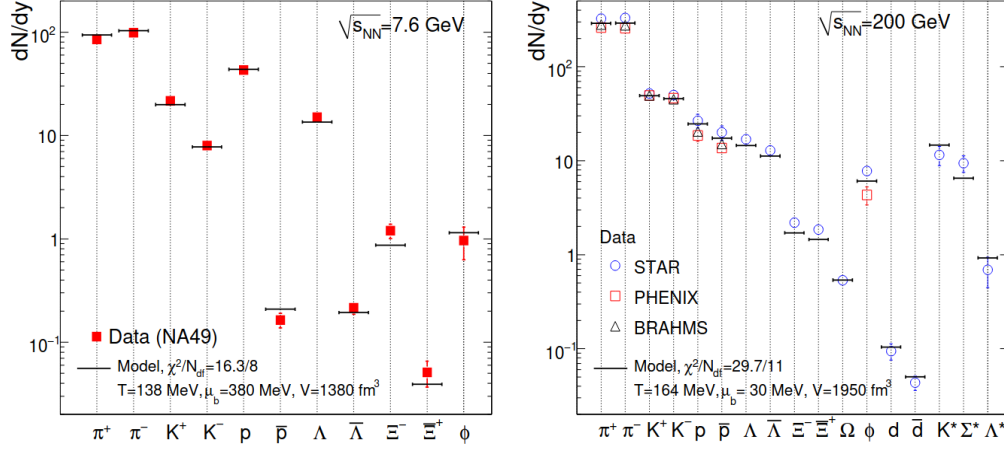


Figure 2.4: Experimental hadron yields and model calculations for the parameters of the best fit at the centre-of-mass energies of 7.6 GeV by the NA49 Collaboration (left panel) and 200 GeV by the STAR Collaboration (right panel) [146].

where the involved parameters are: the chemical freeze-out temperature T_{ch} , the chemical potentials $\mu_i = \mu_B B_i + \mu_S S_i + \mu_{I_3} I_{3i}$ (for a particle i with baryon number B_i , strangeness S_i and third component of the isospin vector I_{3i}), the volume V , the strange quarks fugacity γ_S , the spin J_i and the number of strange quarks S_i of a given particle type i .

The constraints of global strangeness and charge conservation govern the chemical potentials μ_S and μ_{I_3} , whereas the baryonic potential μ_B is a free parameter. The $+$ sign in the denominator is acceptable for fermions while the $-$ sign is used for bosons to account for different statistical behaviour. This model is commonly used in this context in a variety of implementations. To account for a possible non-equilibrium of the non-strange particles, some models incorporate an additional factor γ_q .

Three basic options can be distinguished [143] depending on the assumption made on the equilibration of the system:

2.2.3.1 Chemical Equilibrium ($\gamma_S = \gamma_q = 1$)

It is assumed that both strange and non-strange particles are in a full equilibrium. In other implementations, as it is in [144], the particle ratios are fitted, the

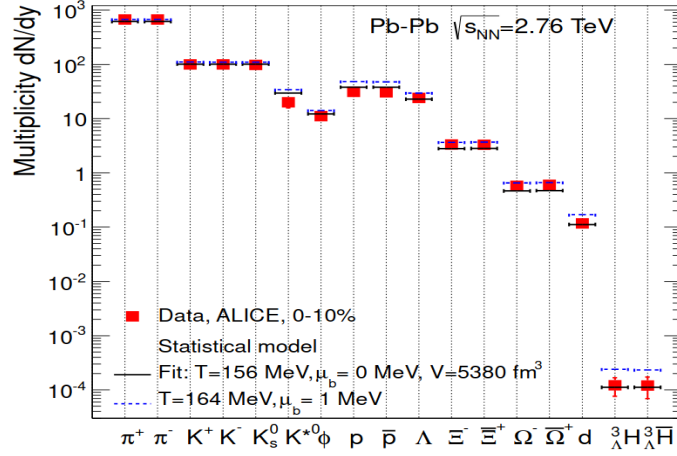


Figure 2.5: Integrated yields at mid-rapidity dN/dy in central (0–10%) Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV with results from a thermal fit. [147]

volume parameter drops out and the only free parameters are T_{ch} and μ_B . Figure 2.4 illustrates how well these models match observations made at SPS and RHIC. Similar fits have been made in Figure 2.5 by ALICE in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV using preliminary data. In the fits of Figure 2.5, μ_B is set to 1 MeV because the antibaryon over baryon ratios [145] indicate a vanishing baryon-chemical potential at the LHC.

There are some discrepancies when comparing the model predictions and the measured particle yields which are present and the system’s temperature is projected to be lower ($T = 156$ MeV) than that determined from STAR’s data ($T = 164$ MeV) [146]. Proton and antiproton yields, in particular, are underestimated by the model (by deviations of 2.7σ and 2.9σ respectively, with negligible experimental errors), whereas cascade yields, on the other hand, lie above the model (by roughly 2σ).

The results of a statistical model calculation with $T = 164$ MeV and $\mu_B = 1$ MeV, which increased the antiproton disagreement to nearly 50%, are also shown in the same figure. It should be noted that leaving protons out of the fit, the temperature would increase by 2 MeV to $T = 158$ MeV with an otherwise perfect fit with a reduced χ^2 less than one. As a result, one could speak of a proton anomaly, which is another

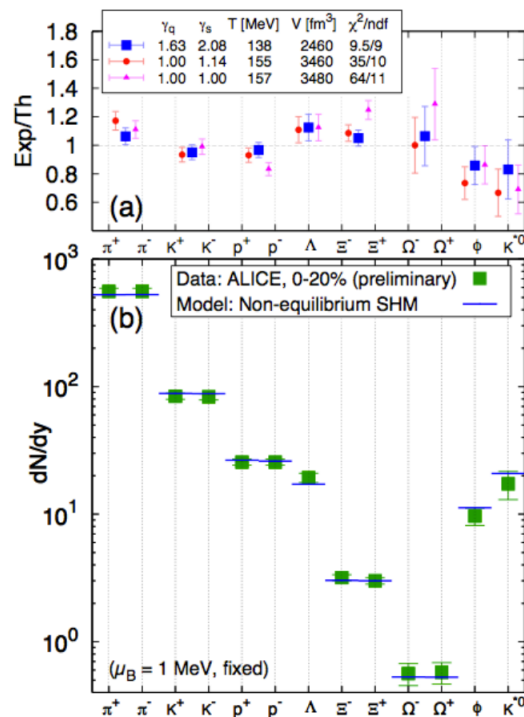


Figure 2.6: The lower panel shows the non-equilibrium SHM fit to experimental data measured by the ALICE experiment in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV for 0–20% centrality. The upper panel shows the ratio of model values to experimental data for the three SHM variants and presents the key parameter values for: chemical non-equilibrium (solid squares), chemical semi-equilibrium (solid circles) and chemical equilibrium (solid triangles). [154]

way of expressing that the proton-to-pion ratio is around 50% overpredicted [148, 149].

2.2.3.2 Chemical Semi-Equilibrium ($\gamma_q = 1, \gamma_s \neq 1$)

To allow for the possibility that strangeness may not be entirely equilibrated, the strange quark fugacity is employed ($\gamma_s < 1$). Some implementations explicitly calculate the multiplicities in accordance with Equation 2.6, as it is done in [150].

2.2.3.3 Chemical Non-equilibrium ($\gamma_q \neq 1, \gamma_s \neq 1$)

Additional implementations allow the non-strange quarks to also be out of equilibrium, which introduces a new factor γ_q [151, 152]. The SHARE code [153], for instance, presupposes an instantaneous hadronization process and a QGP in thermal equilibrium. Although additional free parameters are required in this situation, predictions from this model, as shown in Figure 2.6, are in good agreement with the

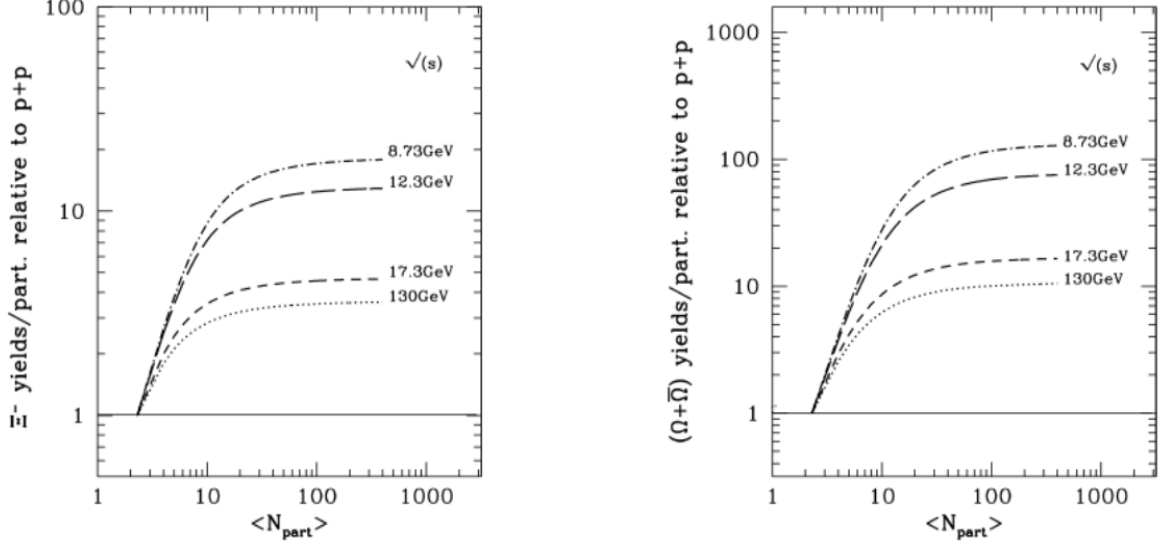


Figure 2.7: Centrality dependence of the relative enhancement normalized to $\langle N_{\text{part}} \rangle$, for Ξ^- (left) and $\Omega^- + \bar{\Omega}^+$ (right), at different collision energies, obtained in the statistical approach with canonical formalism [156].

results at LHC energy.

The conservation law of a quantum number, such as strangeness, can be implemented on average in a large system with a large number of produced particles by applying the relevant chemical potential within the Grand Canonical (GC) formulation. Conservation laws must be applied locally on an event-by-event basis in a small system, like a pp collision, with small particle multiplicity, necessitating a Canonical formulation (C). The phase space accessible for particle generation is known to be significantly constrained by the C conservation of quantum numbers. This is the canonical suppression (CS) mechanism [155, 156].

This implies that $E(h)$ increases in Equation 2.2 if the denominator is decreased by CS. For various colliding energies, the predictions for the enhancement as a function of centrality are illustrated in Figure 2.7. These predictions exhibit the same energy relationship as the observations in Figure 2.2 and so suggest that the canonical suppression will decrease as collision energy increases.

The argument of the canonical suppression does not adequately address every aspects of the problem. Figure 2.2 shows that this model does not replicate the centrality dependency of the strangeness enhancement experimentally observed at the given energy (Figure 2.2). The model specifically forecasts a saturation of the enhancement for the most central collisions.

The core-corona model, which postulates that the system is made up of two parts—a core made up of a hadron gas in chemical equilibrium and a corona where strange particle production occurs through nucleon-nucleon collisions—offers a potential solution to the participants dependence of the enhancement [157]. The weight of the core relative to the corona heading into most central collisions is increased to account for the centrality dependency of the strangeness enhancement. Measurements of the strangeness enhancements by STAR with Au–Au and Cu–Cu [158] at $\sqrt{s_{NN}} = 200$ GeV have been found to be in agreement with the model’s prediction.

The question of whether the phase space suppression in smaller systems (like pp to p-A) is the only reason for the strangeness enhancement in larger systems or whether it is actually caused by increased strangeness production from a deconfined partonic medium is still being debated.

2.3 Production Mechanisms

Hard scattering via parton fragmentation into hadrons is one of the recognised particle generation mechanisms in elementary collisions. In this illustration, the single parton spectrum is convoluted by the likelihood that a parton would hadronize into a hadron, which carries a fraction of the momentum of the parent parton. Using pQCD methods, the cross sections of these interactions can be calculated.

In this conventional view of hadron creation at high transverse momentum, experimental evidence such as the baryon-to-meson enhancement (Figure 2.3 for Λ/K_S^0 , also apparent for non-strange particles) is not adequately explained. While parton

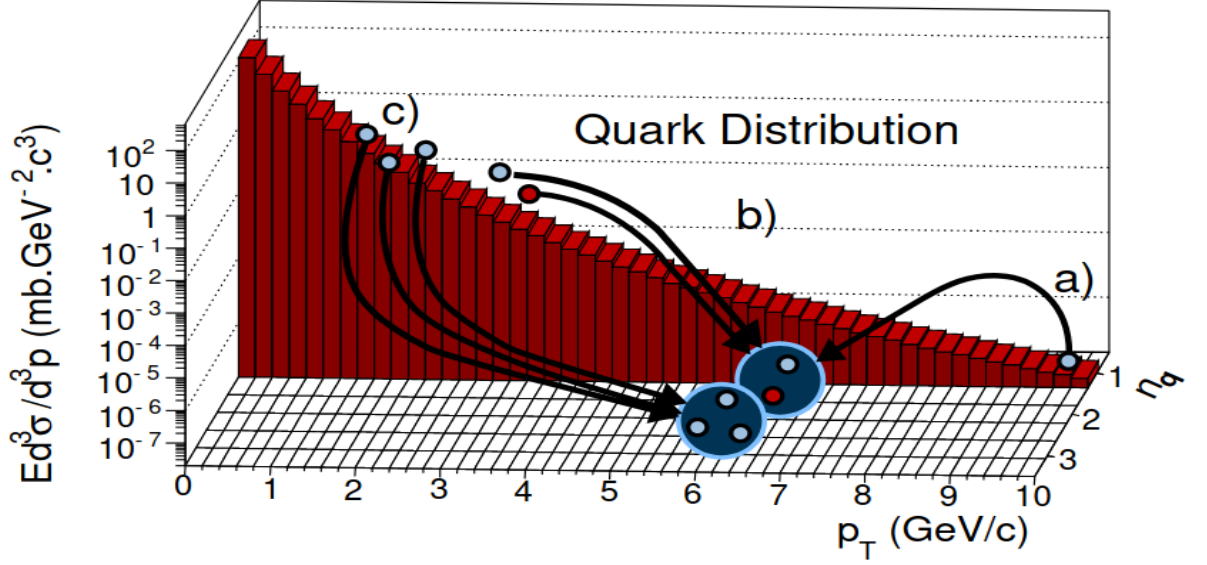


Figure 2.8: Illustration of the possible competition between fragmentation and coalescence at mid p_T . The depicted processes corresponds to: a) high p_T quark fragmentation into a meson, b) and c) creation respectively of a meson and baryon by recombination. [160]

energy loss during parton propagation through the hot, dense matter produced by the nuclear collision is responsible for the observed suppression of the meson yield, the lack of a similar effect in the proton spectrum is perplexing.

This “baryon-meson dilemma” may have a potential answer, according to one theory put forth by [159]: in environments with large parton densities, hadron creation may take place through parton recombination rather than fragmentation at momenta of a few GeV/c .

The fundamental presumption of recombination model is that the partons that form the particles in the intermediate p_T zone are a mixture of partons coming from various transverse momentum regions. On one hand, a high p_T parton breaks up into lower p_T hadrons, as shown in Figure 2.8. While on the other hand, partons produced softly coalesce into a hadron with a p_T that is sum of the p_T of the constituent partons. This is possible for both mesons and baryons, although for the latter, lower p_T partons are required for a fixed p_T of the generated particle. However, the thermal description

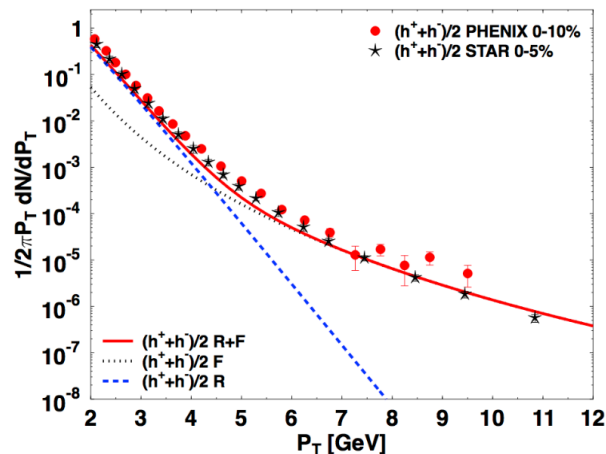


Figure 2.9: Transverse momentum spectrum of charged hadrons $(h^- + h^+)/2$ at mid-rapidity for central Au–Au collisions produced at RHIC, measured by PHENIX (in red, 0–10% central events) and by STAR (in black, 0–5% central events). Prediction from fragmentation model (dotted black line) and recombination model (dashed blue line) and the sum of both (solid red line) are showed. [161]

of the system at low p_T informs us that partons are (exponentially) more likely to exist at low p_T based on the exponential distribution and, as a result, to make a baryon via coalescence is more probable rather than a meson for a given p_T .

This mechanism is not present in the absence of a QGP, in which case fragmentation will disfavor baryons over mesons, thus suggesting an increased baryon over meson ratio in the QGP scenario if compared to HG system. The puzzle can therefore be qualitatively explained by a straightforward analysis of the shape of the parton spectrum.

The two processes switch their roles for the dominant production mechanism in the region of intermediate transverse momentum and can be observed in Figure 2.9

2.3.1 Nuclear modification factors

Comparing the particle spectra acquired in A–A collisions to those obtained from data where QGP is not anticipated, such as pp or p–Pb collisions, is one technique to learn about particle production mechanisms in nuclear matter. One is particularly

interested in comprehending how the hot and dense medium affects particle formation. In this situation, it is expected that there will be considerable energy loss by gluon radiation as the partons move through the medium, especially for high p_T partons [162]. Usually, this is done by measuring the nuclear modification factors, R_{AA} , defined as follows:

$$R_{AA} = \frac{1}{\langle N_{coll} \rangle} \frac{(d^2 N^{AA}/dp_T dy)}{(d^2 N^{pp}/dp_T dy)} \quad (2.7)$$

equation

where $d^2 N^{AA}/dp_T dy$ is the invariant p_T spectrum of nucleus-nucleus (AA), $d^2 N^{pp}/dp_T dy$ is the invariant p_T spectrum in pp collisions and $\langle N_{coll} \rangle$ is the average number of binary nucleon-nucleon collisions at a given centrality and is estimated by the Glauber model of the nuclear collision geometry [163].

Hard processes are predicted to dominate particle production at high p_T , which means that the number of binary nucleon-nucleon collisions N_{coll} should scale with the number of A–A collisions (binary collision scaling). The R_{AA} should be unity for high p_T ($p_T \geq 2\text{--}3$ GeV/ c for hadrons) if there are no impacts of the generated medium.

In Figure 2.10, where the measurements in A–A collisions at RHIC and ALICE have been shown, it can be observed that the production of charged hadrons in the high p_T region is suppressed compared to the expectation from an independent superposition of nucleon-nucleon collisions. The Figure 2.10a demonstrates how the suppression increases with collision energy [164].

Additionally, similar observables can be measured in p–A collisions, which are regarded as a reference for the investigation of the initial and final-state effects. Figure 2.10b shows the comparison of nuclear modification factor obtained by ALICE collaboration in Pb–Pb and p–Pb collision systems. According to this measurement,

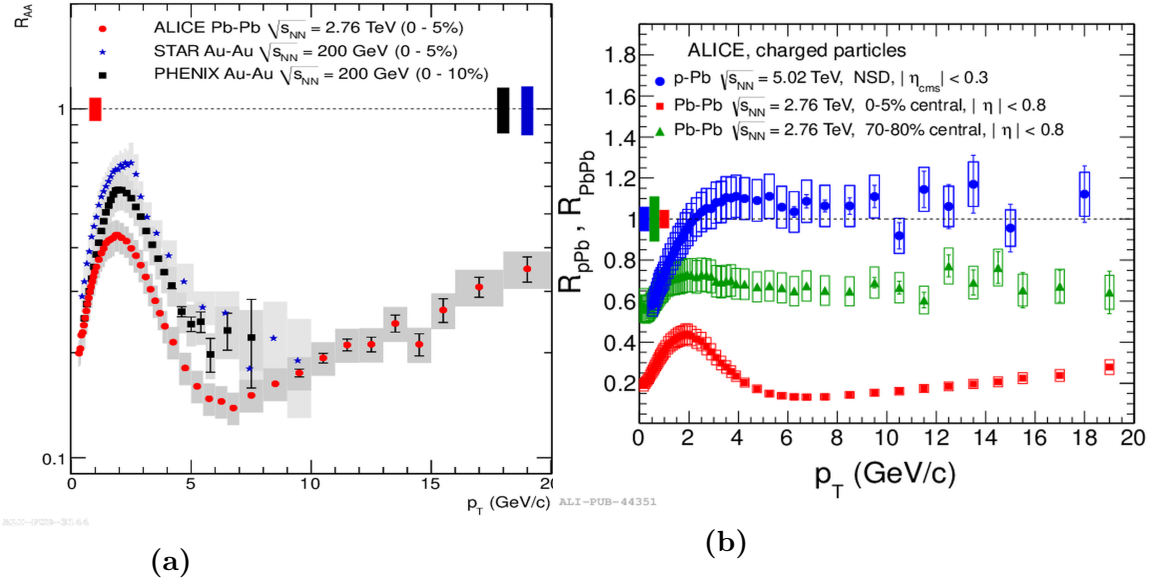


Figure 2.10: (a) Comparison of R_{AA} for unidentified particles in central Pb–Pb collisions at LHC to measurements at $\sqrt{s_{NN}} = 200$ GeV by the PHENIX and STAR experiments at RHIC. For both ALICE and PHENIX, the statistical and systematics errors are separated and represented by the lines and the boxes, respectively. For STAR the two errors are combined and represented by boxes. The vertical bars around R_{AA} equal to unity indicate the p_T independent scaling errors on R_{AA} . [164] (b) The nuclear modification factor of charged particles as a function of p_T in minimum bias (NSD) p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, compared to measurements in collisions at $\sqrt{s_{NN}} = 2.76$ TeV by the ALICE Collaboration. The errors are represented as in plot (a) for ALICE measurements. [165]

initial-state effects, which are present in both p–A and A–A collisions [165], are not responsible for the substantial suppression of hadron generation at high p_T observed in Pb–Pb collisions at the LHC. The distinct behaviour of mesons and protons at low and intermediate p_T ($p_T < 6$ GeV/c) could be interpreted as an evidence of the baryon-to-meson anomaly.

Chapter 3

Several striking signals of collective behaviour have been observed in small systems. These includes flow-like correlations, radial flow effects, enhancements in the production of multi-strange baryons etc. The possible interpretations of these unexpected results in small systems will be discussed in this chapter.

Chapter 3

LHC Results in Small Systems: Hints of QGP?

3.1 pp or p–Pb collisions: more than a reference for heavy ions

The number of particles produced in pp and p–A collisions is considerably smaller than the one reached in A–A as shown pictorially in Figure 3.1. For this reason and considering the mere size of the projectiles involved, these collisions are commonly labeled as “small systems” when compared to A–A. The main purpose of pp collisions in the field of heavy-ion physics is to provide baseline measurements for those performed in heavy-ion collisions. A modification of the observables in heavy-ion collisions with respect to the pp collisions suggests a presence of the QGP.

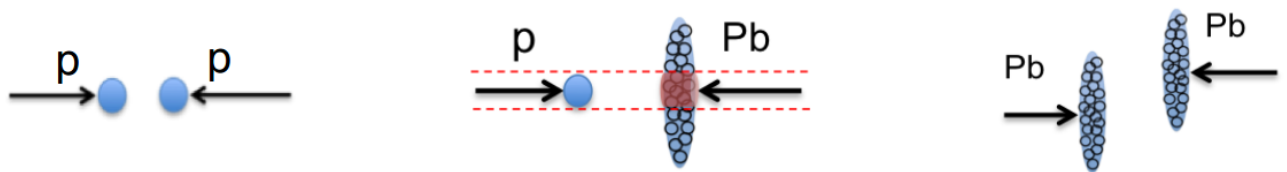


Figure 3.1: Schematic sketch of pp, p–Pb and Pb–Pb collisions

However, initial state effect arising from gluon interactions inside the nuclei during the early stages of a collision might survive during the whole system evolution and affect final observables, both in pp and heavy-ion collisions. For this purpose, asymmetric collisions of proton and a heavy nucleus were performed in order to identify and characterise the so-called Cold Nuclear Matter effects. During the recent years, this simple picture of small collision systems has been challenged. Measurements considered as an indication of the emergence of the QGP have been revealed in collisions of small systems, too. This has triggered lots of discussions about the nature of the observed collective phenomena in small collision systems and how it affects our understanding of large collision systems. A brief discussion of the new findings in small systems is presented in this chapter.

3.2 Experimental results in small systems

The study of particle production in high multiplicity pp and p-Pb collisions at the LHC has revealed unexpected new collective-like phenomena. The understanding on the origin of these effects is ongoing, for the majority of the heavy-ion community that are interpreted as the evidence of the QGP formation in small systems [166, 167]. Although, another fraction of the community is exploring new explanations in terms of initial state effects [168–170] i.e. without involving the formation of a small drop of QGP.

In recent years the number of studies going in this direction increased significantly revealing unexpected features. Below mentioned are some of the studies which have been performed:

3.2.1 Collective effects

Several effects, like near-side long-range correlations and mass-dependent hardening of p_T distributions, which in nuclear collisions are typically attributed to the formation of a strongly-interacting collectively expanding quark-gluon medium, have

been observed in high-multiplicity pp and p-Pb collisions at the LHC [171–174, 176–181, 194]. An enhanced production of strange particles as a function of the charged-particle multiplicity density ($dN_{\text{ch}}/d\eta$), originally considered to be another signature of QGP formation in nuclear collisions [182–184] has also been recently observed in pp collisions [185]. The study of small collision systems at high multiplicity is undoubtedly of considerable interest. Further studies are essential to assess whether the combination of these observations can be interpreted as signal of the progressive onset of a QGP medium, which starts developing already in small systems.

3.2.1.1 Two particle angular correlations:

Figure 3.2 shows the $\Delta\eta$ - $\Delta\phi$ correlations measured by the CMS experiment. The first significant results are shown in Figure 3.2 (b) for pp collisions at CMS experiment. Here the two-particle angular correlations ($\Delta\eta = \eta_1 - \eta_2$, $\Delta\phi = \phi_1 - \phi_2$) are shown as measured in events with high multiplicity [186]. The same type of correlation is reported as measured in p-Pb collisions [187] in Figure 3.2 (c). In this case the correlation obtained from low multiplicity events is subtracted to highlight the one specific to the high multiplicity events.

Correlations between pairs of charged particles plotted as a function of their relative difference in azimuthal angle $\Delta\phi$ and pseudorapidity $\Delta\eta$ show distinct features in heavy-ion collisions. From the two-particle correlation measured in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV is shown in Figure 3.2 (d) one can identify three main structures: the central peak at $(\Delta\eta; \Delta\phi) = (0; 0)$ originates from particles produced in the same jet; there is a long range (extending over a wide $\Delta\eta$ interval) correlation in particle emission in the “near side” ($\Delta\phi \sim 0$) due to the presence of collective flow; the “away side” ($\Delta\phi \sim \pi$) is characterized by a broad correlation (jet quenching). The ridge structure spanning over a large range in pseudorapidity at near-side ($\Delta\phi$), considered to be the effect of collective expansion of the matter during the QGP phase in heavy-

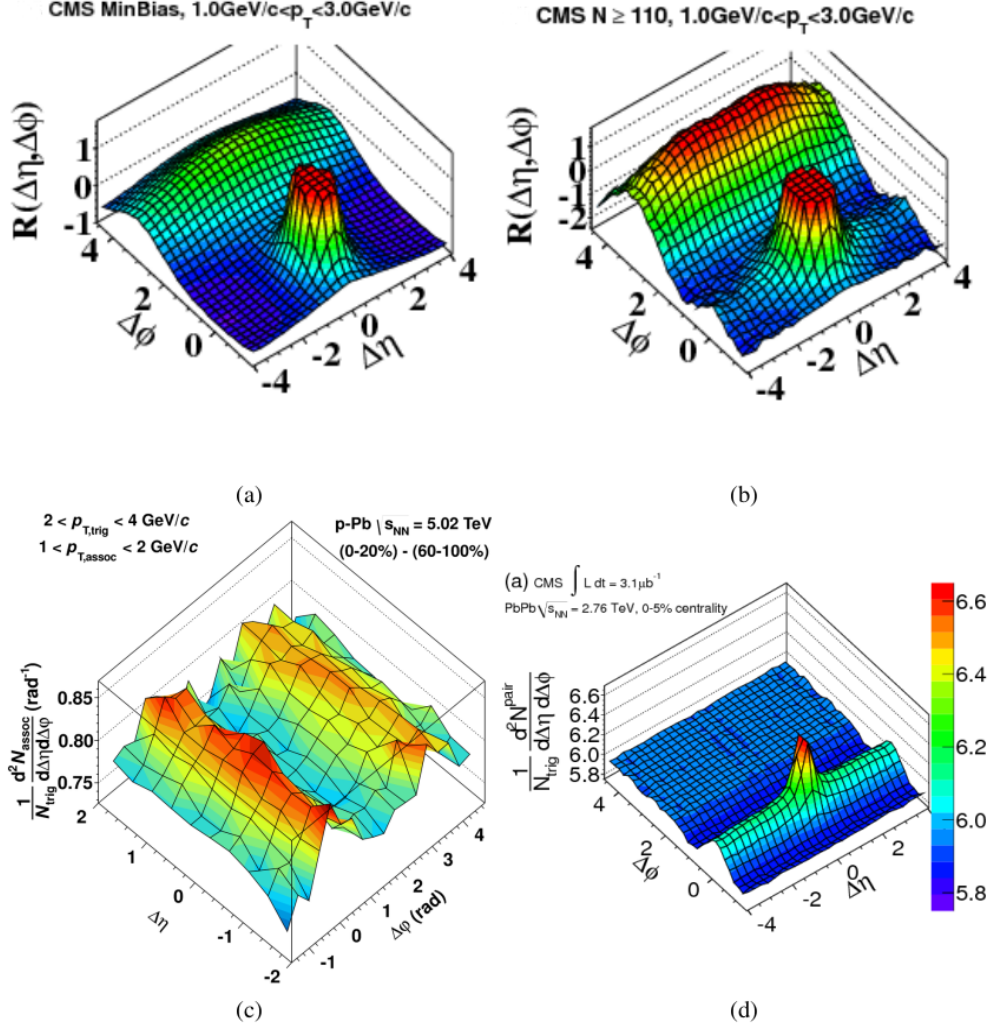


Figure 3.2: $\Delta\eta$ - $\Delta\phi$ angular correlation as measured (a) and (b) in pp collisions at $\sqrt{s} = 7 \text{ TeV}$ [186] for different reconstructed charged particle multiplicities (c) in p-Pb collisions at $\sqrt{s} = 5.02 \text{ TeV}$ [187] for the 0-20 % multiplicity class (the correlation of 60-100 % was subtracted) and (d) in most central Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76 \text{ TeV}$ [188]. One can clearly see the long range ridge at $\Delta\phi \sim 0$ and extending along $\Delta\eta$ for the events where large charged particle densities are produced

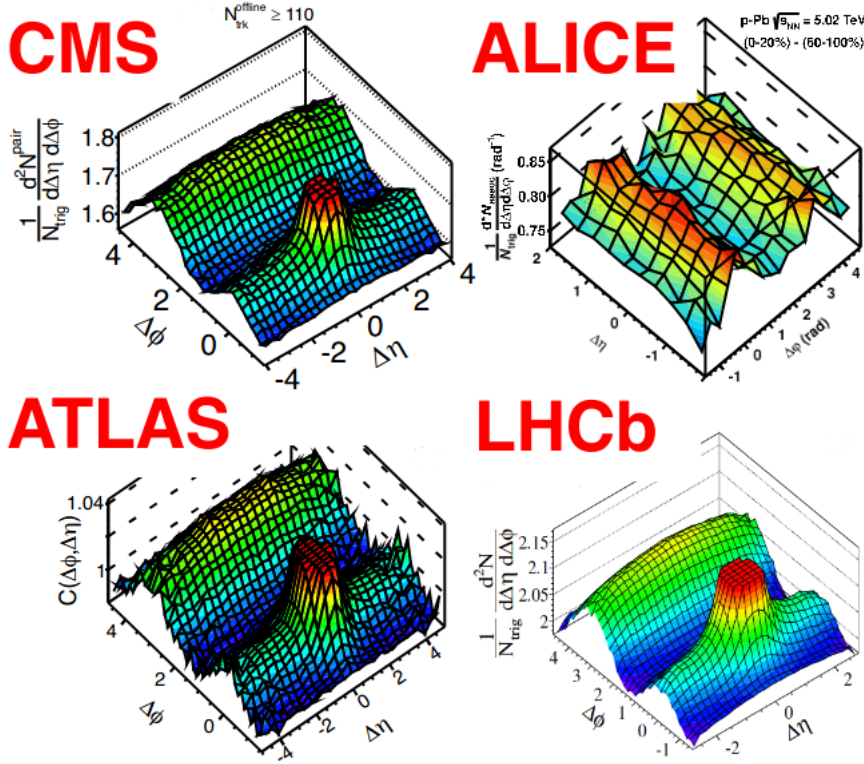


Figure 3.3: $\Delta\eta$ - $\Delta\phi$ angular correlation as measured in p-Pb collisions at CMS [190], ALICE [191], ATLAS [192] and LHCb [193]. One can clearly see the long range ridge at $\Delta\phi \sim 0$ and extending along $\Delta\eta$ for the events where large charged particle densities are produced

ion collisions, is not revealed in low multiplicity pp collisions. However, looking at high multiplicity pp collisions, the ridge appears at near-side for the first time in 7 TeV [189], and also persists at 13 TeV. Even though the ridge is significantly smaller compared to that observed in heavy-ion collisions, it was definitely not expected and needs to be investigated further. The same measurement was performed in low multiplicity and high multiplicity p-Pb collisions (Figure 3.2 (c)) with similar conclusions as in pp collisions: a clear ridge at near-side can be only seen at high multiplicity p-Pb collisions. Moreover, it has more pronounced signal than in pp collisions.

The two-particle correlation measurements are made in different experiments at LHC. The results from all experiments are consistent with each other in Figure 3.3. The ridge appears at near-side in p-Pb collisions when $\Delta\eta$ - $\Delta\phi$ angular correlations

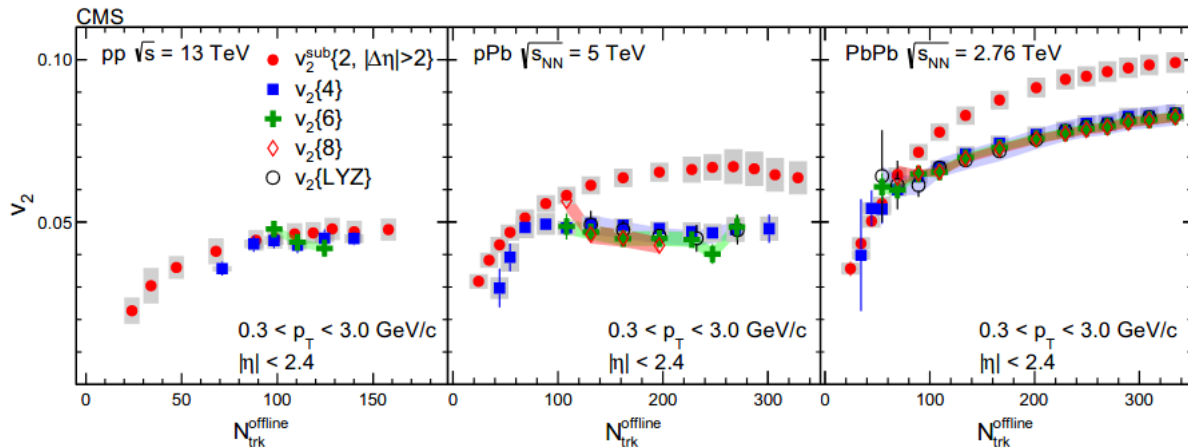


Figure 3.4: Second-order azimuthal anisotropy Fourier harmonics v_2 measured by CMS in pp, p–Pb and Pb–Pb collisions based on correlations among multiple charged particles [194, 195].

are measured in CMS [190], ALICE [191], ATLAS [192] and LHCb [193].

After these observations, it was clear that the possible presence of collectivity in small collision systems needs to be probed further. Whether the observed long-range two-particle correlations come from collective effects, which are understood as effects common to essentially all particles in the bulk of a collision or some other processes are at interplay and mimicking these effects, can be studied with the measurements of multi-particle correlations involving more than just pairs of two particles.

3.2.1.2 Multi-particle correlations

A fundamental question is whether the two-particle azimuthal correlation structures observed at large relative pseudorapidity in pp and p–Pb collisions result from correlations exclusively between particle pairs, or if it is a multi-particle, genuine collective effect. A strong hint, originally named as evidence by the CMS Collaboration, for collective multi-particle correlations in pp and p–Pb collisions was reported already [194, 195]. Figure 3.4 shows the second-order azimuthal anisotropy Fourier harmonics v_2 measured in pp, p–Pb and Pb–Pb collisions over a wide pseudorapidity range based on correlations among up to eight particles. The v_2 values obtained with

3.2. EXPERIMENTAL RESULTS IN SMALL SYSTEMS

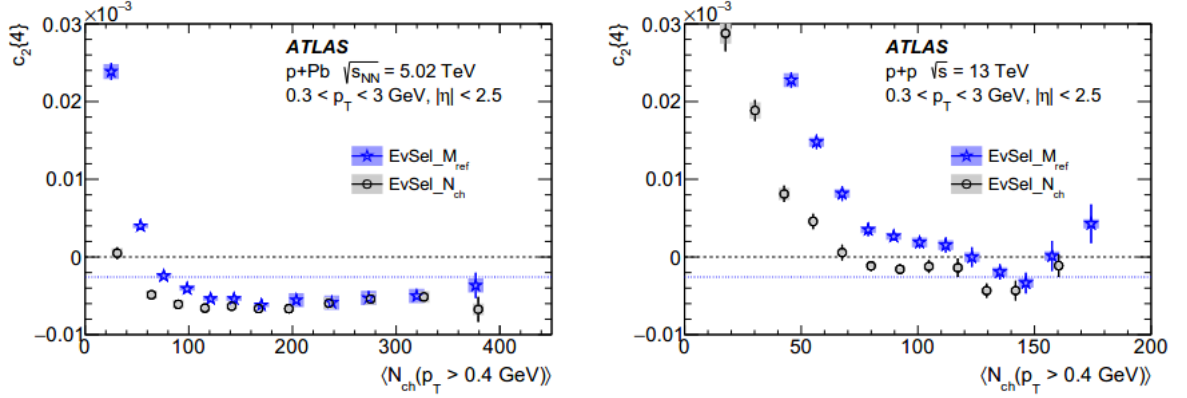


Figure 3.5: Four-particle cumulants $c_2\{4\}$ measured by ATLAS in p–Pb collisions (left) and in pp collisions (right) using different multiplicity selection methods [196].

correlations among more than four particles are consistent with four-particle results and with comparable magnitude to those from two-particle correlations. These observations support the interpretation of a collective origin for the observed long-range correlations in high-multiplicity pp and p–Pb collisions. There is a note of care that must be taken seriously into account. While multi-particle correlation measurements have the advantage of suppressing short-range two-particle correlations such as jets and resonance decays, they are not totally insensitive to such or other so-called “non-flow” effects. This has been stressed in a publication by the ATLAS Collaboration [196], where it is shown how different multiplicity selection methods may yield different four-particle cumulants ($c_2\{4\}$) results. The measurements of $c_2\{4\}$ are shown in Figure 3.5 for pp and p–Pb collisions. The results confirm the evidence for collective phenomena in p–Pb and low-multiplicity Pb–Pb collisions (not reported in Figure 3.5). On the other hand, the pp results for $c_2\{4\}$ do not demonstrate collective behaviour, indicating that they may be biased by contributions from non-flow correlations. Reliably suppressing non-flow correlations in pp collision is a central issue that needs to be solved in order to be able to unveil the actual underlying collectivity. Fortunately, several new methods have been recently developed which

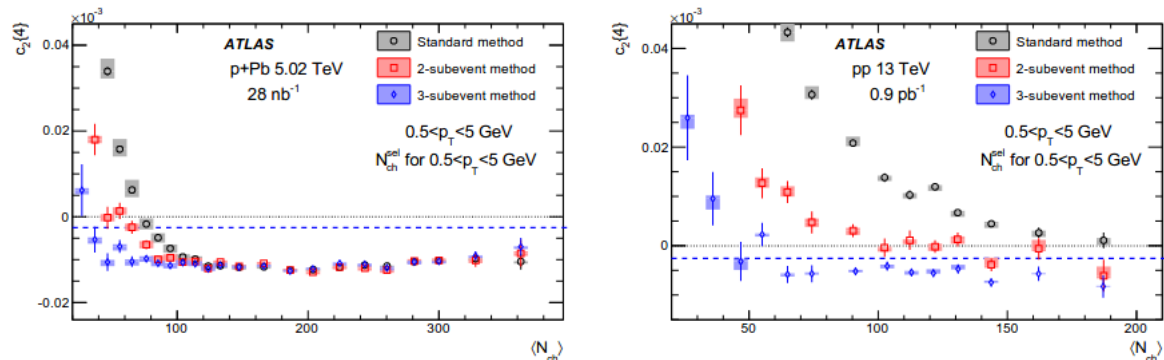


Figure 3.6: Four-particle cumulants $c_2\{4\}$ measured by ATLAS in p–Pb collisions (left) and in pp collisions (right) with the “standard”, “two-subevent” and the “three-subevent” methods [197].

aim at tackling this issue. As an example, the ATLAS Collaboration has reported on the measurement of multi-particle correlations in pp and p–Pb collisions with the so-called “two-subevent” and the “three-sub event” methods [198]. Figure 3.6 shows $c_2\{4\}$ measured in pp and p–Pb collisions with such methods. $c_2\{4\}$ from the standard method is sensitive to the choice of particles used to form the event classes. The sensitivity is greatly reduced in the two-subevent method and is almost fully removed in the three-subevent method, suggesting that the three-subevent method is much more robust against non-flow effects. The three-subevent method shows significant flow in pp collisions in a broad N_{ch} range.

3.2.1.3 Identified hadron spectra

The p_T distributions of identified hadrons are another tool to probe the collective behaviour of particle production. If collective radial flow develops, this would result in a characteristic dependence of the shape of the transverse momentum distribution on the particle mass. The p_T distributions in pp and p–Pb collisions show a clear evolution, becoming harder as the multiplicity increases, with a change which is most pronounced for protons and lambdas. The stronger multiplicity dependence of the spectral shapes of heavier particles is evident when looking at the p_T dependent ratios,

3.2. EXPERIMENTAL RESULTS IN SMALL SYSTEMS

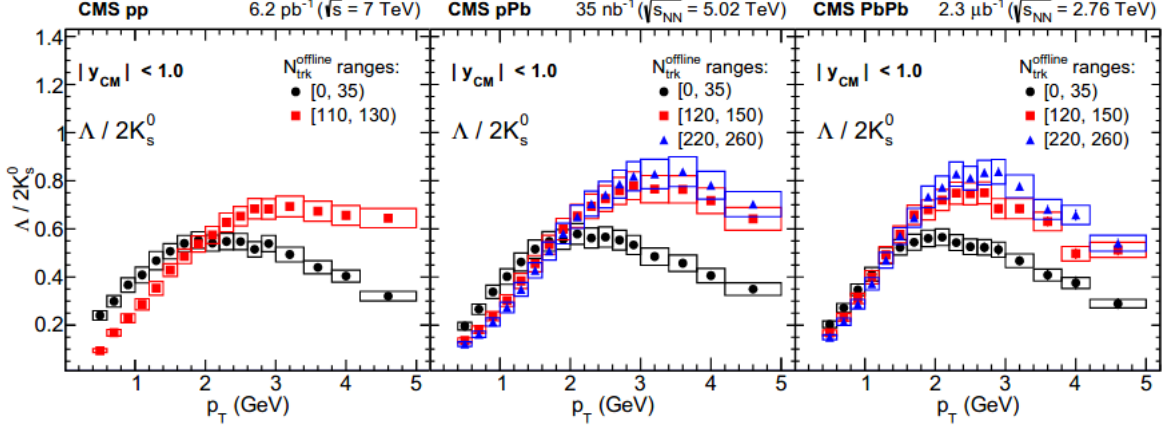


Figure 3.7: Ratios of p_T spectra of $\Lambda/2K_S^0$ measured by CMS in pp (left), p–Pb (middle), and Pb–Pb collisions (right) [177].

such as the Λ/K_S^0 ratio measured by the CMS Collaboration [177], shown in Figure 3.7. The ratios in pp and p–Pb collisions show a significant enhancement at intermediate p_T (~ 3 GeV/ c), quantitatively similar to that measured in Pb–Pb collisions at similar multiplicity. Results from the ALICE Collaboration are shown in Figure 3.8. The figure clearly shows that the modification of the p_T dependent ratios of p/π and of Λ/K_S^0 follows a rather smooth evolution at low and mid p_T across different systems, pointing towards a universal soft mechanism driven by final state $dN_{ch}/d\eta$ [178]. The high- p_T part of the ratio, which is dominated by hard fragmentation, is unchanged.

The Blast-Wave model [199] is a useful tool to characterize the spectral shapes of identified hadrons and test data against the radial flow picture. A simultaneous fit to pions, kaons and protons is performed by the ALICE Collaboration following the approach discussed in [178]. The best fit describes the data with good accuracy in the low- p_T domain, stressing the consistency of particle production from a thermal source expanding with a common transverse velocity. The results of the fit to pp, p–Pb and Pb–Pb data are shown in Figure 3.9 as a function of $dN_{ch}/d\eta$. In the context of the Blast-Wave model, when comparing the parameters of different systems at similar $dN_{ch}/d\eta$, T_{kin} values are similar, whereas β_T is larger for small systems. This might

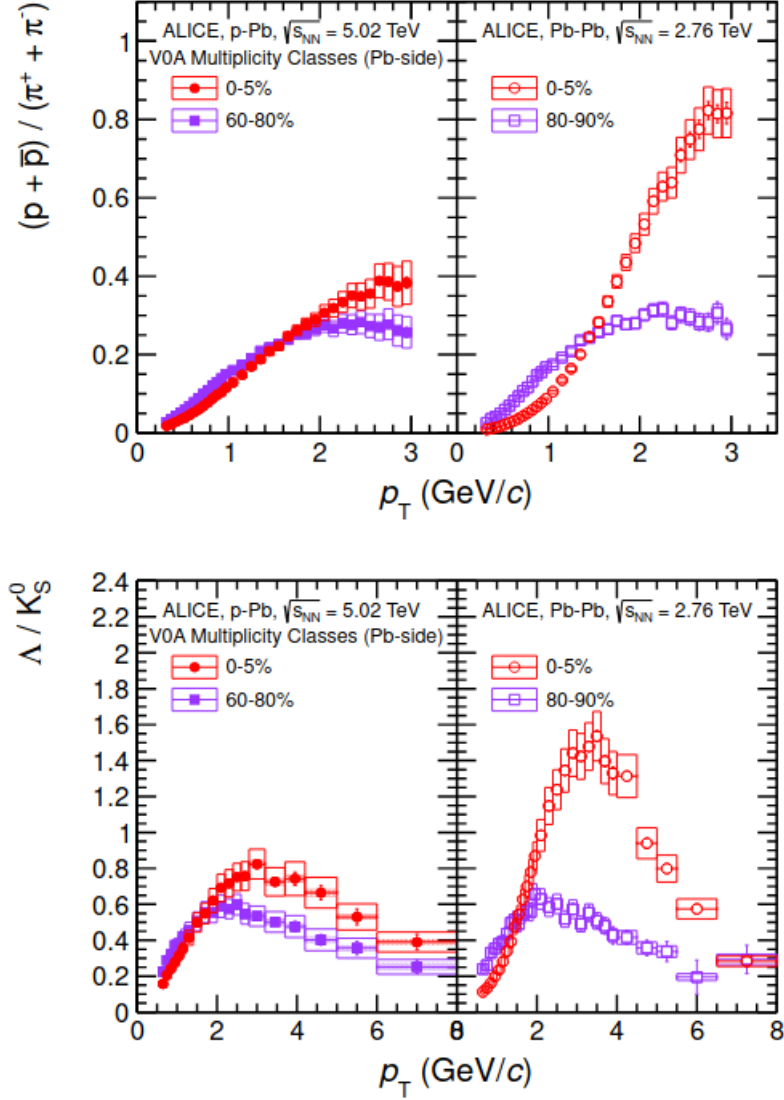


Figure 3.8: Ratios $p/\pi = (p + \bar{p})/(\pi^+ + \pi^-)$ (top plots) and Λ/K_S^0 (bottom plots) as a function of p_T in two multiplicity bins measured in the rapidity interval $0 < y_{CMS} < 0.5$. The ratios are compared to results in Pb–Pb collisions measured at midrapidity, shown in the right panels. The empty boxes show the total systematic uncertainty; the shaded boxes indicate the contribution uncorrelated across multiplicity bins (not estimated in Pb–Pb) [178].

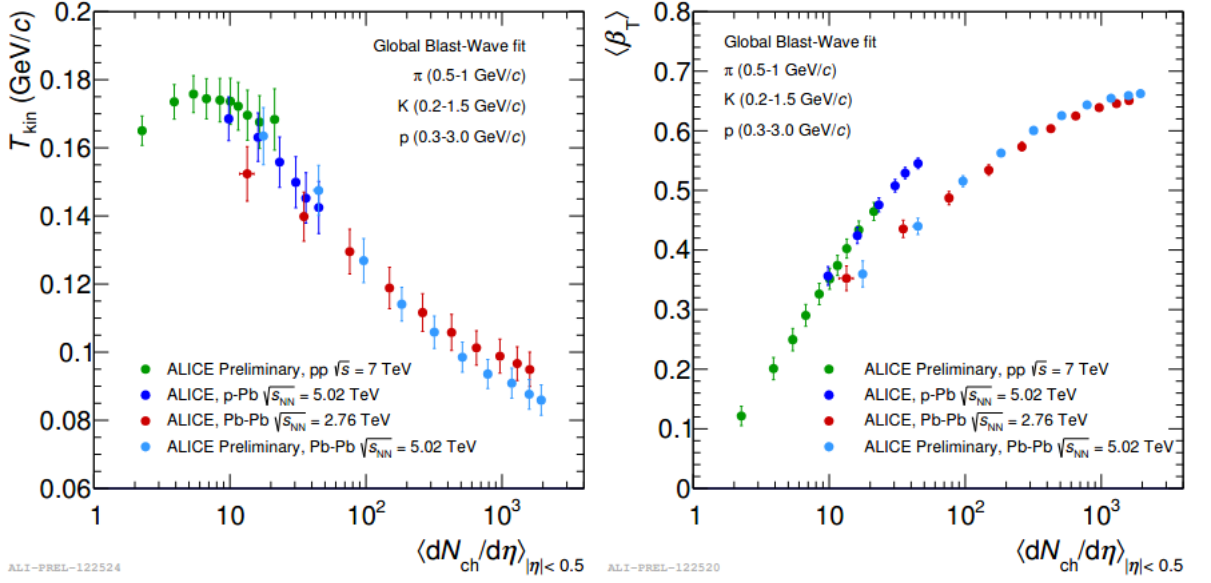


Figure 3.9: Multiplicity dependence of the T_{kin} and β_T Blast-Wave [199] parameters fitted to pion, kaon and proton p_T spectra measured by ALICE in pp, p–Pb and Pb–Pb collisions [200] .

be an indication of a larger radial flow in small systems as consequence of stronger pressure gradients. The ALICE results also show that both T_{kin} and β_T parameters obtained from pp and p–Pb fits are consistent for similar $dN_{ch}/d\eta$ [200].

3.2.2 Multiplicity dependent enhancement of strange and multi-strange hadron production:

In Figure 3.10, the ratios of strange particles (K_S^0 , $\Lambda + \bar{\Lambda}$, $\Xi^- + \bar{\Xi}^+$, $\Omega^- + \bar{\Omega}^+$) to pion yields as a function of the charged particle multiplicity density $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ for different collision systems (pp, p–Pb and Pb–Pb) are shown as measured by ALICE [222]. The lowest concentration of particles containing strangeness is observed in low multiplicity pp events (as expected from the canonical suppression). As the multiplicity increases the particle production seems to evolve without any evident discontinuity, independently on the collision system. The enhancement of strangeness is clearly visible as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ as particles with larger strangeness content are relatively more enhanced (e.g. $\Omega^- + \bar{\Omega}^+$ with respect to K_S^0). This increase

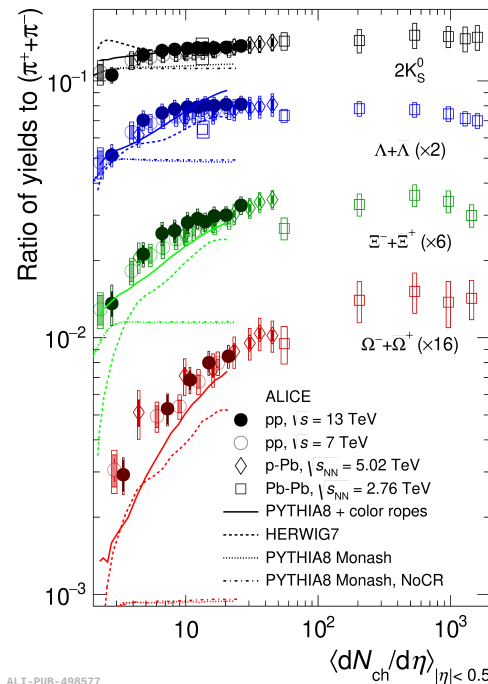
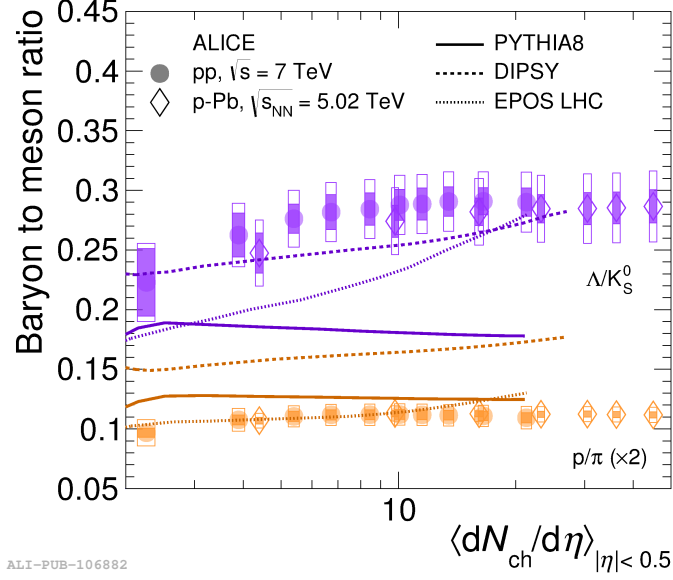


Figure 3.10: p_T integrated yield ratios of strange and multi-strange hadrons to $\pi^+ + \pi^-$ as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ measured in the rapidity interval $|y| < 0.5$ as measured in pp, p-Pb and Pb-Pb collisions [222]

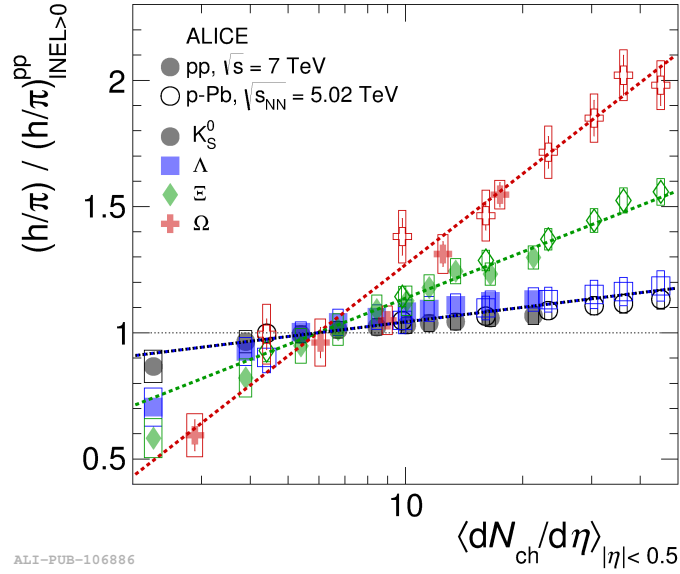
is evident from Figure 3.11b, where the hadron-to-pion ratio is shown scaled to the one measured in the MB pp events [228]. This effect is purely due to the chemical composition of the hadrons under consideration as mesons and baryons with equal strangeness content show the same increase (e.g. K_S^0 and Λ in Figure 3.11a) [228]. Standard event generators such as PYTHIA [223], DIPSY [224] and EPOS LHC [225] are compared with the measured data. DIPSY describes the data well while other models disagree with the data. Modification to the usual mechanisms of particle production (such as color re-connection [226], color ropes [227], string shoving [229]) were proposed to explain such behaviour. These mechanisms show better agreement with the experimental data even.

3.2.2.1 Canonical suppression

The abundances of strange particles in heavy-ion collisions are compatible with those of a hadron gas in thermal and chemical equilibrium and can be described



(a)



(b)

Figure 3.11: (a) $p/(\pi^+ + \pi^-)$ and $(\Lambda + \bar{\Lambda})/(2K_S^0)$ yield ratios as a function of $\langle dN_{ch}/d\eta \rangle$ as measured in pp collisions at $\sqrt{s} = 7$ TeV [228]. Baryon to meson ratios of hadrons with same strangeness show no significant increase. (b) Double ratios $h/(\pi^+ + \pi^-)$ over $h/(\pi^+ + \pi^-)$ in MB collisions[228].

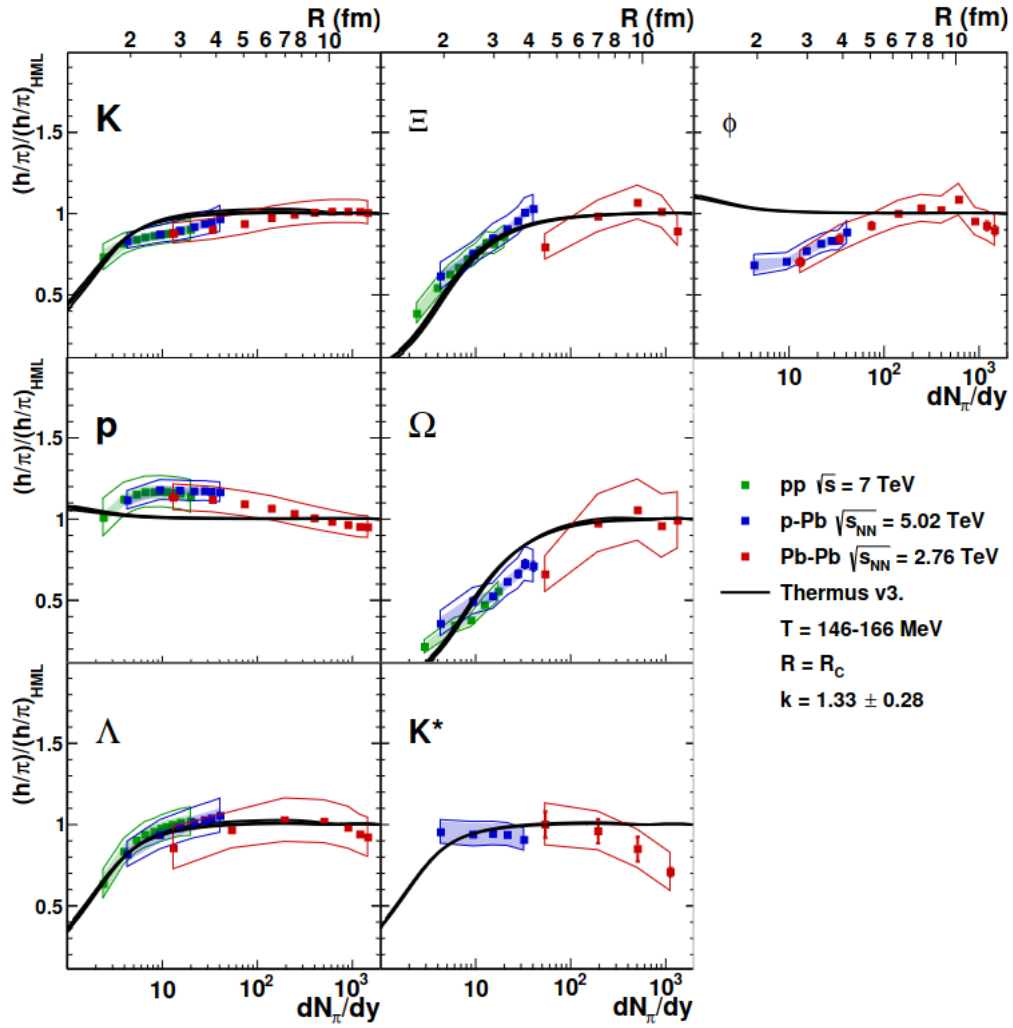


Figure 3.12: Ratios of several particle species measured by ALICE as a function of the midrapidity pion yields for pp, p-Pb and Pb-Pb collisions compared to the THERMUS strangeness canonical suppression model prediction [202].

3.2. EXPERIMENTAL RESULTS IN SMALL SYSTEMS

using a grand canonical statistical model [204, 205]. Extensions of the statistical description, like the ones employed in the strangeness canonical suppression [206] and the core-corona superposition [207, 208] models, can effectively produce a suppression of strangeness production in small systems. The authors in [202] have studied the multiplicity dependence of light flavour hadron production at LHC energies in the strangeness canonical suppression picture using THERMUS [209].

The details can be found therein. As can be observed from Figure 3.12, the study shows that the model provides a good description of the experimental data, except for the ϕ meson. As a matter of fact, the strangeness canonical suppression model is only sensitive to the strangeness quantum number of the hadron. The observation of an enhancement (suppression) of ϕ meson production with increasing (decreasing) multiplicity clearly signals that the particle production rate is not consistent with the one expected from a hadron with zero strangeness quantum number. This can be observed also in Figure 3.13, where the ϕ/K ratio also follows a similar trend in the three collision systems. It is fairly constant as a function of $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta| < 0.5}$, although there is an apparent small increase with $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta| < 0.5}$ from the lowest multiplicities up to $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta| < 0.5} \approx 400$. EPOS-LHC somewhat overestimates the ϕ/K ratio, but is closer to the measured values than PYTHIA, which significantly underestimates ϕ/K . While PYTHIA6 and DIPSY underestimate the ϕ/K ratio, both results exhibit small increases with increasing multiplicity, which is qualitatively similar to the measured trend. The CSM calculation does not describe the behavior of the measured ϕ/K ratio for the $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta| < 0.5}$ range spanned by the ALICE pp measurements.

In addition to comparing the yields of ϕ to pions and kaons, the yields of ϕ is also compared to Ξ . These two particles contain the same number of strange valence (anti-)quarks: ϕ is a $s\bar{s}$ bound state and Ξ contains two strange valence quarks.

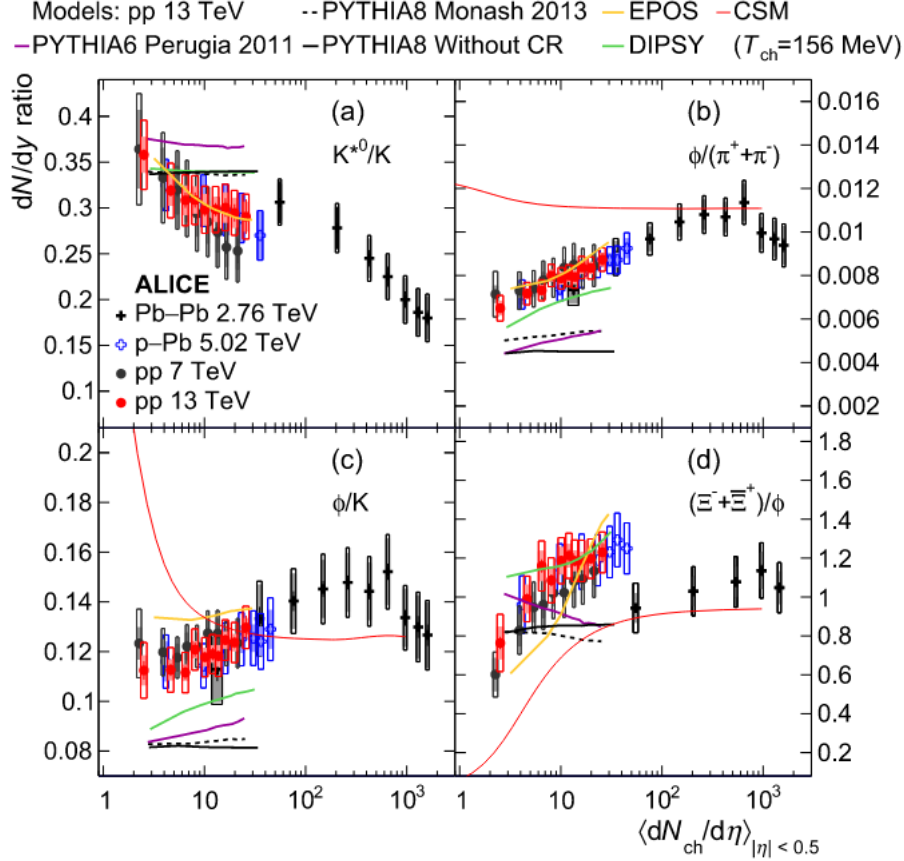


Figure 3.13: Ratios of p_T integrated particle yields K^{*0}/K , ϕ/π , ϕ/K and Ξ^-/ϕ in pp collisions at $\sqrt{s} = 13$ TeV as functions of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ [210, 211]. These measurements are compared with data from pp collisions at $\sqrt{s} = 7$ TeV [212], p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [213, 214], and Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [215, 216]. The widths of the boxes for pp collisions at $\sqrt{s} = 13$ TeV and p-Pb collisions do not represent the uncertainties of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$. The measurements for pp collisions at $\sqrt{s} = 13$ TeV are also compared to results from common event generators [217–220] and a Canonical Statistical Model calculation [212].

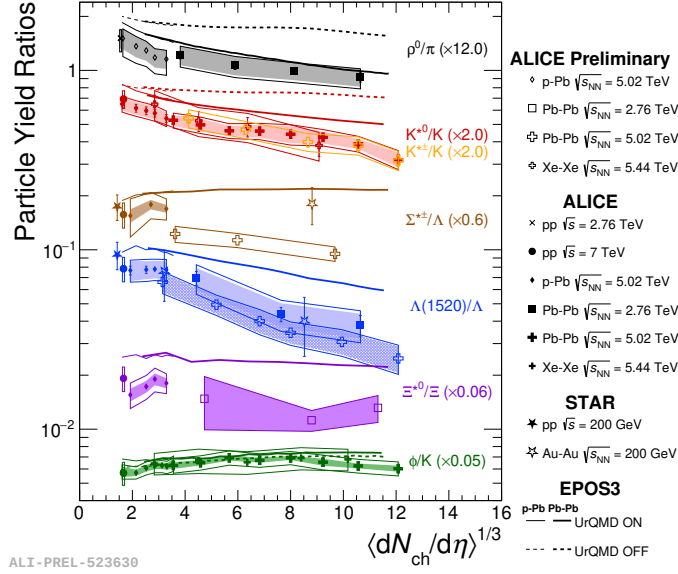


Figure 3.14: Resonance yield ratio to long-lived hadrons for pp, p-Pb, Xe-Xe, and Pb-Pb collisions, with comparisons to EPOS3 predictions and STAR data as a function of charged particle multiplicity ($\langle dN_{ch}/d\eta \rangle$). The bars and the lines represent the statistical and systematic errors, respectively[243].

However, Ξ would be subject to canonical suppression, unlike the strangeness-neutral ϕ . Figure 3.13 also shows the Ξ/ϕ ratio in pp, p-Pb, and Pb-Pb collisions. The ratio increases with increasing $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ for low-multiplicity collisions and is then fairly constant for a wide range of multiplicities: from pp and p-Pb collisions at $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5} \approx 7$ to central Pb-Pb collisions. There is a possible small increase in the Ξ/ϕ ratio from $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5} \approx 7$ to the highest-multiplicity p-Pb collisions. Nevertheless, there is no clear increase in the ratio for $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5} \geq 7$. The decrease in Ξ/ϕ with decreasing $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ for low multiplicities could be interpreted as evidence of canonical suppression in small systems; the canonical statistical model predicts a decrease in the $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ ratio with decreasing $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ that is qualitatively similar to the measured data. However, canonical suppression would also result in an increase in the ϕ/K ratio with decreasing $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$, which is not observed. Given that Ξ and K have different numbers of strange valence (anti)quarks, it is expected that Ξ would be more affected by

canonical suppression [212]. Thus, the multiplicity evolution is not consistent with simple descriptions of canonical suppression. It will be interesting to extend the study of the ϕ/K ratio to lower multiplicities to test if there is any increase in this ratio due to canonical suppression of kaon yields.

3.2.3 Re-scattering effects:

In order to look for the presence of re-scattering effects in high multiplicity p-Pb collisions; the resonance yields relative to long-lived hadrons is studied as a function of the cube root of the average mid-rapidity charged particle density. This ratio as a function of the cube root of the average mid-rapidity charged particle density for ρ^0/π , K^{*0}/K , $\Sigma^{*\pm}/\Lambda$, Λ^*/Λ , Ξ^{*0}/Ξ^- and ϕ/K in pp, p-Pb, Xe-Xe, and Pb-Pb collisions are shown in Figure 3.14. This plot shows the particle ratios along with MC model comparisons. Significant suppression is observed for ρ^0/π , K^{*0}/K , and Λ^*/Λ with charged particle multiplicity in central Pb-Pb collisions. This is qualitatively explained by the EPOS generator with UrQMD, which attributes the suppression to the decay daughters in the hadronic phase of the collision [240, 241]. For pp collisions, the decreasing trend of K^{*0}/K with charged particle multiplicity hints at the possible presence of a hadronic phase with non-zero lifetime in small collision systems [242] and has been explained as a consequence of a re-scattering of K^{*0} decay daughters in the hadronic phase. It is worth noticing that a similar trend is also observed in p-Pb collisions.

3.3 Sources of the QGP-like effects in MC

Similar observations in small collision systems do not necessarily have the same origin as in heavy-ion collisions. Effects of the initial stage, as well as interactions in the hadronic stage, might also result in final particle anisotropy. In order to distinguish various explanations of the experimental observations, a careful comparison between data and theory is in place. Extensive studies in heavy-ion collisions proved

the existence of the QGP and its fluid-like character to be responsible for the reported results. It is one of the main interests nowadays to put similar efforts into the understanding of the phenomena observed in (high multiplicity) collisions of small systems. In case collectivity in small collision systems is established, it is natural to continue with the exploration of its origin and establish whether it arises from initial state effects, final state effects, or a combination of the two.

Several mechanisms were proposed to explain the origin of these collective-like effects, among these the colour re-connection [244], the jet-medium [245] and multi-parton induced [246, 247] interactions, or the possibility to form a high-density system in these collisions which could cause collective effects [248, 249]. These discoveries started a very rich research field which is finding more and more similarities between small and large systems even connecting observables specific to the QGP such as the particle composition, collective phenomena and the strangeness enhancement.

Hydrodynamic calculations reproduce many of the observations qualitatively [250]. However, it has also been found that multi-parton interactions (MPI) [251] and color reconnection (CR) as implemented in Pythia [252] produce radial flow patterns via boosted color strings [253]. Moreover, within the dilute-dense limit of the color glass condensate, it has been demonstrated that the physics of fluctuating color fields can generate azimuthal multi-particle correlations [254]; and the mass ordering of elliptic flow when the fragmentation implemented in Pythia is included [255]. The same observable has been studied using the multi-phase transport model [256], where the ridge structure can be generated assuming incoherent elastic scatterings of partons and the string melting mechanism. Other mechanisms like “color ropes”, which are formed by the fusion of color strings close in space, can increase both the strangeness production and the radial flow-like effects [257]. The measurements of the transverse momentum (p_T) spectra of identified particles as a function of event

multiplicity in pp collisions at the LHC [258, 259] have shown that models fail to describe the data quantitatively. Therefore, the results of those comparisons alone are not enough to give desired information about the origin of the observed effects (i.e. radial flow-like patterns). In order to extract more information, it has been proposed for the implementation of a differential study based on the classification of the events according to event multiplicity and the jet content.

Chapter 4

This thesis work is carried out in the context of one of the experiments at LHC namely the ALICE (A Large Ion Collider Experiment). The extensive PID capabilities of this detector to study various observables related to particle production will be elaborated in this chapter. Moreover, the LHC Run3 upgrades and its physics motivation will also be discussed.

Chapter 4

The ALICE Experiment

4.1 The LHC Project

The Large Hadron Collider (LHC) [260] at CERN is the biggest particle accelerator in the world. The first discussions that led to the project started in 1984 and was approved in 1994. The construction works in the existing underground tunnel started in 2001 after the dismantling of the LEP collider [261], which had previously been built in the tunnel which is located under the Swiss-French border area close to Geneva at a depth of 50 to 175 m. The LHC has a circumference of 27 km. By design, its maximum achievable energies are 7 TeV for beam of protons and for beam of lead ions it is 2.76 TeV per nucleon, thus providing collisions at $\sqrt{s} = 14$ TeV and $\sqrt{s_{\text{NN}}} = 5.5$ TeV, respectively. These are the largest energies ever achieved in particle collision experiments.

The LHC is a synchrotron that accelerates two counter-rotating beams in separate parallel beam pipes. In each of them, the bunches of particles travel many times around the accelerator ring before the collision energy is reached. The accelerator bends the beams around the ring, keep the bunches focused and accelerate them to their collision energy. Finally, the spatial dimension of the bunches has to be

CHAPTER 4. THE ALICE EXPERIMENT

minimized, which ensures a higher probability of collisions per time interval at the collision points, i.e. a high luminosity. A combination of magnetic and electric field components performs the mentioned tasks. Despite the high luminosity reached, only a very small fraction of the particles of two bunches collide in a single bunch crossing. Most of them leave the interaction region essentially uninfluenced, are defocused, and continue to circulate in the accelerator.

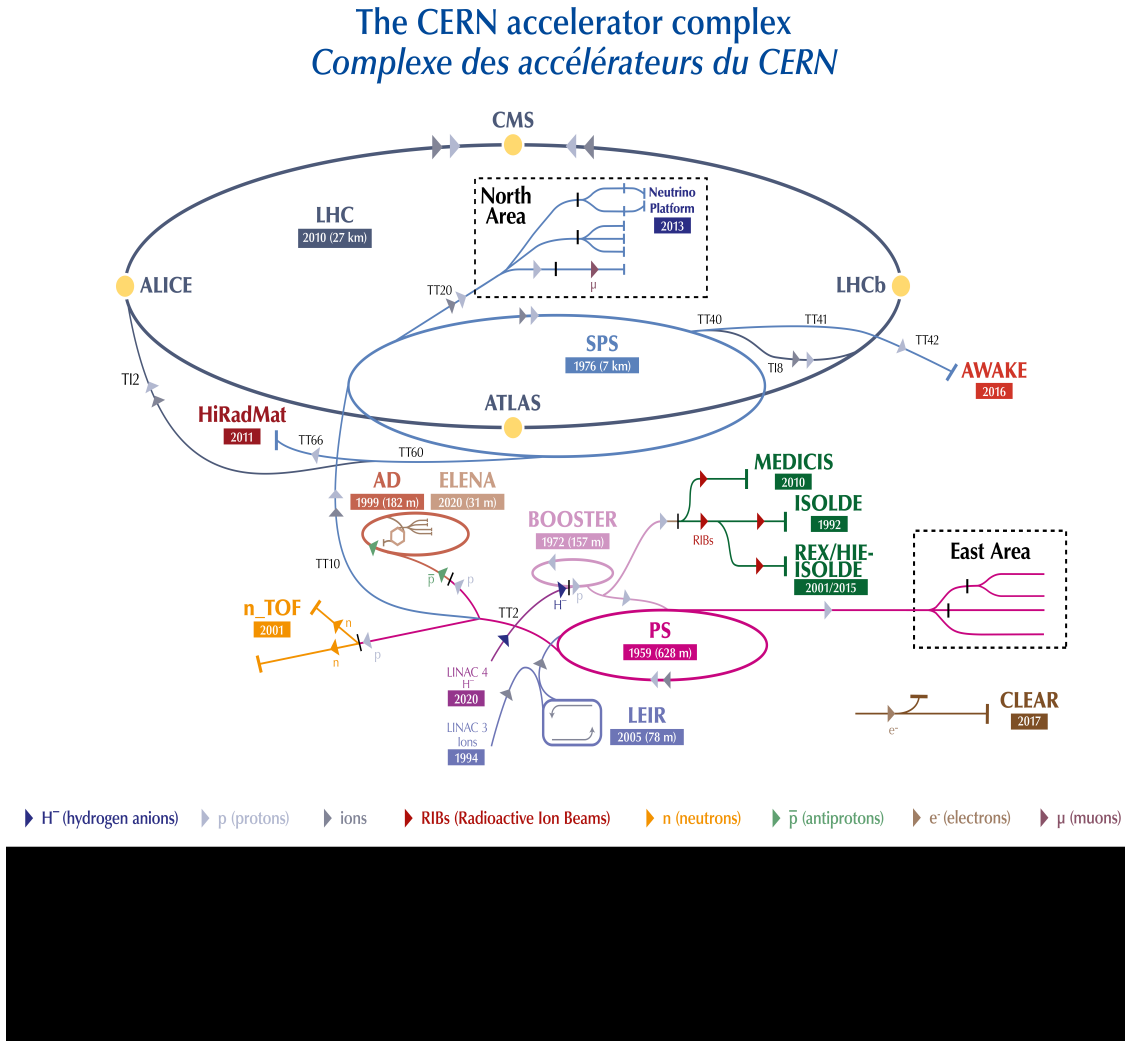


Figure 4.1: The LHC is the last ring (dark blue line) in a complex chain of particle accelerators. The smaller machines are used in a chain to help boost the particles to their final energies and provide beams to a whole set of smaller experiments, which also aim to uncover the mysteries of the Universe [262]

The process of injection of bunches into the LHC (Figure 4.1) is preceded by acceleration in the LINAC2, PS booster, PS, and SPS accelerators. The acceleration sequence is slightly different for heavy-ions, in which case the bunches pass the LINAC3, LEIR, PS, and SPS accelerators (more information can be found in [263]). Several injections to the LHC are needed until all bunches of both beams are filled.

On 10th of September, 2008 the first pp collisions at 900 GeV centre-of-mass energy were delivered by the LHC. Nine days later, the operations were interrupted due to some failure in an electrical connection between two magnets. The machine operators spent over a year repairing and consolidating the accelerator. On 20th of November, 2009 low energy proton beams circulated again, and a few days later, by achieving the energy of 1.18 TeV per proton beam, LHC became the most powerful accelerator in the world. The first pp collisions at centre-of-mass energy of 7 TeV were delivered in March 2010, and the first Pb–Pb collisions at centre-of-mass energy of 2.76 TeV per nucleon pair in November 2010.

In 2010 the integrated luminosity delivered by the LHC was $\sim 48 \text{ pb}^{-1}$ for pp collisions at $\sqrt{s} = 7 \text{ TeV}$ ($\sim 0.5 \text{ pb}^{-1}$ in ALICE) and $\sim 10 \mu\text{b}^{-1}$ for Pb–Pb at $\sqrt{s_{\text{NN}}} = 2.76 \text{ TeV}$ ($\sim 10 \mu\text{b}^{-1}$ in ALICE). In 2011 the beam energy was the same as in 2010 both for pp and Pb–Pb. The performance of the LHC improved in terms of luminosity with $\sim 5.61 \text{ fb}^{-1}$ for pp ($\sim 2 \text{ pb}^{-1}$ in ALICE) and $\sim 166 \mu\text{b}^{-1}$ for Pb–Pb collisions ($\sim 143.62 \mu\text{b}^{-1}$ in ALICE). In 2012, the centre-of-mass energy for pp collisions was brought to 8 TeV and the integrated luminosity (up to December 2012, end of the pp programme) was $\sim 23.3 \text{ fb}^{-1}$ ($\sim 10 \text{ pb}^{-1}$ in ALICE).

A pilot p–Pb run operated at $\sqrt{s_{\text{NN}}} = 5.02 \text{ TeV}$ on September 2012, followed by a long p–Pb run on February 2013 with a delivered luminosity of 31.2 nb^{-1} . A very short pp run at $\sqrt{s} = 2.76 \text{ TeV}$ ended the Run1 of the LHC programme, marking the start of the first long shutdown (LS1) until the end of 2014.

A long shutdown (LS1) then commenced where many of the detectors, and the LHC itself, were upgraded. The LHC then restarted in April 2015 for Run 2 which continued until October 2018. Higher energy collisions of $\sqrt{s} = 13$ TeV for pp collisions, $\sqrt{s_{\text{NN}}} = 5.02$ TeV for Pb–Pb collisions and $\sqrt{s_{\text{NN}}} = 8.16$ TeV for p–Pb collisions, were carried out during this run, as well as a brief run of Xe–Xe collisions at $\sqrt{s_{\text{NN}}} = 5.44$ TeV.

On July 5, 2022 [264], Large Hadron Collider reached the highest energy proton–proton collisions at 13.6 TeV, after more than three years of upgrade. The data taking started with a fully revamped ALICE detector recording the first 13.6 TeV collisions of Run 3.

4.1.1 Proton Acceleration

Hydrogen is the source of protons for the LHC. In the presence of electric field, the electrons are ripped off from the hydrogen atoms and protons are obtained. The first stage for acceleration is accomplished by the LINAC2 linear accelerator, which accelerates the proton beam up to the energy of 50 MeV. The beam is then fed into the Proton Synchrotron Booster (PSB), a circular synchrotron accelerator of 157 m in circumference, which accelerates the proton beam up to 1.4 GeV. Afterwards, this boosted beam is injected into the Proton Synchrotron (PS) (628 m in circumference), which pushes the beam energy up to 25 GeV. This 25 GeV proton beam is further fed into the Super Proton Synchrotron (SPS) (7 km in circumference) where it is accelerated up to 450 GeV. Finally, the beam is injected into the two beam pipes of the LHC, the final accelerator of 26.7 km in circumference. The beam circulate clockwise in one pipe and anticlockwise in the second pipe. It takes 4 minutes and 20 seconds to fill each LHC ring with protons and 20 minutes for the LHC to accelerate the protons to attain their maximum energy of 6.5 TeV. Both beams circulate for hours under normal operating conditions and finally collide at centre-of-mass energy

$\sqrt{s} = 13$ TeV at the four interaction points where the experimental detectors are located. The detectors installed at four interaction points are ALICE at Point 2 [273], ATLAS at Point 1 [274], CMS at Point 5 [275] and LHCb [276] at Point 8 [265].

4.1.2 Lead Ions Acceleration

The LHC has so far accelerated only one isotope of Pb, the one that contains 126 neutrons and 82 protons. The lead-ion beam is generated starting from a pure lead piece of 2 cm in length and 500 milligrams in weight, in the Electron Cyclotron Resonance (ECR) source. In the ECR source, this $^{208}\text{Pb}_{82}$ piece is heated up to 500°C to form Pb atom vapors. At the ECR stage, an electric current is used to remove a few electrons from the lead atoms to form Pb^{27+} ions. The newly created ions begin their ride after being transferred into the linear accelerator LINAC 3. At this stage, the Pb^{27+} ions are accelerated up to 4.4 MeV/nucleon and more electrons are stripped off in the carbon foil stripping chamber to convert them into Pb^{54+} ions. These Pb^{54+} ions are then injected into a first circular accelerator, the Low Energy Ion Ring (LEIR) of 78.54 m in circumference, where they are accelerated up to 72 MeV/nucleon. The accelerated Pb^{54+} ions are injected into the PS for further acceleration of up to 5.9 GeV/nucleon, where a stripping Al foil (~ 1 mm in thickness) converts them to fully ionized Pb^{82+} ions. The ion beam is then fed into the SPS for further acceleration of up to 177 GeV/nucleon and finally transferred into the beam pipes of the LHC ring. It takes 10 minutes to fill the LHC ring with Pb^{82+} ions, where the ion beams attain maximum energy of 1.38 TeV/nucleon and finally collide at the interaction points at centre-of-mass energy, $\sqrt{s_{\text{NN}}} = 2.76$ TeV [266–270].

4.1.3 Lead Ion Collisions

In the year 2012, the protons and lead ions were made to collide with each other at $\sqrt{s_{\text{NN}}} = 5.02$ TeV for the first time and the collision data has been recorded for the

physics investigation [271]. The beam preparation for protons and Pb ions are the same as explained in the previous two sections. In the year 2015, the p-Pb collision data has been recorded at $\sqrt{s_{\text{NN}}} = 8.16$ TeV in the LHC for the second time. The centre-of-mass energy of 5.02 and 8.16 TeV in p-Pb collisions are achieved with the LHC magnet setting for proton beam at 4 and 6.5 TeV, respectively [272].

4.1.4 LHC Experiments

The LHC produces collisions in four so called Interaction Points (IPs) in correspondence to which there are eight experiments at the Large Hadron Collider (LHC) which use detectors to analyse the myriad of particles produced by collisions in the accelerator. These experiments are run by collaborations of scientists from institutes all over the world. Each experiment is distinct, and characterised by its detectors.

ALICE (A Large Ion Collider Experiment) [273] is a dedicated heavy-ion experiment designed to study strongly-interacting matter at very high energy density. It explores the phase transition to the QGP, its phase diagram, and its properties. Furthermore, ALICE also studies collisions of protons, on one hand as a baseline for heavy-ion measurements and on the other hand it contributes to measurements of identified particles by making use of its excellent particle identification capability and its acceptance at very low transverse momenta.

ATLAS (A Toroidal LHC Apparatus) and CMS (Compact Muon Solenoid) [274, 275] are general-purpose detectors that are built to cover the widest possible range of physics at the LHC. Specific topics are the studies related to the Higgs boson and physics beyond the Standard Model, e.g. new heavy particles postulated by supersymmetric extensions (SUSY) of the Standard Model and evidence of extra dimensions.

LHCb (The Large Hadron Collider beauty experiment) [276] is a dedicated experiment for the study of heavy flavour physics at the LHC. In particular, the

experiment focuses on the study of CP violation and rare decays of beauty and charm particles, to test the Standard Model and to search for evidence of New Physics. The LHCb physics programme is complementary to the flavour physics studies conducted at the B-factories and to the direct searches for new particles performed at ATLAS and CMS.

LHCf (Large Hadron Collider forward experiment) [277] measures forward particles created during LHC collisions to provide further understanding of high energy cosmic rays. The detector is placed close to the ATLAS experiment.

TOTEM (TOTAl Elastic and diffractive cross-section Measurement) [278] measures the total cross-section, elastic scattering, and diffractive processes. The detector is located close to the CMS experiment.

MoEDAL (Monopole and Exotics Detector at the Large Hadron Collider) and FASER (Forward Search Experiment)[279] MoEDAL uses detectors deployed near LHCb to search for a hypothetical particle called the magnetic monopole. FASER, the newest LHC experiment, is situated 480 metres from the ATLAS collision point in order to search for light new particles and study neutrinos.

4.2 The ALICE Project

The ALICE detector is optimised to study the high multiplicities produced in heavy-ion collisions[280]. It is designed to study nuclear matter at extreme conditions of temperature and energy density in which a QGP is produced. It is a 10,000 tonne detector with dimensions of 16 by 16 by 26 m^3 . It is made up of eighteen different detector sub-systems which are designed to provide a broad range of measurements that can be used to analyse the heavy-ion collisions at the LHC. It consists of a central barrel encased in a solenoidal magnet with a field of 0.5 T and a forward muon spectrometer. The layout of the detectors within ALICE is shown in Figure 4.2.

THE ALICE DETECTOR

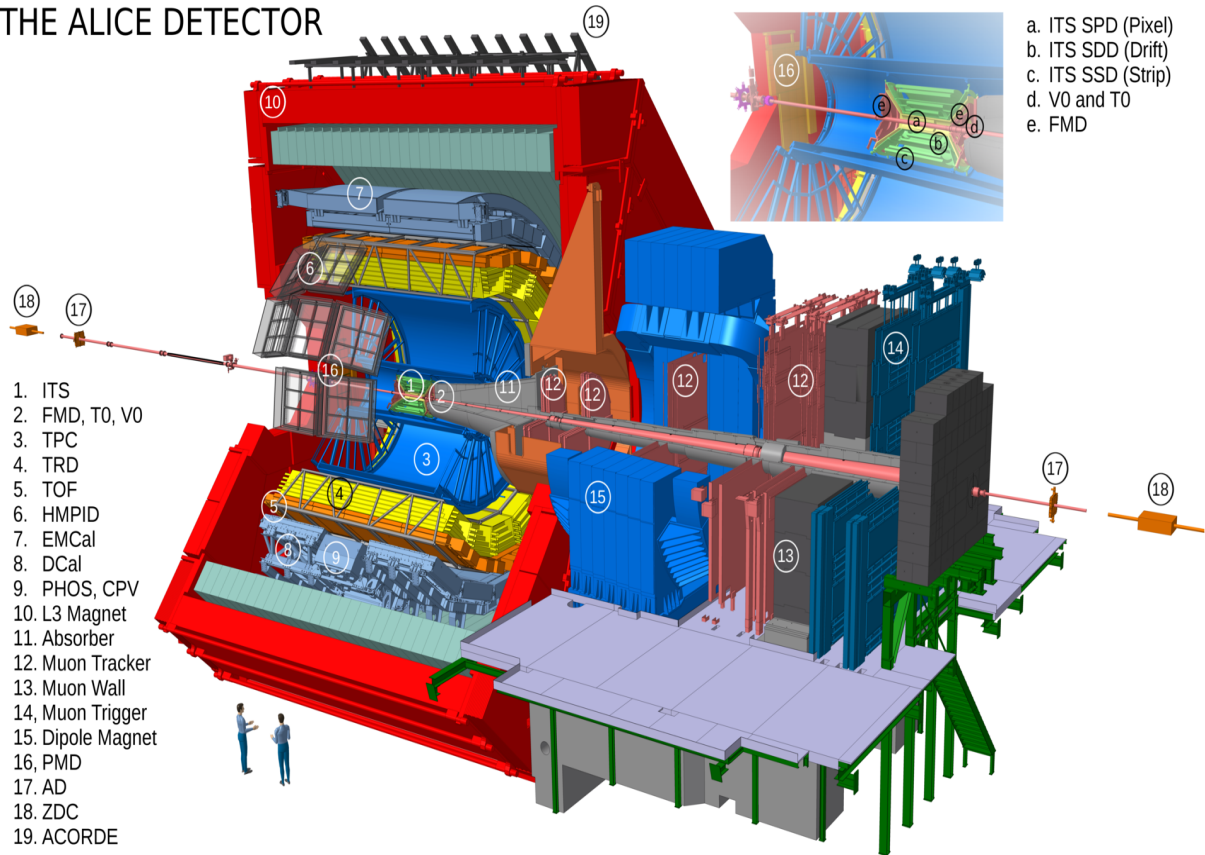


Figure 4.2: Schematic view of ALICE Experiment

4.2.1 ALICE Coordinates

The ALICE detector follows a right-handed Orthogonal Cartesian coordinate system [281] with the origin $(x, y, z = 0)$ at the very centre of the detector. The x axes is defined to be pointing to the centre of the LHC, the y axes points up and the z axes is parallel to the beam pipe pointing away from the muon spectrometer, as indicated in Figure 4.3. The variables studied can be divided into transverse and longitudinal, where the longitudinal components refer to the z -components and the transverse components combine the x and y components. For example, the transverse momentum is defined by $p_T = \sqrt{p_x^2 + p_y^2}$. As the p_T of the colliding nucleons is approximately 0, transverse momentum in the final state is of interest as it is created

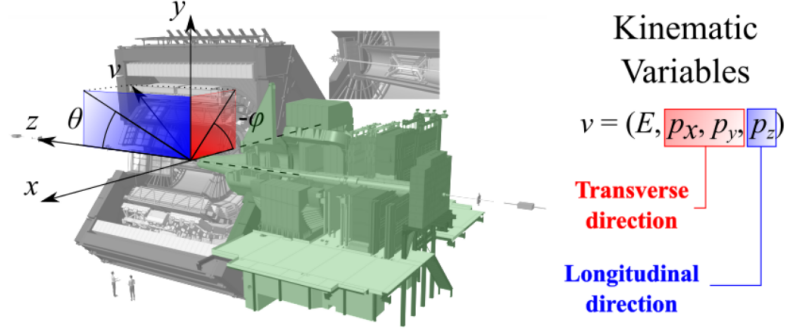


Figure 4.3: ALICE coordinates system and relevant variables. Adapted from [282]

in the collision. The longitudinal variable pseudorapidity is given by

$$\eta = \frac{1}{2} \ln \left(\frac{|\mathbf{p}| + p_z}{|\mathbf{p}| - p_z} \right) = -\ln \left[\tan \left(\frac{\theta}{2} \right) \right] \quad (4.1)$$

where $\mathbf{p}$ is the momentum vector for a particle, p_z is the momentum along the z axis and θ is the angle between the particle's momentum and the z axis. $\eta = 0$ when $\theta = 90^\circ$ ie. along the y axis. At very high relativistic energies, η is a good approximation of the invariant variable rapidity defined as

$$y = \frac{1}{2} \left(\frac{E + p_z}{E - p_z} \right) \quad (4.2)$$

These variables can be used to define different rapidity regions. Particles produced close to the beam line with small θ have $\eta \geq 3$ (where $\theta = 10^\circ$ for $\eta = 2.44$) which is known as the forward region. Particles associated with the transverse region have $|\eta| < 0.5$, which is known as the mid-rapidity region. The side of ALICE with the positive z -axis and y , $\eta > 0$ is referred to as the A-side and the opposite region with y , $\eta < 0$ is referred to as the C-side. In this thesis the analysis takes place on p-Pb collisions in which the protons travel from A to C and the lead ions travel from C to A. In symmetric systems like Pb-Pb and p-p, the centre of mass rapidity y_{CM} and lab rapidity y_{lab} are the same. However, for p-Pb collisions in the lab frame, the

centre of mass frame is moving with a shift in rapidity given by

$$\Delta y \approx \frac{1}{2} \ln \left[\frac{Z_1 A_2}{Z_2 A_1} \right] \approx \frac{1}{2} \ln \left[\frac{82 \times 1}{1 \times 208} \right] \approx 0.465 \quad (4.3)$$

where Z and A are the atomic and mass numbers of the colliding particles [283]. For this system, the shift is towards negative rapidities so at $y_{CM} = 0$ (midrapidity), $y_{lab} = -0.465$.

4.2.2 ALICE Subdetectors

The ALICE detector consists of 18 different subdetectors divided into three groups: central barrel detectors, forward detectors and the muon arm. The central barrel covers the pseudo-rapidity interval $-0.9 < \eta < 0.9$ and contains detectors used for vertex reconstruction, tracking, particle identification (PID) and momentum measurements. It consists of the inner tracking system (ITS), the time projection chamber (TPC), the transition radiation detector (TRD) and the Time of Flight (TOF). In the mid-rapidity region there is also the High Momentum Particle Identification Detector (HMPID), the Photon Spectrometer (PHOS) and the Electromagnetic Calorimeter (EMCAL).

The forward detectors are in the high pseudo-rapidity region and therefore at a small angle with respect to the beam pipe. They are used for triggering and measuring global event characteristics. They consist of the Time Zero (T0), VZERO (V0), Forward Multiplicity Detector (FMD), Photon Multiplicity Detector (PMD) and Zero Degree Calorimeter (ZDC). The T0 measures the event time with a precision of the order of 10^{-11} s. The FMD and PMD can provide multiplicity information and the ZDC is used to differentiate peripheral and central collisions by measuring the energy of the spectator nucleons that do not participate in the collisions.

In the forward pseudo-rapidity region ($-4 < \eta < -2.5$) also lies the muon spectrometer, which has a front absorber, a dipole magnet as well as tracking and trigger

chambers, and is used to study heavy quarkonia, open charm, W and Z bosons via their muon decays. Finally, the ACORDE detector is an array of scintillators on top of the ALICE solenoidal magnet that is used for studies of cosmic rays. In the following sections more detailed descriptions of the detectors used in the analysis in this thesis are given.

4.2.2.1 ITS

The Inner Tracking System (ITS) [285, 286] is a high-resolution tracking detector made up of six layers of silicon detectors (2 SPD, 2 SDD and 2 SSD) illustrated in Figure 4.4. The ITS can locate the primary and secondary vertices with a high resolution, which is important for strangeness analyses that measure the decays of hyperons with decay lengths of the order of centimetres. The four outermost layers also provide dE/dx measurements for low-momentum particles that can be used for particle identification (PID). The ITS surrounds the beam pipe between $4 < r < 43$ cm, covering the rapidity range $|\eta| < 0.9$.

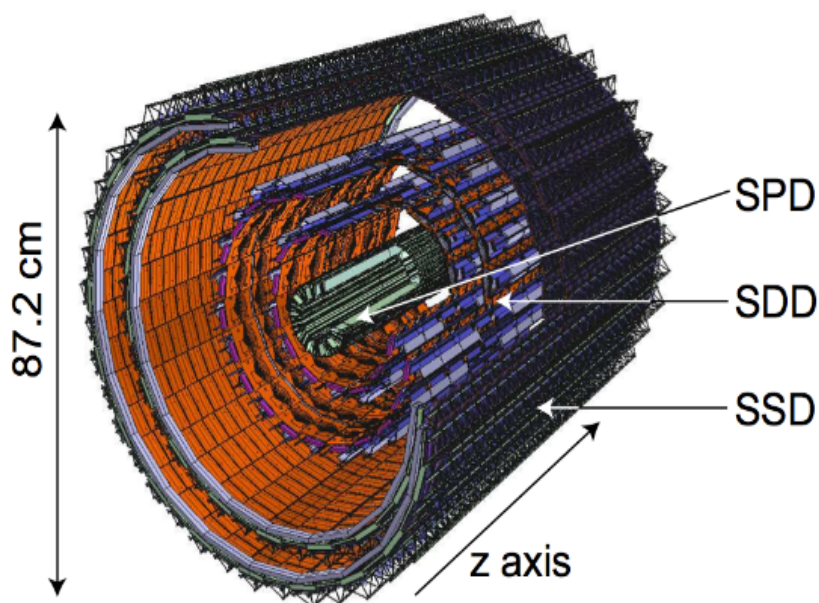


Figure 4.4: Schematic view of ITS [289]

Silicon Pixel Detector (SPD): In order to handle the high particle density close to the beam pipe (up to 50 particles per cm^2), the innermost two layers are Silicon Pixel Detectors (SPD) which have a spatial resolution of $12\ \mu\text{m}$ in the $r\phi$ direction and $100\ \mu\text{m}$ in the z direction. The SPD is made of a matrix of 2D reverse-biased silicon detector diodes connected by a conductive solder bump onto readout chips. The readout is binary meaning the output generates a pulse when the signal in a cell increases above a set threshold. The detector modules are made of half-staves that contain two ladders of silicon sensor matrices made of 256×160 cells. The whole sensor matrix has an active area of $12.8\ \text{mm}$ ($r\phi$) by $70.7\ \text{mm}$ (z). The ladders are attached and bonded to a pixel bus which carries data/control lines and power/ground planes. A multi-chip module (MCM) is bonded to the pixel bus at the end of the half-stave and controls the front-end electronics and is connected to the readout system via optical fibre links. Two half-staves are connected to a carbon-fibre support sector with an MCM at each end to form a stave. The SPD contains 60 staves in total with 240 ladders of 1200 chips and 9.8×10^6 cells. The SPD is surrounded by an aluminium coated carbon-fibre shield that prevents heat radiating to the SDD layers. The concentration of pixels in the SPD allows the primary vertex to be identified with a resolution better than $100\ \mu\text{m}$, as well as measuring secondary vertices with very good accuracy.

Silicon Drift Detector (SDD): The middle two layers are Silicon Drift Detectors (SDD) which have a spatial resolution of $35\ \mu\text{m}$ and $25\ \mu\text{m}$ in the $r\phi$ and z directions respectively. They provide good multiparticle tracking and two of the four dE/dx samples used for ITS PID. They are made from $300\ \mu\text{m}$ thick silicon with a sensitive area of $70.17\ (r\phi) \times 75.26\ (z)\ \text{mm}^2$. The sensitive area is split into two drift regions by a central cathode strip with a high bias voltage of $-2.4\ \text{kV}$ applied, which generates a drift field. At the bias voltage of $-2.4\ \text{kV}$ a drift velocity of 8.1

$\mu\text{m}/\text{ns}$ is provided. There are 89.1×10^3 cells per layer of $294 \times 202 \mu\text{m}^2$ each, which are readout by 512 electronic channels. Electrons produced by ionisation of charged particles passing through the detector, drift towards the 256 anodes at the ends of the detector. Measurement of the drift velocity allows the position of the original particle with respect to the anode to be identified. Energy loss of the particles as they interact with the detector is also measured, providing PID information.

Silicon micro-Strip Detector (SSD): The track density is reduced significantly at the outermost layers (< 1 particle per cm^2) so they are made with double-sided Silicon micro-Strip Detectors (SSD). They provide the necessary information to match tracks between the ITS and the TPC, as well as providing additional dE/dx information for low-momentum PID. The SSD has a resolution of $20(830) \mu\text{m}$ in the $r\phi(z)$ direction. It is made of 1770 individual $300 \mu\text{m}$ thick sensors with 768 sensitive strips on each side. The strips are tilted by 17.5 mrad with respect to the sensor in two different directions as shown in figure 3.6. When two strips on opposite sides detect a signal, the position of the charged particle passing through the sensor can be located to the position of the two strips crossing.

4.2.2.2 TPC

The Time Projection Chamber (TPC) [284, 285, 287, 288] (Figure 4.5) is the main tracking detector of the central barrel, optimized to provide, together with the other central barrel detectors, charged particle momentum measurements with good two-track separation, particle identification and vertex determination. It allows the complete reconstruction of tracks over a wide range of p_T from about $0.1 \text{ GeV}/c$ up to $100 \text{ GeV}/c$ for $|\eta| < 0.9$. It consists of a field cage filled with 90 m^3 of gas made up of neon, carbon dioxide and nitrogen with the ratio (90:10:5). This gas mixture is optimised for drift speed, reducing multiple scattering and stability over time. The drift velocity in the gas is highly dependent on temperature meaning a

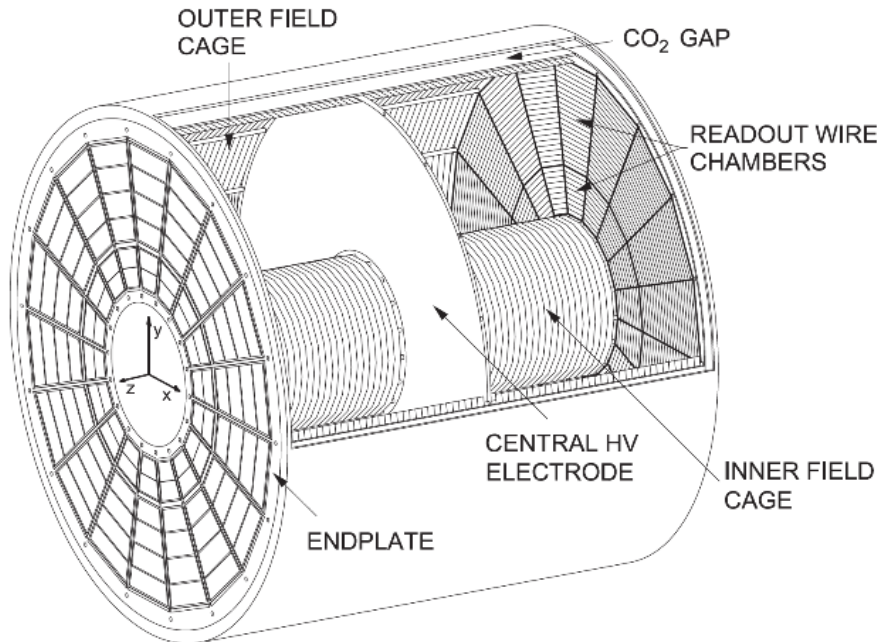


Figure 4.5: Schematic view of TPC [284]

thermal stability of $\delta T \leq 0.1$ K is required, which is provided by heat screens and cooling circuits. A uniform electric drift field of 400 V/cm is applied along the z axis by a central high-voltage electrode at 100 kV. When particles traverse the gas they ionise it, causing electrons to drift along the chamber. When the electrons reach the Multi-Wire Proportional Chambers (MWPC) at the end plates, they cause an electron avalanche which reaches the cathode pads on the readout planes. The coordinates of the pad give the $r\phi$ coordinates of the track. The time taken to cross the chamber gives the z coordinate, where a constant drift velocity is assumed. This allows a position resolution of 1100(800) μm in the $r\phi$ direction and 1250(1100) μm in the z direction to be achieved for tracks at the inner (outer) radii. The p_T of a particle can also be extracted from the curvature of its track in the 0.5 T magnetic field. A momentum resolution of 0.7% can be achieved at 1 GeV/c and up to 6.5% at around 10 GeV/c.

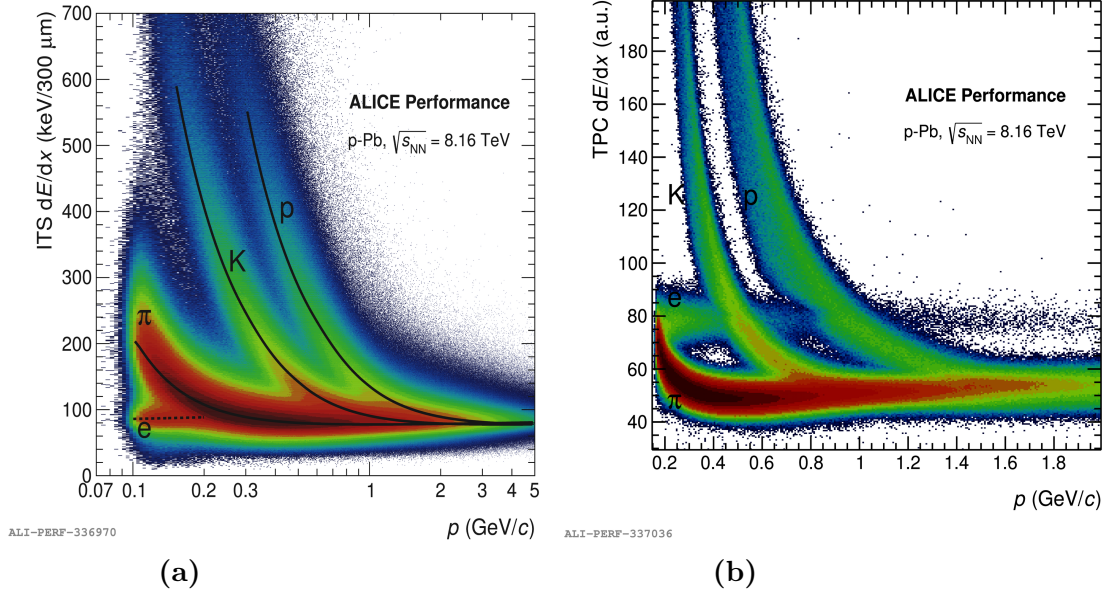


Figure 4.6: (a) Specific energy loss measured by the ITS as a function of p . (b) Specific energy loss measured by the TPC as a function of rigidity for positive tracks. The curves represent parameterizations of the expected mean energy loss

PID is also performed by the TPC by measuring the energy loss per unit length or specific energy loss, dE/dx . This is described by the Bethe-Bloch formula which shows that

$$-\frac{dE}{dx} \propto \frac{z^2}{\beta^2} \ln(1 - \beta^2)^{-1/2} \quad (4.4)$$

where Z is the atomic number of the atoms in the material being traversed and β is equal to v/c [290]. Therefore, the energy loss is dependent on the velocity of the particle. For a given momentum, particles of different masses will have a different velocity, and therefore a different dE/dx . By plotting the measured dE/dx against the measured momenta of the particles detected, particles of different masses can easily be separated for p below ~ 1 GeV, as shown in figure 4.6.

Particle Identification with the TPC [291] Many of the ALICE detector subsystems provide particle identification information, using all the PID techniques known nowadays, in their state of art implementation. A synopsis is given in Ta-

ble 4.1 and 4.2.

Table 4.1: PID techniques provided by the ALICE barrel and single arm detector subsystems.

Central barrel subsystem		Single arm subsystem	
Technique	Detector	Technique	Detector
Specific energy loss	ITS,TPC,TRD	Calorimetry	PHOS,EMCal,DCal
Transition radiation	TRD	Muon spectrometry	MUON arm
Time of Flight	TOF	Cherenkov radiation	HMPID

Table 4.2: Decay topology PID methods implemented by the ALICE barrel detector subsystems.

Decay Topology	
Particle/Method	Detector
$K^\pm$ via kinks	TPC
K_s^0, Λ, Ξ^-, D	ITS,TPC
π^0, η via photon conversion (PCM)	ITS,TPC,TOF

The identification can be performed in two different ways: by direct identification and by the reconstruction of the topology of the disintegration process of a particle. In the direct identification, usually five mass hypotheses are considered, that correspond to the particles that live long enough to be identified at track level: $e^\pm, \mu^\pm, \pi^\pm, K^\pm, p^\pm$. The specific energy loss distributions measured in LHC Run 2 are shown in Figure 4.6a for the ITS as a function of p and in Figure 4.6b for the TPC in p-Pb collisions as a function of the rigidity p/z . A second way to identify particles is via invariant mass computation for its daughter particles, as in the case, for instance, for a strong decay of a resonance and also for a weakly decayed particle. In this latter case, reconstruction can be done via topological reconstruction, since the tracks originate from a decay point that is not the primary interaction vertex. In this second case the direct identification of the daughter particles ($e^\pm, \mu^\pm, \pi^\pm, K^\pm, p^\pm$, that are the main products of these decays) has an important role in the background reduction. The strange baryons are within this second group.

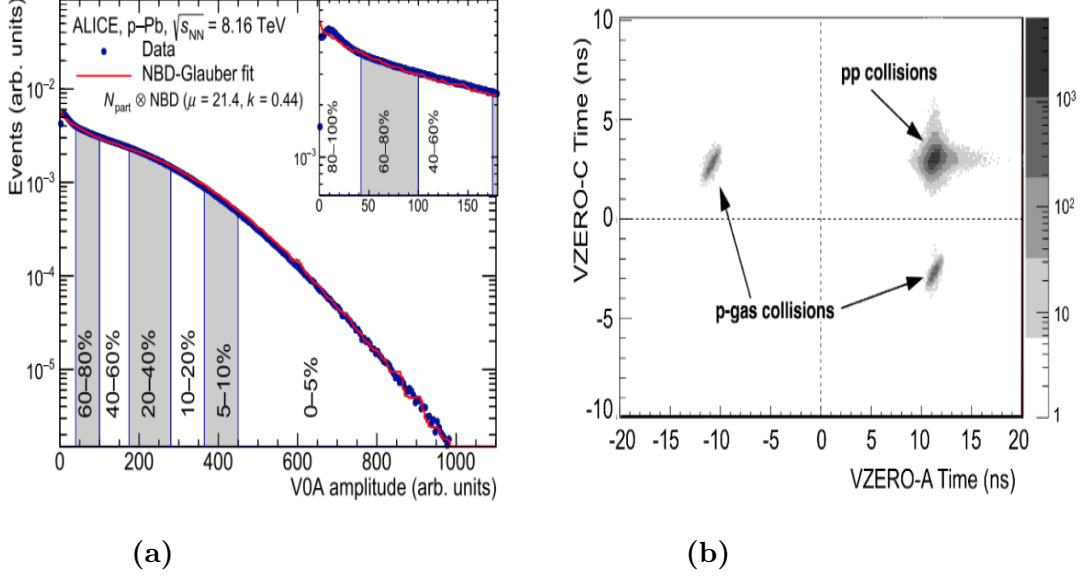


Figure 4.7: (a) Distribution of measured V0A amplitude with NBD-Glauber model fit shown in red. The multiplicity classes are indicated by the vertical lines [292]. (b) Correlation of weighted average time of flight of particles that are detected in the V0A and V0C detectors [293].

4.2.2.3 V0 Detectors

The V0 detectors [292] are plastic scintillators located on either side of the interaction point named V0A and V0C. They are used for triggering and measuring the charged-particle multiplicity for an event. The detectors are placed at $z = 340$ cm (V0A) and $z = -90$ cm (V0C) covering the pseudorapidity ranges of $2.8 < \eta < 5.1$ (V0A) and $-3.7 < \eta < 1.7$ (V0C). They consist of a disk with 32 counters in four rings, each divided into eight sectors. Particles that reach the detectors and traverse them cause the atoms of the scintillator to excite and release energy through photons. These photons are shifted down in frequency via wavelength-shifting fibres and transferred to photomultiplier tubes (PMTs) with optical fibres, where the total energy of the photons is recorded. This causes a signal which is proportional to the number of particles that have hit the detector and therefore the number of particles produced in the collision. The distribution of the amplitude of the V0 signal can be used to

calculate the multiplicity of the collision using a NBD-Glauber fit to the data [294]. For pp and Pb-Pb collisions the sum of the amplitudes in the V0A and V0C detectors is used. For p-Pb collisions the data from the Pb-going side are used, which for this analysis is the V0A detector. The data are divided up into percentile classes as shown in Figure 4.7a, for which the average charged-particle multiplicity per unit rapidity ($\langle dN_{ch}/d\eta \rangle$) can be calculated.

Measurements in the V0 detectors are also used to define the Minimum Bias (MB) trigger used in this analysis, V0AND, which corresponds to an event with at least one hit in the scintillators of both the V0A and the V0C detectors [293]. The V0 detectors are also important for rejection of events caused by beam-gas interactions i.e. interactions of the beam with residual gas found in the beam pipe. This is done by measuring the time of flight of particles by each of the V0 detector arrays. A beam-gas interaction can clearly be distinguished from beam-beam events by looking at the correlation between the timing signals in the two detectors as shown in Figure 4.7b.

4.2.3 Event Reconstruction

Using the raw data obtained from the detectors, which includes information on the activated detector elements and timing information, the track trajectories of particles need to be reconstructed for analysis. This allows useful variables to be extracted such as the trajectory coordinates, particle momentum, particle energy loss and distance of closest approach (DCA) to the primary vertex (PV). When particles traverse the detectors they leave behind signals in adjacent detector cells. These signal clusters give the position information of the tracks. Using cluster finding reduces the effect of random noise in position and timing measurements. The Kalman filter technique is used on these clusters to reconstruct the tracks [295]. First a ‘seed’ cluster is found in the outer TPC and then tracked back through the inner TPC and the ITS, where the

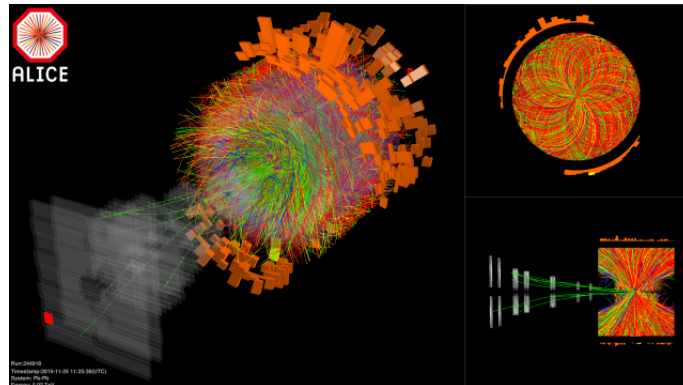


Figure 4.8: Event display for Pb–Pb collision recorded by ALICE during one of the first runs at centre of mass energy of 5.02 TeV per nucleon.

location of the PV can be estimated. Information from other detectors, if available, is then added to produce global tracks. Secondary vertices are located by combining tracks sufficiently far from the PV and utilising the topology of the decays of unstable particles [296]. After reconstruction, the useful information can be stored in compact Event Summary Data (ESD) files, which the analysis in this thesis is performed on.

Events visualisation in ALICE is a very important aspect for both experts and general public. It provides early feedback on the data quality and reconstruction correctness, it is an important tool for visual inspection by experts and it produces publicity pictures that are widely used for the outreach. The aim of all event visualisation software in high energy physics experiments is to display the data collected by detectors in a human-understandable way. An illustration of ALICE event display is shown in Figure 4.8.

4.3 LHC Run 3 Upgrades

In November 2018, the LHC finished its second running period (Run 2), which started in September 2015. In this period, a peak interaction rate of 7.5 kHz was reached for Pb–Pb collisions with centre-of-mass energy of 6.3 TeV per nucleon pair [297]. As a comparison, in the heavy ion run in 2011 the interaction rate reached was 3 kHz [298]. There was a two year planned shutdown to prepare for Run 3 starting

in 2021, although started in 2022. Many of the planned measurements for Run 3 will involve complex probes at low transverse momentum, where traditional methods of triggering will not be applicable. Therefore, the ALICE collaboration planned to upgrade the current detector by enhancing its low-momentum vertexing and tracking capability, and allowing data taking at substantially higher interaction rates. Run 3 has an expected peak luminosity of $6 \times 10^{27} \text{cm}^{-2} \text{s}^{-1}$ and interaction rates of about 50 kHz for Pb–Pb collisions [299]. As per the proposed plan, the ALICE detector is modified such that all interactions will be inspected (trigger less mode). This implies a major upgrade of the TPC detector and a modification of the readout electronics of other detectors to comply with the high readout rate.

The envisaged upgrades preserve the excellent particle identification capability while enhancing the vertexing and tracking capabilities, and include:

- ❑ A new, high-resolution, low-material-thickness Inner Tracking System (ITS).
With this new detector the resolution of the distance of closest approach between a track and the primary vertex will be improved by a factor of about 3, and the standalone ITS tracking performance will be significantly enhanced.
- ❑ Upgrade of the Time-Projection Chamber (TPC) with replacement of the readout multi-wire chambers with GEM (Gas Electron Multiplier) detectors and new pipelined readout electronics. This is needed in order to operate the TPC in ungated mode and make its readout dead-time free.
- ❑ Upgrade of the readout electronics of: Transition-Radiation Detector (TRD), Time-Of-Flight (TOF) detector, PHOTON Spectrometer (PHOS) and Muon spectrometer. This upgrade will allow for high rate data taking with these detectors.
- ❑ Upgrade of the forward trigger detectors and trigger system for high rate oper-

ation.

- Upgrade of the Online Systems: High-Level Trigger (HLT), data acquisition (DAQ) and trigger system, to adapt for high rates. The rate for heavy-ion events handled by the online systems up to permanent data storage should be increased up to 50 kHz corresponding to roughly two orders of magnitude with respect to the present system.
- Upgrade of the offline data processing software to cope with the reconstruction and analysis of a much larger number of events. The strategy for increasing the offline event processing rate up to two order of magnitude

4.3.1 Motivation for ITS Upgrade

A crucial limitation of the present ITS detector is given by its limited read-out rate capabilities. The ITS can run at a maximum rate of 1 kHz (with dead time close to 100%), irrespective of the detector occupancy. For all physics channels that cannot be selected by a trigger, this rate limitation restricts ALICE to use only a small fraction of the full Pb–Pb collision rate of 8 kHz that the LHC presently can deliver and prevents the collection of required reference data in pp collisions. Clearly the present ITS is inadequate to fulfil the required rate capabilities.

The primary focus of the ITS upgrade is on improving the performance for detection of heavy-flavour hadrons, and of the thermal photons and low-mass dileptons emitted by the QGP. The new ITS will greatly improve the distance of closest approach to the primary vertex, tracking efficiency at low transverse momenta, and read-out capabilities. The baseline solution for the layout of the ITS upgrade is to replace the existing ITS detector in its entirety with seven concentric layers of pixel detectors as shown in Figure 4.9. Monolithic Active Pixel Sensors (MAPS) implemented using the 0.18 μm CMOS technology of TowerJazz have been selected as the

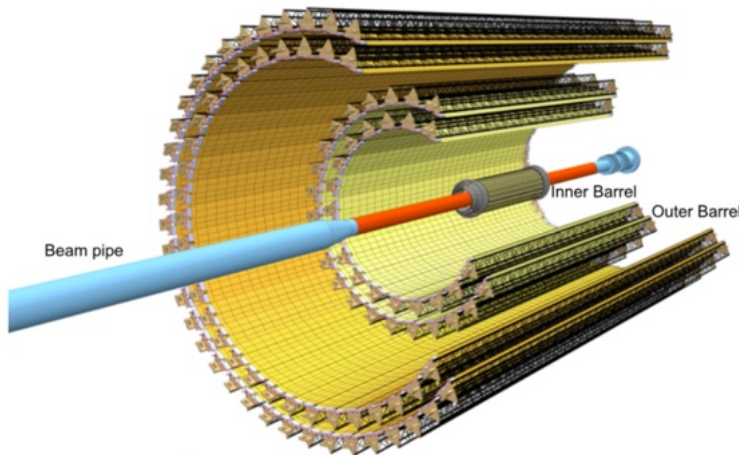


Figure 4.9: Layout of the new ITS detector.

technology for all layers. The basic MAPS element is the Pixel Chip. It consists of a single silicon die of about $15 \text{ mm} \times 30 \text{ mm}$, which incorporates a high-resistivity silicon epitaxial layer (sensor active volume), a matrix of charge collection diodes (pixels) with a pitch of the order of $30 \mu\text{m}$, and the electronics that perform signal amplification, digitisation and zero-suppression. Only the information on whether or not a particle was crossing a pixel is read out.

A new silicon tracker will enable the track position resolution at the primary vertex to be improved by a factor of three or greater. The standalone tracking efficiency would be comparable to what can be presently achieved by combining the information of the ITS and the TPC, but extended to much lower values of transverse momentum. The relative momentum resolution of the silicon tracker standalone would be about 4% up to $2 \text{ GeV}/c$ and remain below 6% up to $20 \text{ GeV}/c$.

4.3.2 Motivation for TPC Upgrade

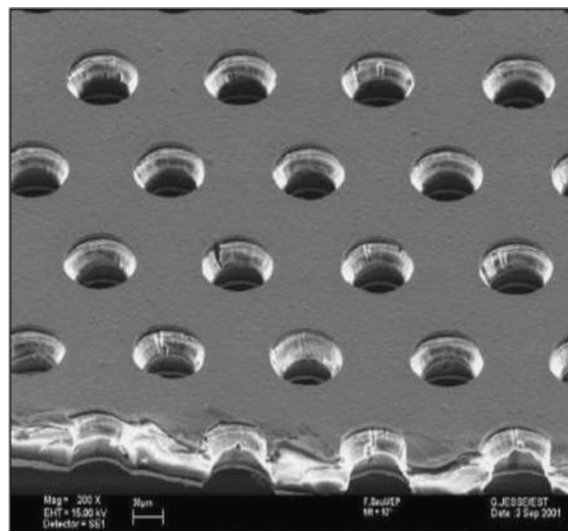
The drift time of electrons, from centre of the detector to the end plates, in the currently used Ne-CO₂ (90-10) gas in the TPC is $\sim 100 \mu\text{s}$. The drift time of positive ions from the amplification region around the MWPC anode wires back to the central electrode is $\sim 180 \mu\text{s}$. To avoid any ions drifting back from the amplification region

to the drift region, a gating grid is in place, which is enabled after the initial electron drift time, preventing any back drift, but also preventing electrons to pass into the amplification region. This leads to an intrinsic dead time of $\sim 280 \mu\text{s}$ and a limitation in the maximum interaction rate of 3.5 kHz. If the gating grid structure were to be permanently disabled and the TPC ran at the targeted interaction rate of 50 kHz ($20 \mu\text{s}$) for Run 3, then space charge distortions would occur due to the accumulated ions in the drift region, which would render track reconstructions useless. Operation of the TPC at 50 kHz can thus not be accomplished with the current gating scheme [287]. It will therefore be replaced by a multi-stage Gas Electron Multiplier (GEM) system [300]. GEMs have been proven to operate reliably in high-rate applications and provide intrinsic ion blocking capabilities, therefore enabling the TPC to operate in a continuous, ungated readout mode where the readout is dead-time free [301].

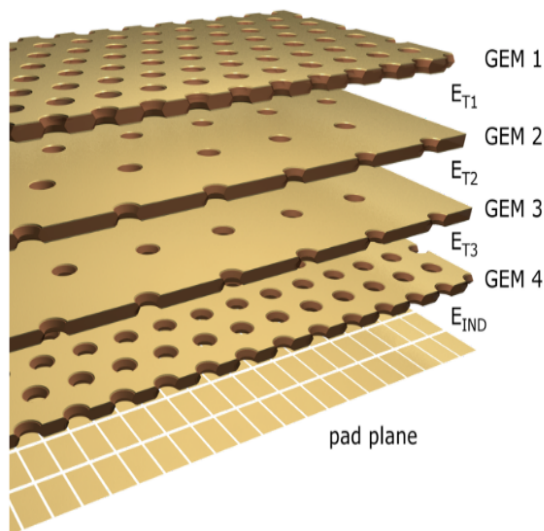
As the current readout electronics is based around a triggered readout, a full redesign of the complete front-end chain is needed [302]. The new electronics must implement a continuous readout scheme and should be able to handle the resulting higher readout data rate. In addition, it should accommodate both the negative signal polarity of the new GEM detectors and the lower gas gain, which demands a low noise design.

4.3.2.1 Principal Operation of GEMs

The GEM detector is constructed of a thin foil of Kapton, clad in copper on both sides and chemically pierced with small holes a fraction of a millimetre apart, see Figure 4.10a. A high voltage potential is set up between the two sides of the foil, making large electric fields in the holes. Electrons in the gas entering the hole will create an avalanche of hundreds of electrons. By having several layers of GEMs the number of electrons can be increased even more by providing additional stages of amplification. By offsetting the holes between the layers, the ion backflow can be



(a)



(b)

Figure 4.10: (a) Close-up of GEM construction. (b)CTPC GEM stackup

decreased significantly [300]. The electrons are collected by a suitable device, like a pad plane. A stackup as to be used in the TPC detector can be seen in Figure 4.10b.

4.3.3 Physics Motivations for Run 3

After the ALICE upgrade, the excellent tracking performance in high-multiplicity environment and particle identification over a large momentum range will help us to focus on the below mentioned physics questions:

- ❑ Study of the thermalization of partons in the QGP, with focus on the massive charm and beauty quarks. Heavy-quark azimuthal-flow anisotropy is especially sensitive to the partonic equation of state. Ultimately, heavy quarks might fully equilibrate and become part of the strongly-coupled medium.
- ❑ Study of the low-momentum quarkonium dissociation and, possibly, regeneration pattern, as a probe of deconfinement, and an evaluation of the medium temperature.
- ❑ Study of the production of thermal photons and low-mass dileptons emitted by the QGP. This should allow to assess the initial temperature and the equation of state of the medium, as well as to shed light on the chiral nature of the phase transition.
- ❑ Study of the in-medium parton energy-loss mechanism that provides both a testing ground for the multi-particle aspects of QCD and a probe of the QGP density. The relevant observables are: jet structure, jet–jet and photon–jet correlations, and jet correlations with high-momentum identified hadrons and heavy-flavour particle production in jets. In particular, it is crucial to characterize the dependencies of energy loss on the parton colour-charge, mass, and energy, as well as on the density of the medium.

- Search for heavy nuclear states such as light multi- Λ hyper-nuclei ${}^5_{\Lambda\Lambda}\text{H}$, bound states of $(\Lambda\Lambda)$ or the H dibaryon, (Λn) bound state, as well as bound states involving multi-strange baryons; a systematic study of light nuclei and anti-nuclei production.

Chapter 5

This chapter includes the techniques adopted to identify strange hadrons in the ALICE. The full description of the step-by-step procedure adopted to analyse the strange hadrons including several corrections and systematics computations will be discussed in this chapter.

Chapter 5

Strange hadrons identification at ALICE

5.1 Introduction to strange hadrons

5.1.1 K^0 meson and its anti-particle

The K^0 meson is a pseudoscalar state consisting of a d quark and an $\bar{s}$ antiquark:

$$|K^0\rangle = |d\bar{s}\rangle \quad (5.1)$$

$$|\overline{K^0}\rangle = |\bar{d}s\rangle \quad (5.2)$$

If we define as “strangeness” operator, S , (which counts strange quarks), these states are eigenstates with:

$$S|K^0\rangle = |K^0\rangle \quad (5.3)$$

$$S|\overline{K^0}\rangle = -|\overline{K^0}\rangle \quad (5.4)$$

These are not eigenstates of C , the charge conjugation operator (which changes particles to antiparticles). It is convenient to pick the antiparticle phases such that:

$$C|K^0\rangle = -|\bar{K}^0\rangle \quad (5.5)$$

$$C|\bar{K}^0\rangle = -|K^0\rangle \quad (5.6)$$

If we multiply the C operator by the parity operator P , we have:

$$CP|K^0\rangle = |\bar{K}^0\rangle \quad (5.7)$$

$$CP|\bar{K}^0\rangle = |K^0\rangle \quad (5.8)$$

Thus the neutral kaons are not eigenstates of CP i.e. $CP|K^0\rangle \neq \pm|K^0\rangle$. The eigenstates of CP are (the choice of the nomenclature shall be discussed later):

$$|K_S^0\rangle = \frac{1}{\sqrt{2}}(|K^0\rangle + |\bar{K}^0\rangle), \text{ with } CP = +1 \quad (5.9)$$

$$|K_L^0\rangle = \frac{1}{\sqrt{2}}(|K^0\rangle - |\bar{K}^0\rangle), \text{ with } CP = -1 \quad (5.10)$$

where $\frac{1}{\sqrt{2}}$ ensures that $\langle K_S^0|K_S^0\rangle = 1$ and $\langle K_L^0|K_L^0\rangle = 1$.

We remark that the K^0 (or $\bar{K}^0$) is the lowest mass particle containing the strange quark. Thus, the only permitted decays must be via the weak interaction. To a good approximation (but not exactly!), CP is conserved in the weak interaction (and even more so in the strong and electromagnetic interactions); we shall assume this here.

A neutral K meson (K^0 or $\bar{K}^0$) is observed to decay sometimes to two pions and sometimes to three pions. For example, consider the observed process $K^0 \rightarrow \pi^0\pi^0$. Since all of the particles in this decay are spinless, the decay must proceed with zero orbital angular momentum (“S-wave” decay). Thus, the parity of the $\pi^0\pi^0$ system in

the final state must be positive. But we said that the K^0 is a pseudoscalar particle, i.e., has negative parity. Thus, this is a parity-violating decay. The weak interaction is known to violate parity (i.e., parity is not conserved in the weak interaction), so this is all right. The π^0 is its own anti-particle, hence the $|\pi^0\pi^0\rangle$ final state is an eigenstate of C with eigenvalue $+1$. Thus, the $|\pi^0\pi^0\rangle$ final final state is also an eigenstate of CP with eigenvalue $+1$. Under our approximation that CP is conserved in the weak interaction, we therefore conclude that the observation of a $K^0 \rightarrow \pi^0\pi^0$ decay projects out the K_S^0 component of the K^0 meson (likewise for the $\bar{K}^0$). The 2π decay mode is favored by phase space over decays to greater numbers of pions. However, the $K_L^0 \rightarrow 2\pi$ decay is forbidden by CP conservation. Hence, the $K_L^0 \rightarrow 3\pi$ decay is important for K_L^0 . Because the phase space is considerably suppressed, the K_L^0 decay rate is much slower than the K_S^0 rate. The observed lifetimes of the K_S^0 and K_L^0 are, respectively:

$$\tau = 9 \times 10^{-11} s \text{ (90 ps)} \quad (5.11)$$

$$\tau = 5 \times 10^{-8} s \text{ (50 ns)} \quad (5.12)$$

CP violation by neutral kaons The kaons provide a testing ground for CP invariance. By using a long enough drift distance for an initial beam of K^0 particles (i.e. long enough drift time), an arbitrarily pure beam of the long-lived kaons (K_L^0) can be produced. Then if any decays to $2\pi^0$ are observed (which can only come from a K_L^0 decay) this is an indication of CP violation. This was observed in 1964 [303]. It is a small effect but the result indicates that the long lived kaons are not purely K_L^0 but must have a small admixture of K_S^0 . Since then, additional examples of CP violation have been found. CP violation provides evidence of the unequal treatment of matter and antimatter and may be the reason why matter is dominant over anti-matter in the universe.

5.1.2 K_S^0 meson and Λ ($\bar{\Lambda}$) baryon

In this thesis we will be focusing on K_S^0 , Λ and $\bar{\Lambda}$.

K_S^0 is a meson and has a mass of 497.614 ± 0.024 MeV and $c\tau = 2.6844$ cm. K_S^0 is measured in ALICE in the $\pi^+\pi^-$ channel which has a branching ratio of $(69.20 \pm 0.05)\%$ [304].

The Λ baryon, also called Λ hyperon, has a mass of (1115.683 ± 0.006) MeV and a $c\tau = 7.89$ cm. Its quark composition is: uds . Λ is measured in ALICE in the $p\pi^-$ channel which has a branching ratio of $(63.9 \pm 0.5\%)$ [304].

The $\bar{\Lambda}$ baryon, is the anti-particle of Λ , has a mass of (1115.683 ± 0.006) MeV and a $c\tau = 7.89$ cm. Its quark composition is: $\bar{u}\bar{d}\bar{s}$. $\bar{\Lambda}$ is measured in ALICE in the $\bar{p}\pi^+$ channel which has a branching ratio of $(63.9 \pm 0.5\%)$ [304].

5.1.3 V^0 identification

The reconstruction of K_S^0 and Λ ($\bar{\Lambda}$) makes use of the fact that the both the particles have a special decay topology, the so called V^0 topology. Since they are neutral and live relatively long they produce no visible track in the tracking detectors until they decay and lead into two visible tracks of opposite signed charged tracks. These two tracks form the shape of a V, starting at the decay vertex of these particles, shown in Figure 5.1. V^0 s have been discovered in 1947 by G.D. Rochester and C.C. Butler [305] in a cloud-chamber by observing exactly this special topology. This was the first observation of neutral kaon decays. The Λ was first observed in 1950 by V. D. Hopper and S. Biswas in photographic emulsions [306]. For the detection of these V^0 s two approaches are followed in ALICE. Either the entire event is checked offline for tracks which show a V^0 topology or already while the tracks are reconstructed, the tracking algorithm searches for V^0 candidates. These are the so called Offline and the On-the-Fly V^0 finders respectively.

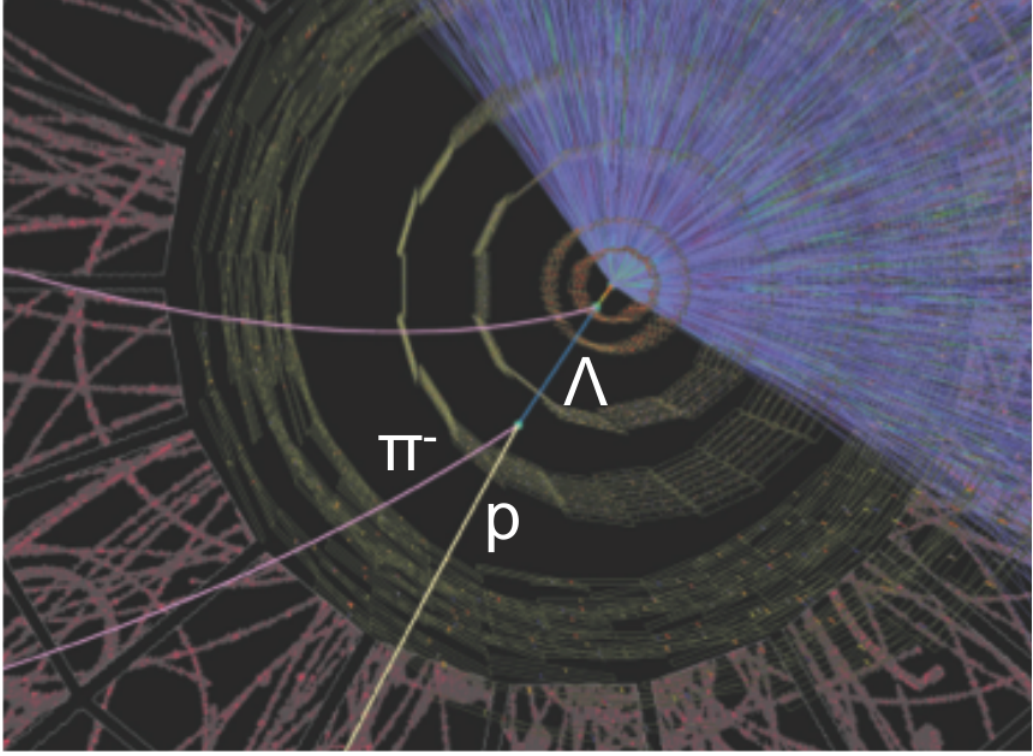


Figure 5.1: Simulated Λ decay in ALICE

5.1.3.1 Topological Reconstruction

The technique which is used for V^0 's identification is known as topological reconstruction because it involves the particle association to get a V^0 based on geometrical and kinematical criteria. The daughter tracks have to fulfill a set of geometrical and kinematical conditions (selection criteria). The cuts considered are discussed below, with several shown graphically in Figure 5.2

The V^0 finding proceeds combining each secondary track with all the other secondary tracks having an opposite charge. Two distinct cuts are applied on the minimum value of the impact parameters to the primary vertex for the two opposite charged tracks, to ensure that these are secondary particles. Such pair of tracks is also rejected if the distance of closest approach (DCA) in space between the two tracks is larger than a given value. Once the V^0 vertex position is defined, only the vertices inside a given fiducial volume are kept; the inner boundary of this fiducial

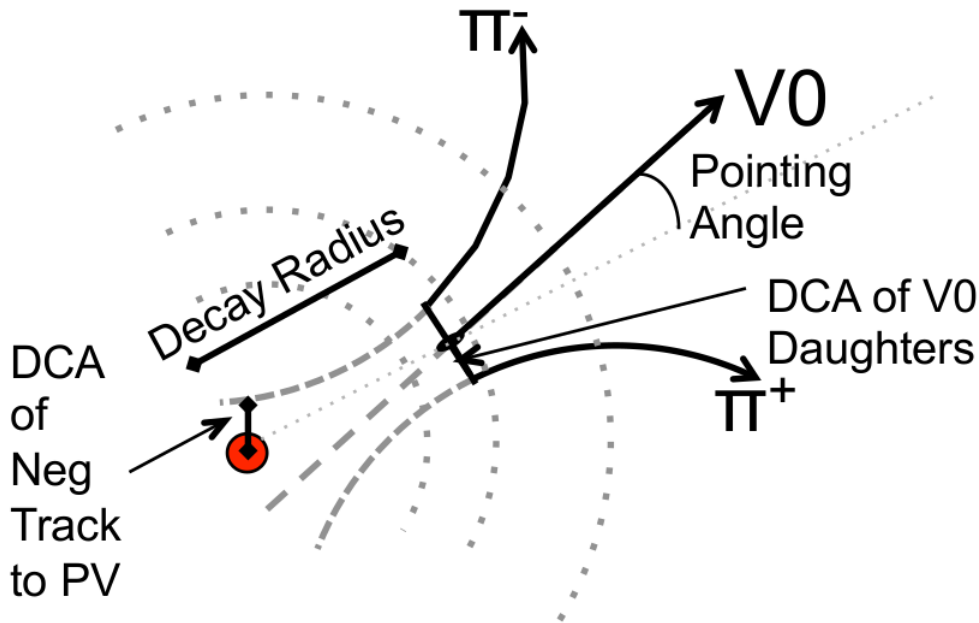


Figure 5.2: Kinematic Cuts - Schematic indicating some kinematic cut variables.

area is limited in order to keep track density and number of miscombinations low. Finally, the V^0 finding procedure checks whether the momentum vector of the V^0 candidate does not point back to the primary vertex. This is performed applying a cut on the cosine of the angle between the V^0 reconstructed momentum vector and a vector connecting the primary vertex to the secondary vertex position (pointing angle)

5.1.3.2 Invariant Mass Calculation

The squared mass value for each single candidate, in the case of a true V^0 particle (suppose Λ), from the kinematic equation, is given by:

$$M_{theor}^2(\Lambda) = (E_{\pi^-} + E_p)^2 - (\vec{p}_{\pi^-} + \vec{p}_p)^2 \quad (5.13)$$

A similar expression gives the squared mass in the hypothesis of the $\bar{\Lambda}$, where the tracks are a π^+ and an $\bar{p}$. Experimentally one has to measure the invariant mass

starting from the measurement of the two daughter tracks. This can be done using the following quantity:

$$M_{experimental}^2(\Lambda) = \left(\sqrt{m_{\pi^-}^2 + \vec{p}_{neg}^2} + \sqrt{m_p^2 + \vec{p}_{pos}^2} \right)^2 + (\vec{p}_{neg} + \vec{p}_{pos})^2 \quad (5.14)$$

Let us suppose we are identifying Λ but since the candidates are built for findable

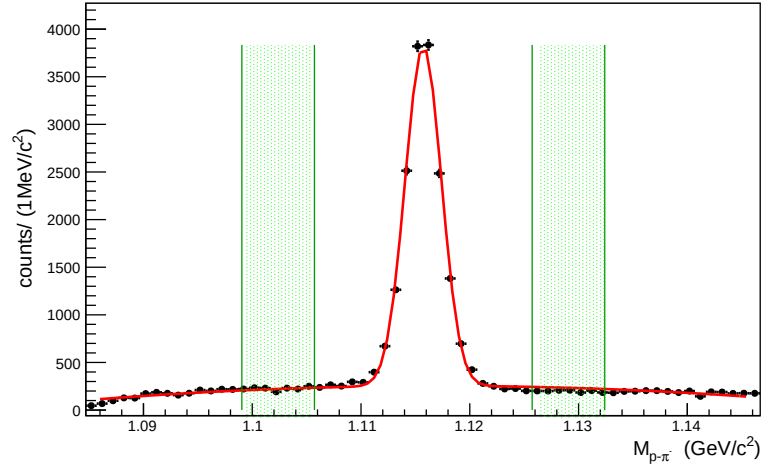


Figure 5.3: Λ Candidate's invariant mass distribution

V^0 decays they do not just consist of true Λ , but also contain true K_S^0 and γ (which have been converted into $e^+ e^-$ pairs) decays. They give a significant contribution to the invariant mass of the Λ which is shown in Figure 5.3. To remove these so-called reflections several ways are possible:

- tuning of the topological cuts,
- usage of particle identification,

The topological cuts are chosen so that they are removing most background but keeping at the same time most of the signal. The particle identification, utilizing a $n\sigma$ cut on the dE/dx in the Time-Projection Chamber, helps a lot to remove fake

Λ 's. But since electrons cross pions at momenta of around 100 MeV/c and protons at around 1.0 GeV/c, and at higher momenta also protons and pions are becoming harder to disentangle, there are still fakes left as shown in previous chapter in Figure 4.6.

5.2 Data Sample and Event selections

The analyses described in this thesis is conducted on an event sample collected at the end of 2016 (Run 2 data) on proton–lead collisions at a centre of mass energy of $\sqrt{s_{NN}} = 8.16$ TeV. The data set is referred to internally by ALICE as LHC16r_CENT, where CENT refers to the CENT cluster which includes all events with the slowest detector in ALICE, the SDD, in the readout (there is also a FAST cluster which includes events collected without the SDD when it is busy). Six runs with the lowest interaction rates were used.

Basic event selections used in p–Pb collisions in ALICE were applied in this analysis. The data selected were collected using a minimum bias trigger, which requires at least one hit in both the V0A and V0C detectors. Also, a selection on the position of the PV of the interaction was made, requiring it to be within ± 10 cm from the centre of the detector along the z axis. Finally, an initial cut was placed to remove in-bunch pileup events using the SPD. Events were rejected if a secondary interaction vertex was found, separated from the first one by more than 0.8 cm along the beam (z) axis.

5.2.1 Event Multiplicity Selections

The multiplicity estimator used for this analysis is an estimator based on the VZERO-A amplitude, which is proportional to charged particle counts in the pseudorapidity range $2.8 < \eta < 5.1$ - the Pb-going direction for this period. The amplitudes are classified in percentiles of the selected events, retrieved via a call to the function:

```
AliPPVsMultUtils::GetMultiplicityPercentile( AliVEvent, "V0A" );
```

The actual multiplicity classes used for this analysis are:

- K_S^0 mult. bins = $\{0, 1, 5, 10, 15, 20, 30, 40, 50, 60, 70, 80, 100\}$ %
- Λ ($\bar{\Lambda}$) mult. bins = $\{0, 1, 5, 10, 15, 20, 30, 40, 50, 60, 70, 80, 100\}$ %

5.3 Kinematic Cuts: Topological and Non-topological

The analysis was performed using tasks, written in C++ and compatible with the Aliroot framework. V^0 candidates are found by combining charged tracks reconstructed in the ITS and TPC volume and imposing topological and kinematic restrictions to select them. The cuts can be categorized as topological and non-topological and are described below:

- **Rapidity Interval** The rapidity interval chosen for analysis is $-0.5 < y < 0$ for both data and Monte Carlo corresponding to the central rapidity region for p-Pb collisions.
- **Pseudorapidity (η)**: A cut is applied on the pseudorapidity in order to include tracks only within full the TPC acceptance. Any tracks outside the pseudorapidity range $|\eta| < 0.8$ might risk not getting fully projected in the TPC, leading to reduced detector efficiency.
- **TPC dE/dx Selection** Additionally, in order to remove background from combinations of tracks that are of the wrong particle species, we further apply a selection on the energy loss while traversing the TPC gas. A candidate is removed from the analysis if any of its two daughter tracks has a dE/dx value that differs by more than 4σ from the expected value.
- **Proper Lifetime** In order to minimize secondary V^0 s in the analysis, a selection on the proper lifetime of the V^0 candidate is applied. The proper lifetime

is calculated via mL/p , where m is the expected V^0 mass under the current hypothesis, L is the linear (3D) distance between the primary vertex and the V^0 decay vertex, and p is the total momentum of the candidate.

- **Primary Selection and MC Association** For the analysis of simulated data, the reconstructed daughter tracks are required to match with particles in the MC record that are of the correct desired species. Furthermore, their mother particle is required to be of the desired type as well. Finally, only primary K_S^0 , Λ or $\bar{\Lambda}$ are accepted using `PrimaryStatus` flag = 1 for analysis.
- **Daughter Track Quality Selections** Finally, in order to ensure that only high-quality V^0 s are included in the analysis, a set of track quality requirements is applied as follows:
 - **Number of TPC clusters:** A minimum number of TPC clusters is imposed on each track in order to select good quality tracks.
 - **The no. of crossed rows over findable:** In order to further reduce dependence on the running condition of the TPC, a cut was proposed on the ratio of the number of TPC pad rows actually crossed by the track to the number of clusters which could potentially be found. A higher value of this ratio would indicate a more reliable track.
- **Competing species rejection window** An additional invariant mass selection cut is applied to the mass of the V^0 , which removes all the candidates for which the mass lies outside a window of $\pm 5 \text{ MeV}/c$ and $\pm 10 \text{ MeV}/c$ around the PDG mass of K_S^0 and Λ respectively.

Furthermore, an additional set of geometrical requirements is applied to ensure that the tracks are configured in space in a manner that is consistent with the expected V^0 decay topology. These requirements will henceforth be referred to

as ‘topological selections’ and are shown in the top part of table 5.1, while the bottom part summarizes the other aforementioned selections.

- **Decay Radius:** In order to be able to separate the secondary vertex from the primary vertex a cut is added to the decay radius; which ensures that all V^0 candidates within this cut value from the primary vertex are excluded from the analysis.
- **DCAPos(Neg)toPV:** A cut is also added to the DCA-Distance of Closest Approach between the daughter particles and the primary vertex. Putting a minimal limit on the DCA will ensure that the daughter particle tracks are not mixed up with any particle coming from the primary vertex.
- **DCAV0Daughters:** To assure that the daughter particles both come from the same secondary vertex, an upper limit is applied to the DCA of the daughter particles. As the detector has a finite track resolution, it is impossible to say definitively that the two tracks originated from a common point, but the smaller the measured distance of closest approach between the two tracks, the greater the likelihood that they did in fact originate from a common decay vertex. It is measured in units of σ , where σ is the position resolution of the tracks at the point of closest approach.
- **Cosine of Pointing Angle:** The pointing angle (P.A.) is defined as the angle between the line connecting the primary to the decay vertex and the reconstructed momentum vector of the decaying particle. This angle increases with the bending of the decaying particle. Since the V^0 s are neutral, they do not bend in the magnetic field and follow a straight trajectory, they do not contribute to a larger pointing angle.

5.3.0.1 Optimization of cuts

All the cuts are fine-tuned in the final analysis in order to find the best compromise between the combinatorial background minimization and the significance of the signals. The criteria adopted to determine the values of the topological selections correspond to a compromise between the best signal over background ratio and the minimization of the signal loss. These values have been chosen and validated via a quality assurance study (cut study). We use three broad p_T bins for K_S^0 , Λ and $\bar{\Lambda}$ which are shown below:

- $K_S^0 = \{(0.2-1), (1-3), (3-10)\}$ (GeV/ c)
- Λ (or $\bar{\Lambda}$) = $\{(0.6-1), (1-3.2), (3.2-10)\}$ (GeV/ c)

Below are the steps followed to choose the cuts:

- Choose the value of the variable (cut) such that it has the highest value of the significance
- Check if at that value signal loss fraction is $<10\%$, if not then choose the next value which has the second highest value of the significance and so on
- After following the above two steps, also check that MC and real data agree within a couple of percentage.

From the Figure 5.12 it is clear that in case of many variables (cuts) p_T dependence has been observed. Table 5.1 shows the values that are chosen after this cut study (optimization).

5.4 Signal Extraction

The signal is extracted directly from the invariant mass distributions, in different p_T and multiplicity bins. Figure 5.3 demonstrates how mass histograms are divided

into side-bands in the process, in a particular p_T bin for the 0-100 % multiplicity class. Both K_S^0 and Λ invariant mass distributions are fitted with a Gaussian function in the peak region. A central peak band is defined in terms of the number of σ from the Gaussian peak, in this case chosen to be within -4σ and $+4\sigma$ (i.e. inside the $[-4; +4]\sigma$ window).

Background Subtraction: The background is the result of contributions from, e.g., wrongly associated tracks that do not originate from a V^0 decay or particles of similar mass that have decayed with a V^0 -like topology. To be able to work with a clean signal, it is necessary to subtract the background from the mass spectrum. The adjacent side-bands to the mass peak in the ranges $[-10; -6]\sigma$ and $[+6; +10]\sigma$, represent the background. This choice was made to ensure that in every p_T interval of every spectrum, the side-bands do not overlap with the mass peak and be fully defined in the background region. The background subtraction method applied is a so-called bin counting approach, which consists of directly subtracting the bin entries inside the two side-bands from the total number of entries registered in the peak region. Such a method forcibly assumes that the background area under the mass peak is the same as it is in the outside regions where no signal is present. In other words, it assumes a linear background. Although the fit was not directly used in the default signal extraction method, an approach in which the integral under that linear function represents the background was pursued in the studies of the systematic uncertainties.

5.4.1 Statistical uncertainty

The subtraction applied in the signal extraction method is performed on two non-correlated quantities: the peak counts, which contain the signal (N_s) and the background (N_b) and the sum of the entries in the adjacent background bands which also constitute the background (N_b). A Poisson uncertainty is attributed to these

three counts. The statistical uncertainty on the signal is obtained through the error propagation of this subtraction and is:

$$\sigma_{stat} = \sqrt{N_s + 2N_b} \quad (5.15)$$

5.5 Signals Corrections

Using MC data, where the all properties of a particle is known, correction factors for inaccuracies such as feeddown and detector acceptance and efficiency can be calculated. Samples of simulated p–Pb collisions were used for efficiency calculations. The simulated data were created using the EPOS-LHC [307] and DPMJET [308] event generators and then propagated through the ALICE detector using GEANT4 [309].

5.5.1 Feed-down Correction

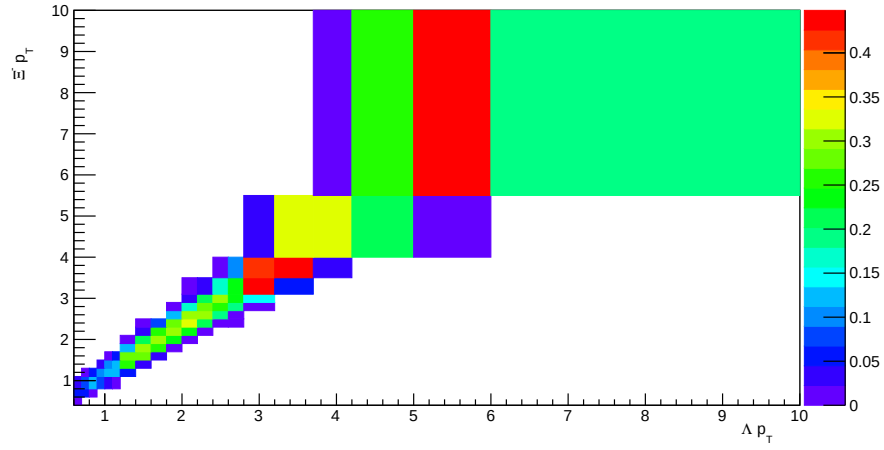
The feed-down contribution to the Λ and $\bar{\Lambda}$ spectra coming from both charged (Ξ^- , Ξ^+) and neutral (Ξ^0) decays is computed by using the same method adopted in previous strangeness analyses [310–312]. The spectrum of candidates from feed-down is calculated as follows:

$$\Lambda_{primary}^{raw}(p_{T_i}) = \Lambda_{measured}^{raw}(p_{T_i}) - \sum_j F_{ij} \int_{p_{T_j}} \frac{dN}{dp_T}(\Xi^-) \quad (5.16)$$

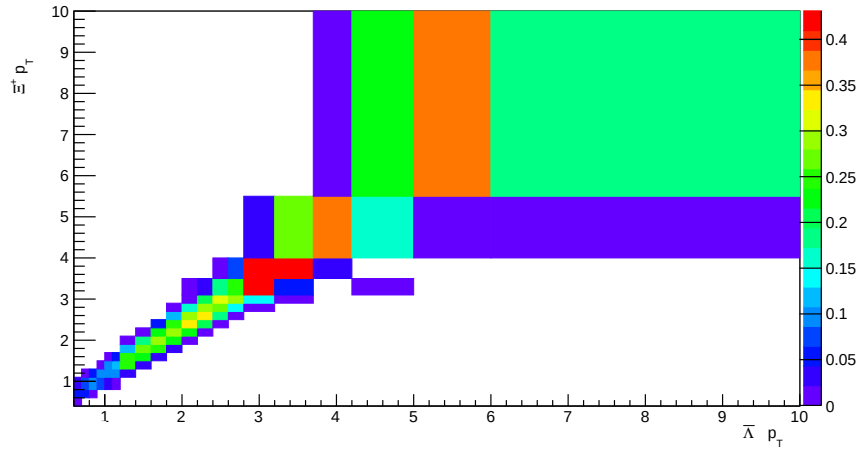
where $\int_{p_{T_j}} \frac{dN}{dp_T}(\Xi^-)$ is the Levy-Tsallis fit (described in Appendix A) to the extrapolated Ξ spectra in our case. F_{ij} is the feed-down matrix defined as

$$F_{ij} = \frac{N_{reco}(\Lambda)_{from \Xi \text{ bin } j}^{in \text{ bin } i}}{N_{gen}(\Xi)_{\Xi \text{ bin } j}} \quad (5.17)$$

The feed-down matrix, F_{ij} , is calculated using MC productions and represents the fraction of reconstructed Λ (or $\bar{\Lambda}$) produced in the p_T bin- i coming from the decay of a Ξ particle (charged or neutral) produced in the p_T bin- j . It is shown in Figure 5.4.

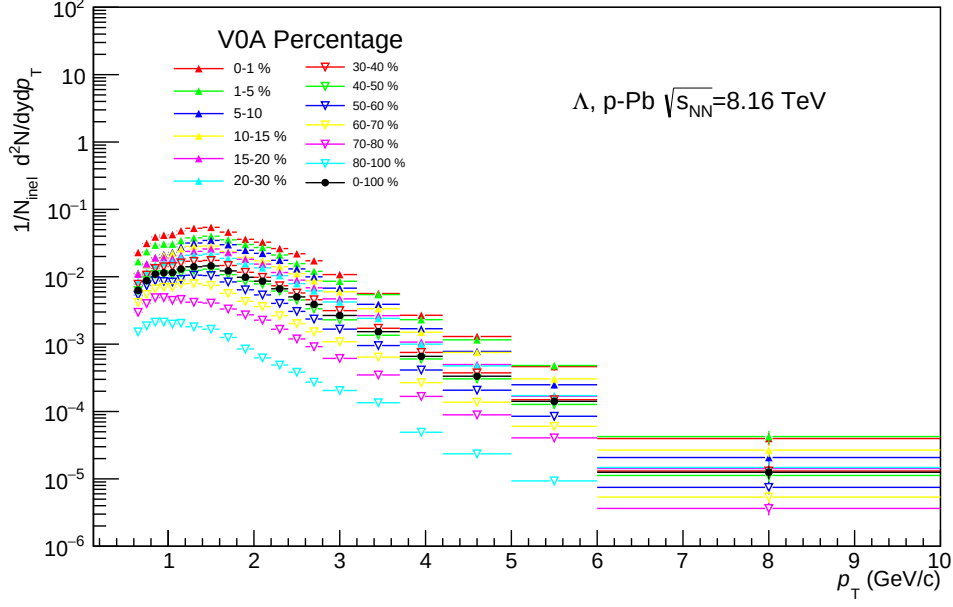


(a)

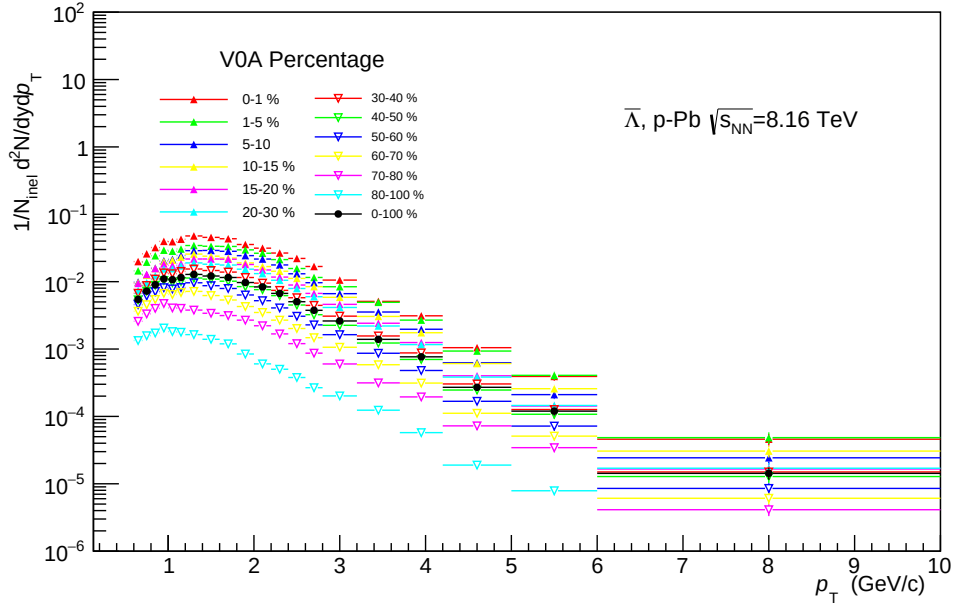


(b)

Figure 5.4: (a) Feed-down matrix for Λ , (b) Feed-down matrix for $\bar{\Lambda}$

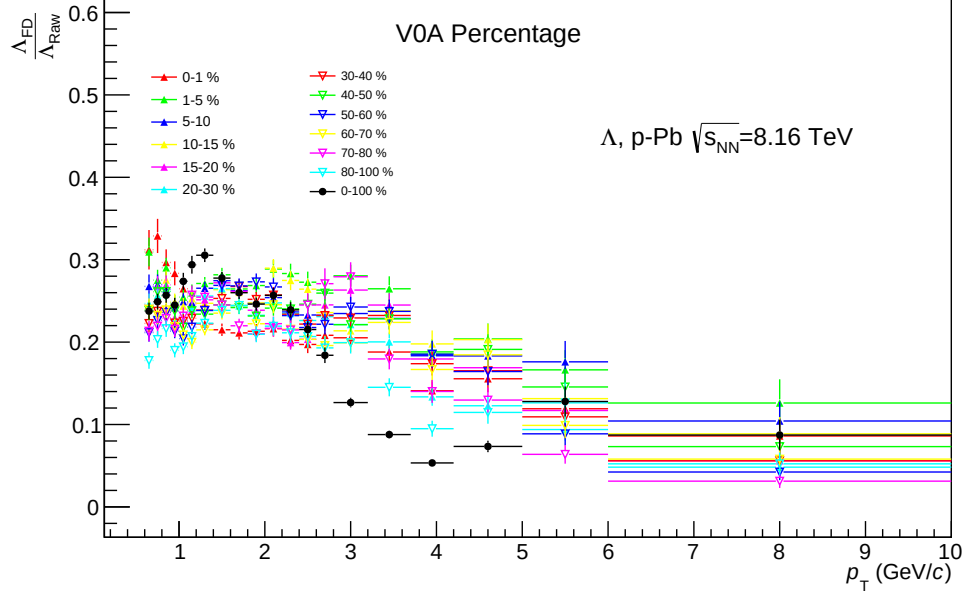


(a)

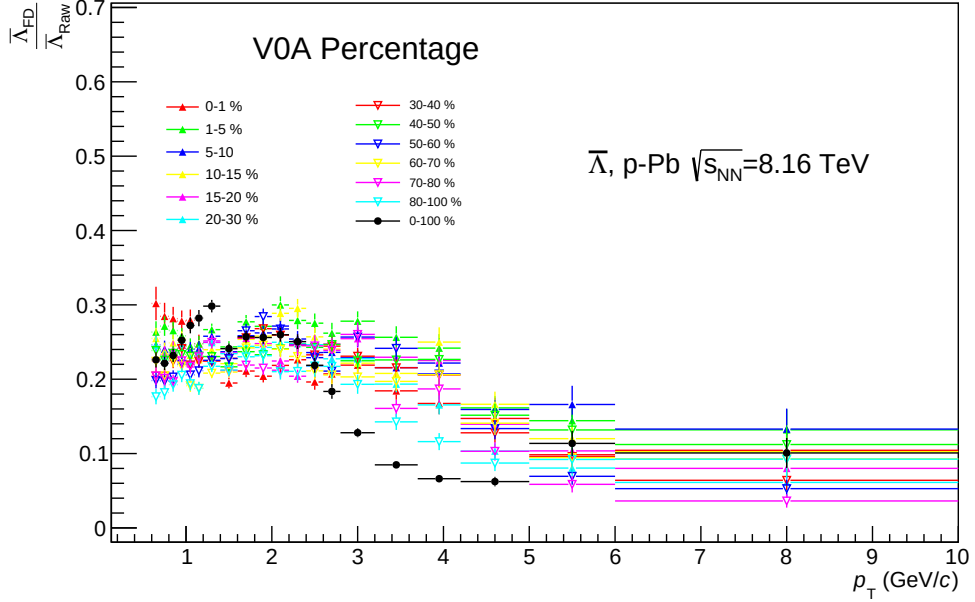


(b)

 Figure 5.5: (a) Feed-down spectra for Λ , (b) Feed-down spectra for $\bar{\Lambda}$



(a)



(b)

 Figure 5.6: (a) Feed-down ratio for Λ , (b) Feed-down ratio for $\bar{\Lambda}$

Two different methods are considered:

- **MC Ratio:** In the MC Ratio (Monte Carlo Ratio) method, the numerator is filled with Λ (or $\bar{\Lambda}$) from decays of both charged and neutral Ξ . In this case the ratio Ξ^- / Ξ^0 as provided by the Monte Carlo generator used to compute the feed-down matrix element F_{ij} is used.
- **Double Counting:** In this method, the numerator is filled with Λ (or $\bar{\Lambda}$) from decays of only charged Ξ , and the resulting charged Ξ feed-down is multiplied by a factor two (assume the production of Ξ^0 is the same as for $\Xi^\pm$).

The $\frac{dN}{dp_T}(\Xi^-)$ at the numerator is obtained by fitting the extrapolated charged Ξ spectra (in our case) by a Levy-Tsallis function. The feed-down spectra removed from raw spectra is shown in Figure 5.5 for the each multiplicity class case using MC Ratio method described above.

We have implemented two different methods for obtaining charged Ξ spectra in different multiplicity bins:

- Scaling the charged Ξ spectra at $\sqrt{s_{NN}} = 5.02$ TeV in p-Pb collision system using the measured proton spectra in p-Pb $\sqrt{s_{NN}} = 5.02$ TeV and 8.16 TeV using the equation below:

$$\left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{at\ 8.16\ TeV} = \left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{at\ 5.02\ TeV} \times \frac{\left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{proton\ spectrum\ at\ 8.16\ TeV}}{\left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{proton\ spectrum\ at\ 5.02\ TeV}} \quad (5.18)$$

- Scaling the charged Ξ spectra at $\sqrt{s_{NN}} = 5.02$ TeV in p-Pb collision system using General Purpose DPMJet production (LHC17f3b cent and LHC17f2b cent) at

$\sqrt{s_{NN}} = 5.02$ TeV and 8.16 TeV using the equation below:

$$\left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{at\ 8.16\ TeV} = \left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{at\ 5.02\ TeV} \times \frac{\left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{MC\ at\ 8.16\ TeV}}{\left(\frac{1}{N_{inel}} \frac{d^2 N}{dy dp_T} \right)_{MC\ at\ 5.02\ TeV}} \quad (5.19)$$

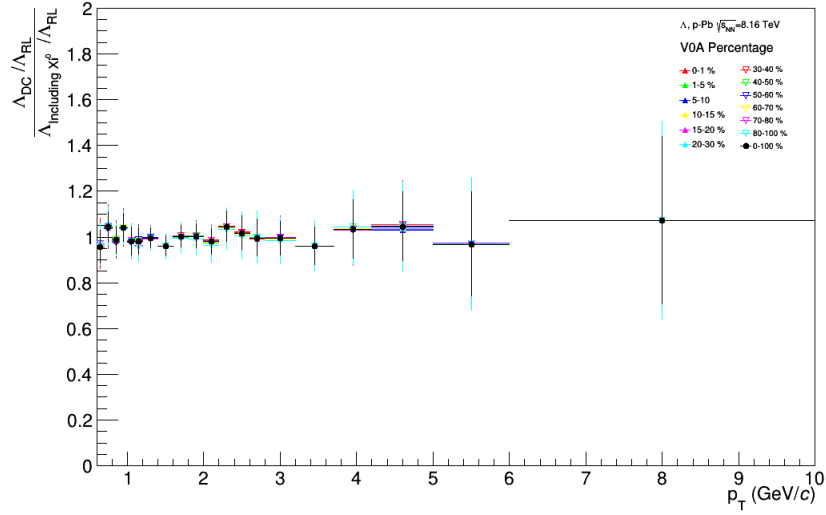
The first method is taken as default and other is used to compute the systematics. The feed-down fraction was further computed as function of multiplicity. The feed-down fractions computed by the “MC Ratio” method as shown in Figure 5.6 .

The total systematics due to feed-down comprises of following components:

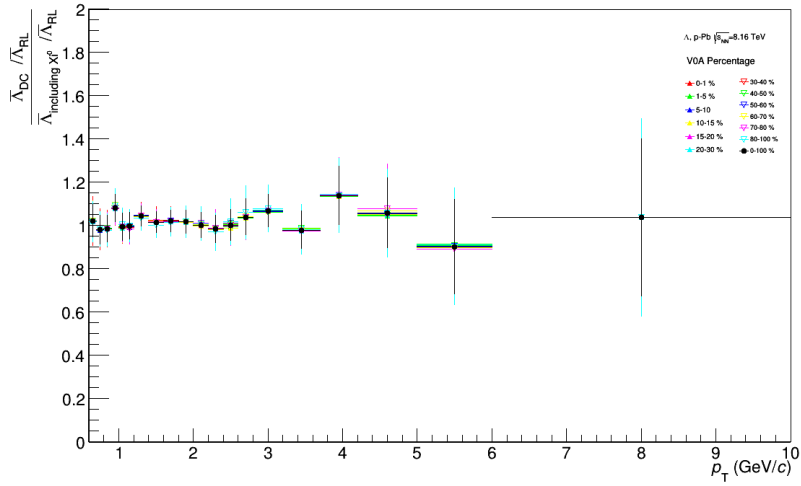
- The feed-down matrix is computed in minimum bias events (Figure 5.4) and is used in this procedure, since it has been observed that this matrix is not multiplicity dependent. For this choice a safety 2% p_T independent systematic has been added on the top of the other systematic uncertainties.
- The uncertainty due to the method used to fill the feed-down matrix (“MCRatio”/“Double Counting”). In Figure 5.7 the ratio of the feed-down spectra removed in the two cases for Λ and $\bar{\Lambda}$ is shown.
- The uncertainty from the Ξ spectra. Scaling using proton spectra is considered the default and DPMJet scaling is used to compute systematics (shown in Figure 5.8)

5.5.2 Efficiency x Acceptance Correction

Since the detector does not detect all of the V^0 particles created in a collision, an efficiency correction has to be made to correct for any shortcomings in the detector performance as well as track reconstruction and selection deficiencies. By efficiency, here, is meant the detector efficiency x detector acceptance and is defined for each p_T



(a)



(b)

Figure 5.7: Ratio of feed-down spectra obtained using DC and MC Ratio method for Λ (a) and $\bar{\Lambda}$ (b)

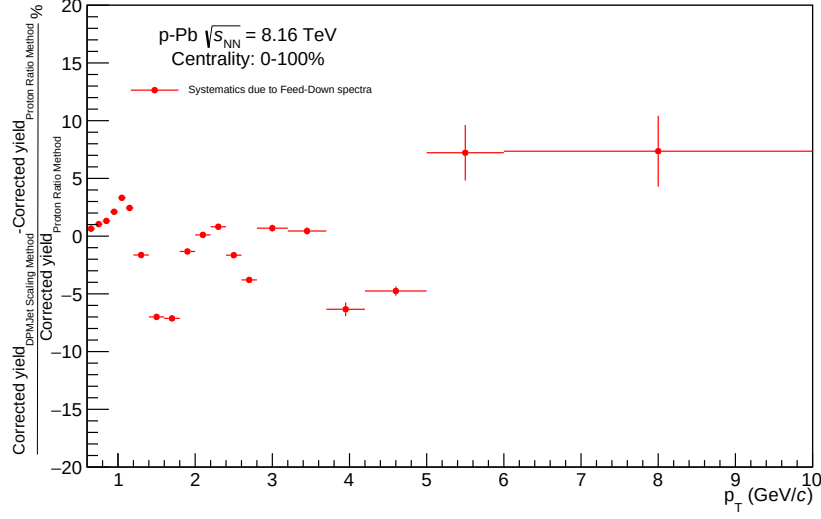


Figure 5.8: Systematics obtained due to charged Ξ spectra used for feed-down correction.

bin as

$$\epsilon(p_T) = \frac{N_{\text{reconstructed+associated}}(p_T)}{N_{\text{generated}}(p_T)} \quad (5.20)$$

The efficiencies are estimated by repeating the same analysis on simulated data and comparing the results to the known generated cascades. The only way the analyses differ is the Monte Carlo (MC) cascade candidates are checked to ensure they are associated with a true generated cascade. An efficiency correction is carried out by multiplying the uncorrected raw K_S^0 and Λ spectra by the efficiency correction factor $1/\epsilon(p_T)$.

5.5.2.1 Dependence on Multiplicity

Since the analysis is performed in multiplicity intervals, an important consideration is if the detection efficiency of V^0 s changes with multiplicity. It was confirmed that, within statistical uncertainties, ϵ is unaltered for K_S^0 (similarly for Λ and $\bar{\Lambda}$), as can be seen in Figures 5.9, 5.10 and 5.11. Consequently, in order to reduce statistical fluctuations in the efficiency determination, the p_T spectra using real data were

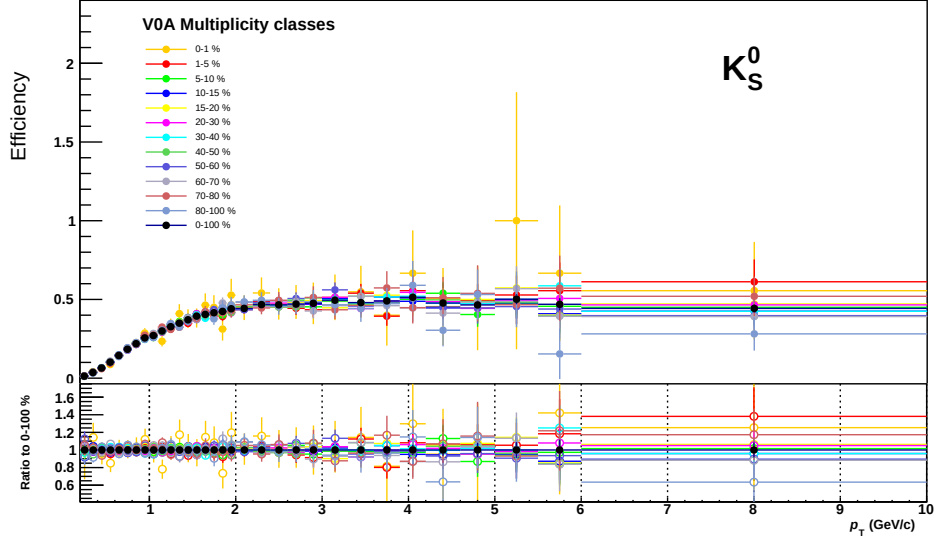


Figure 5.9: (Top panel) Efficiency for K_S^0 in different multiplicity bins as a function of p_T . (Bottom panel) comparison with centrality integrated bin.

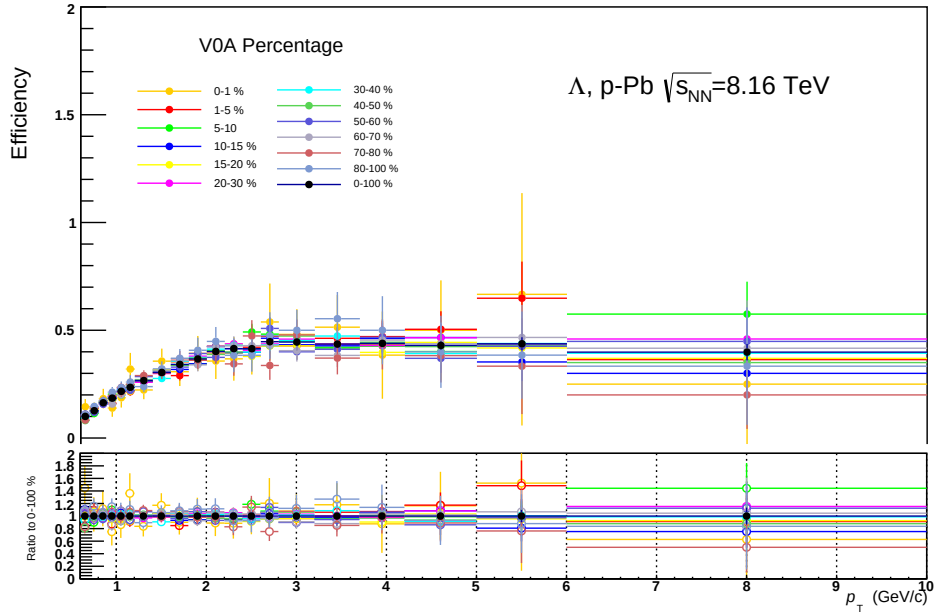


Figure 5.10: (Top panel) Efficiency for Λ in different multiplicity bins as a function of p_T . (Bottom panel) comparison with centrality integrated bin.

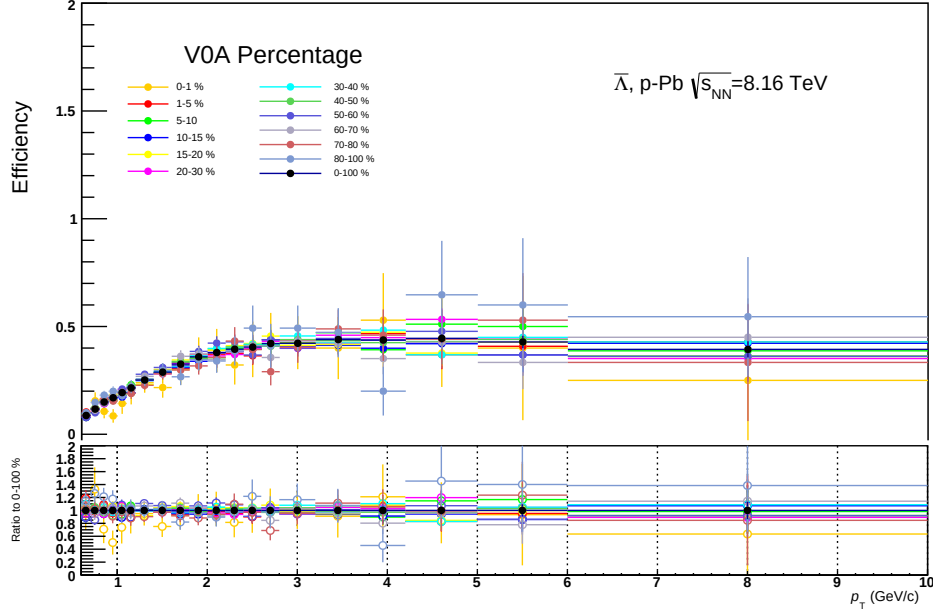


Figure 5.11: (Top panel) Efficiency for $\bar{\Lambda}$ in different multiplicity bins as a function of p_T . (Bottom panel) comparison with centrality integrated bin.

corrected with the multiplicity-integrated efficiency values.

5.6 Systematic Uncertainties

Several sources of systematic uncertainty are considered, following the general criterion followed in previous analyses [310–312]. Three broad p_T ranges are chosen to perform cut study:

For K_S^0 : low p_T bin (0.2-1), mid p_T bin (1-3) and high p_T bin (3-10) in GeV/ c

For Λ (and $\bar{\Lambda}$): low p_T bin (0.6-1.0), mid p_T bin (1.0-3.2) and high p_T bin (3.2-10), in GeV/ c .

Topological cuts are varied in order to evaluate the dependence of the final result on the particular value of the cut chosen in the analysis. Moreover, the dependence of the extracted yield on several other selection variables is taken into account. The variation is tuned using the following procedure:

- signal losses (with respect to the loosest cut considered) are plotted for both

data and MC. A comparison between the two tells if the cut is consistently applied to data and to the MC sample used for efficiency determination: big differences would point to a large discrepancy in the final result. As can be seen in Figures 5.12 and 5.13 (also from Appendix A), the agreement between data and MC is generally very good.

- at most four different cut values for each source are set for systematics estimation (two looser and two tighter with respect to the default one), making use of the signal loss plots in order to avoid losses larger than 10% and regions where real and MC data differ by more than a few percent (3 to 4 %).
- signal losses are then worked out in several p_T bins, with the purpose of avoiding cuts which, in specific p_T ranges, could lead to signal decreases larger than 10%. From Figure 5.12 (also from Appendix A), it is clear that in the case of any sources, the cuts can be found to be p_T dependent.

The final set of cuts used for systematic estimation is reported in Table 5.2 and 5.3. In Figure 5.12 and 5.13, red marker is used for low p_T bin, green marker for mid p_T bin and blue marker for high p_T bin. Circle symbol is used for real data and triangle symbol for MC data. Red vertical line shows the default cut for low p_T bin, green vertical line shows default cut for mid p_T bin and blue vertical line shows default cut for high p_T bin. In Figure 5.14 the relative difference between the corrected spectrum obtained with the new chosen cut and the one obtained with the default cut is shown for few systematic sources considered (plots for other variables is shown in Appendix A). Errors are here Roger Barlow's $\sigma_\Delta = |\sigma_1^2 - \sigma_2^2|$ [313] where σ_1 is the statistical error associated to the data-point obtained with the default cut and σ_2 is the same quantity obtained with the new chosen cut. If a source i is varied N times then the systematics due to the source is given by the formula below:

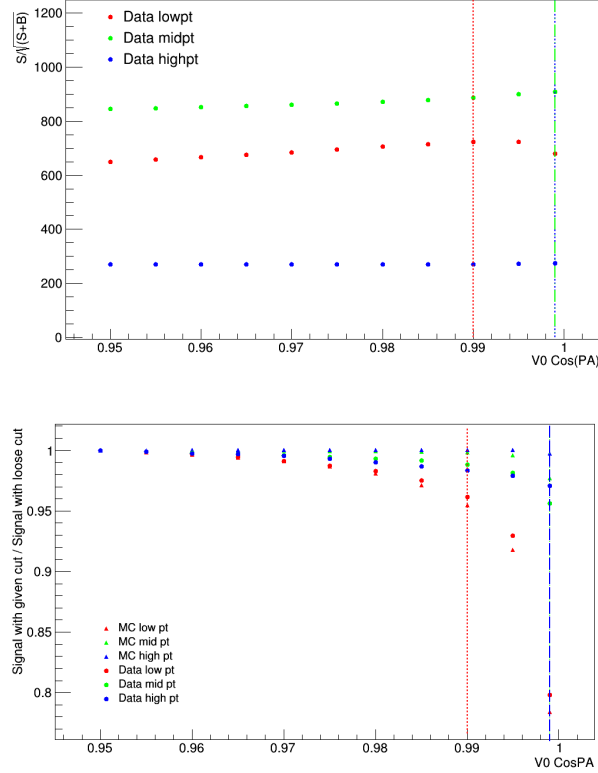


Figure 5.12: Significance and signal losses, with respect to loosest cut considered, in low, mid and high p_T bins for the variable cosine of pointing angle.

$$\delta_{i,sys} = \sqrt{\frac{1}{N} \sum \left(\frac{y_{def} - y_i}{y_{def}} \right)^2} \quad (5.21)$$

The relative deviations are considered significant (and added in quadrature in the computation of the systematic uncertainty) only if they were larger than one Roger Barlow's sigma. The systematics due to multiplicity independent efficiency (2%) is also considered. For Λ and $\bar{\Lambda}$ systematics due to feed-down is added and explained in the previous section. The contribution to systematics due to signal extraction includes: different sigma for signal extraction, choosing fitting functions (Pol1, Pol2 and exponential) for estimating background in the peak region (where default method is bin counting). On top of the above-mentioned sources of systematic error, the uncertainty on the material budget is taken from published analysis and is added in quadrature to all the other sources.

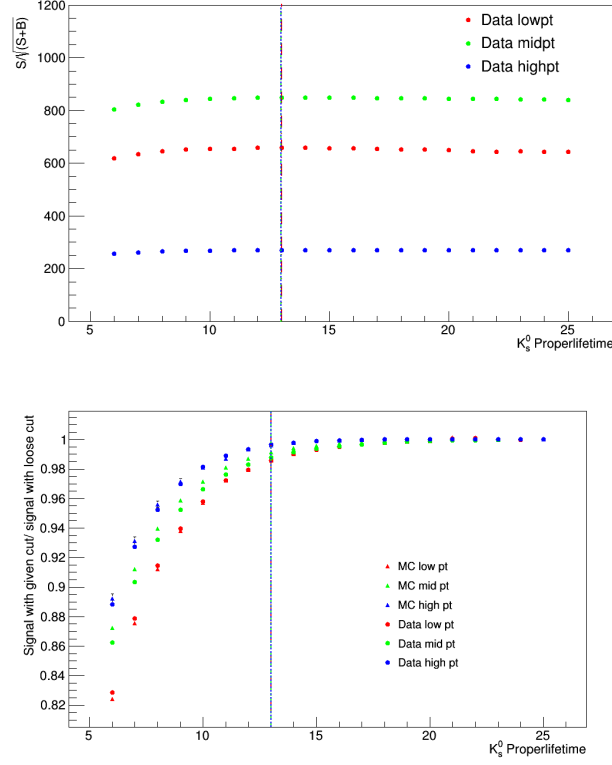
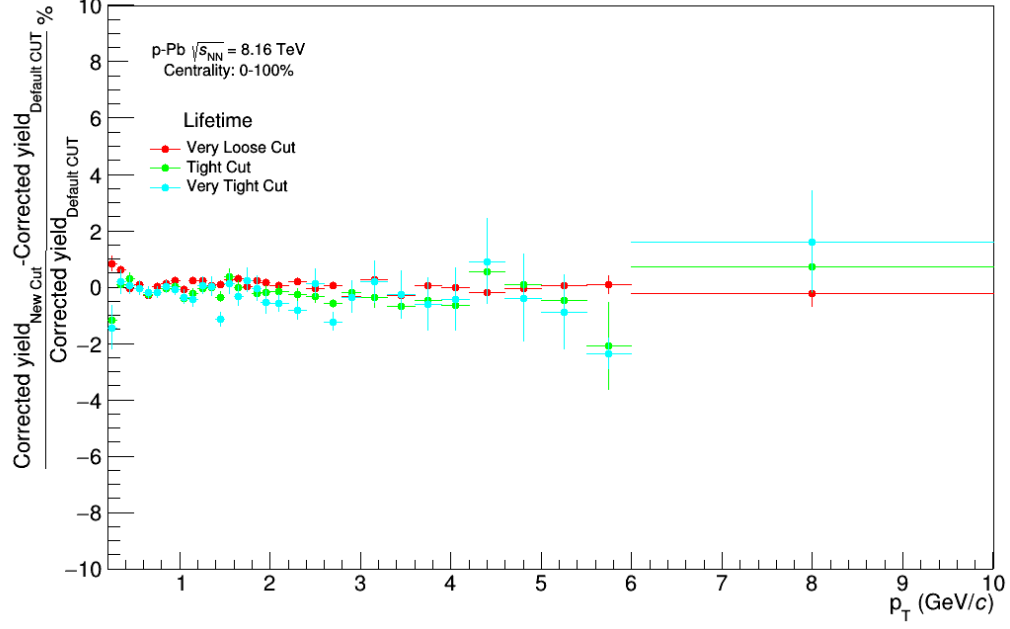


Figure 5.13: Significance and signal losses, with respect to loosest cut considered, in low, mid and high p_T bins for the variable properlifetime.

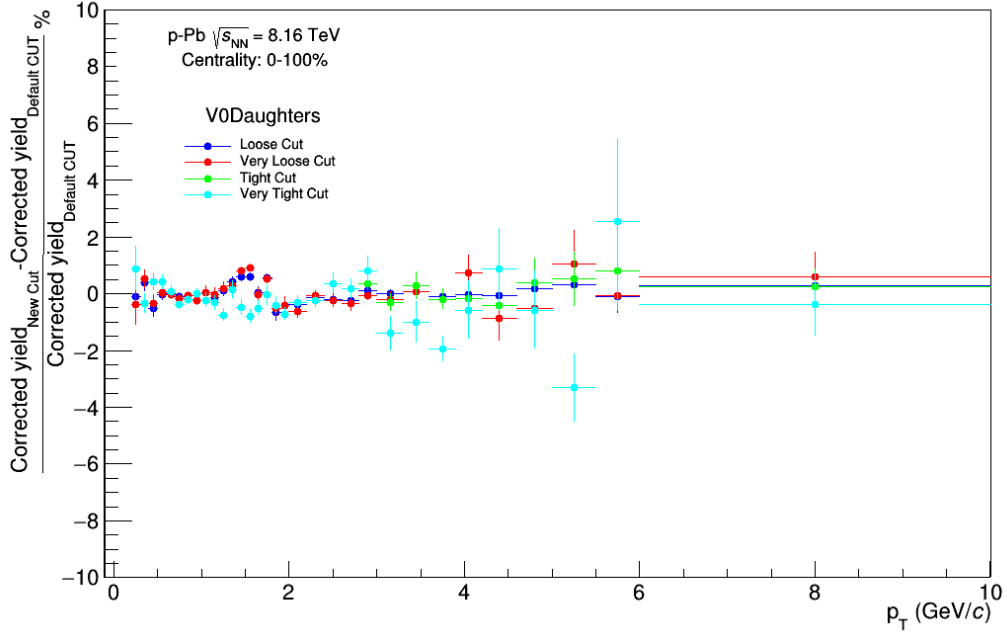
Same criteria as followed for K_S^0 is also followed for Λ and $\bar{\Lambda}$. The total systematic error for K_S^0 and Λ as shown in Figure 5.15 (for $\bar{\Lambda}$ refer to Appendix A)

5.6.1 Smoothing of systematic error

The systematic uncertainties exhibits a few significant fluctuation from bin to bin, while one might expect a smooth behavior of the systematic uncertainties. So we need to smooth the dips and spikes in the systematic uncertainties. To get more uniform systematic uncertainties the fractional systematic uncertainties in a problematic bin has been set to the average value w.r.t the adjacent bins except the first and last p_T bin. The each source of systematic uncertainty is smoothed individually. The smoothening procedure has been applied in all multiplicity case, according to the following approach. For p_T bin j use the uncertainty σ'_j obtained as the average among the initial systematic uncertainty, σ_j and $\sigma_{j\pm 1}$. Let σ_j and $\sigma_{j\pm 1}$ are the uncertainty

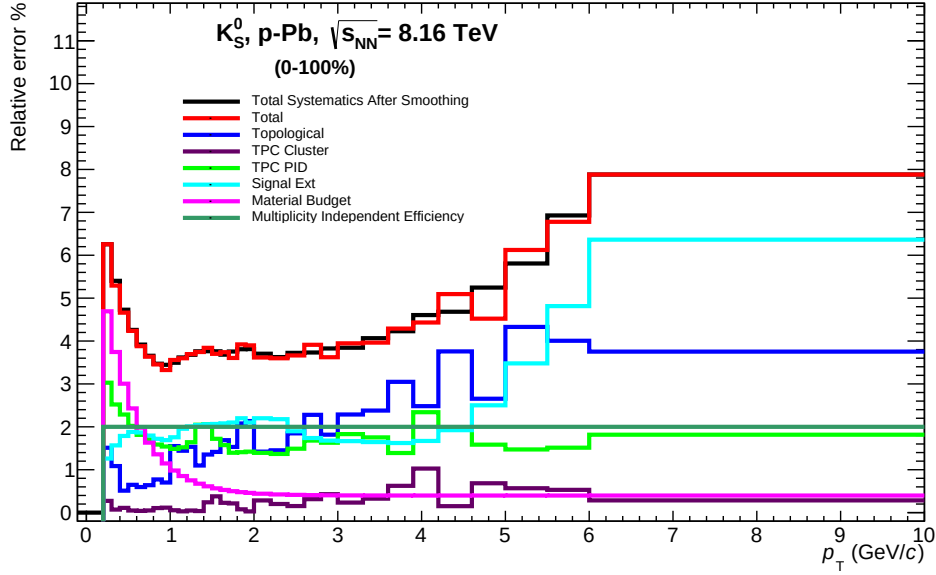


(a)

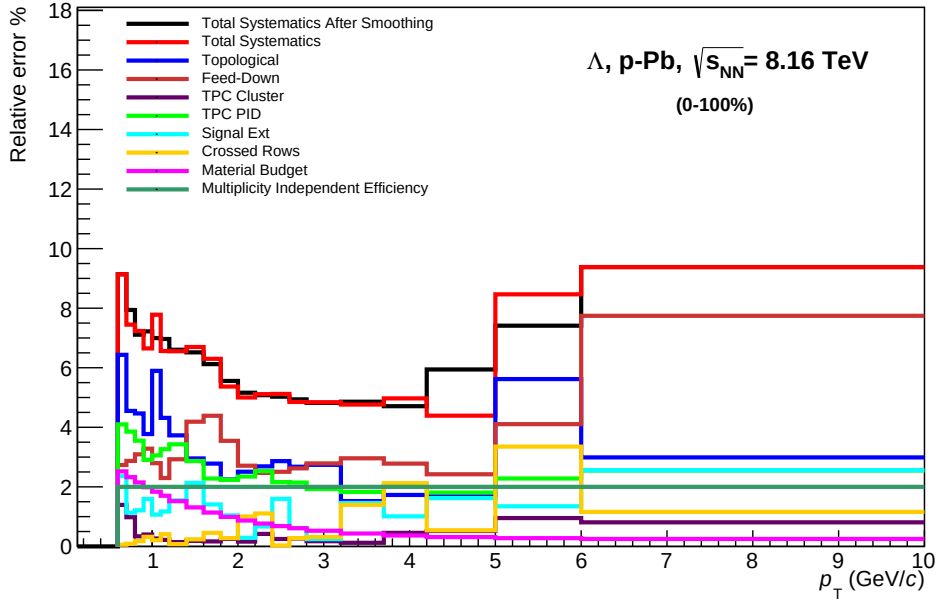


(b)

Figure 5.14: Signal deviations seen in corrected K_S^0 spectra due to cut selection changes for the variable Properlifetime (a) and V0Daughters (b). Error bars are Roger-Barlow errors representing an acceptable deviation due to statistical fluctuations.



(a)



(b)

Figure 5.15: Total systematics obtained in 0-100% multiplicity class for K_S^0 (a) and Λ (b)

of j th, $j+1$ th and $j-1$ th p_T bin. Then, systematics for p_T bin j after smoothing is given by :

$$\sigma'_j = \frac{\sigma_{j-1} + \sigma_j + \sigma_{j+1}}{3.0} \quad (5.22)$$

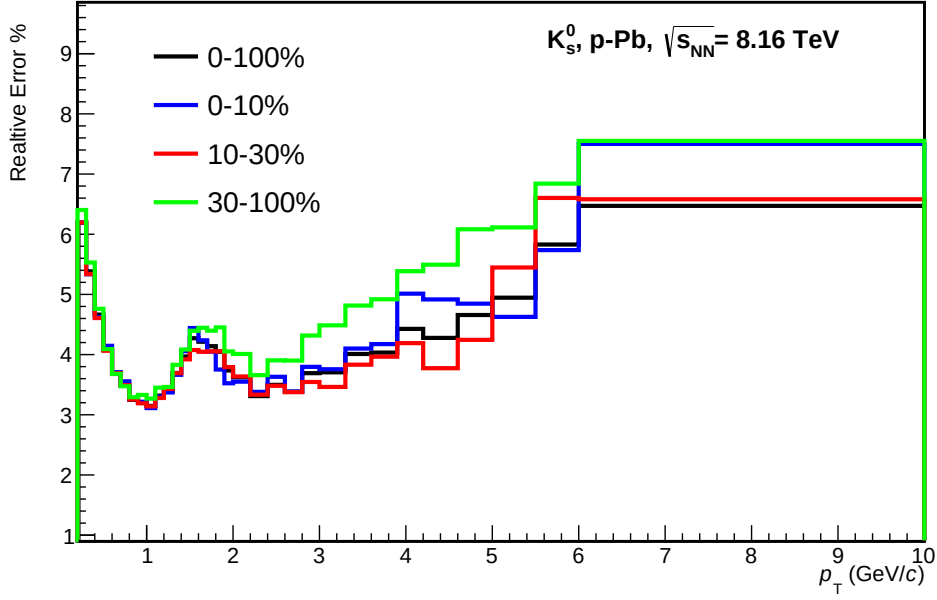
The first and last p_T bin remain unchanged during the smoothening procedure and this procedure can be followed more than once to remove large fluctuations. The total systematics after smoothing process is considered as final systematics.

5.6.2 Multiplicity Dependent Systematics

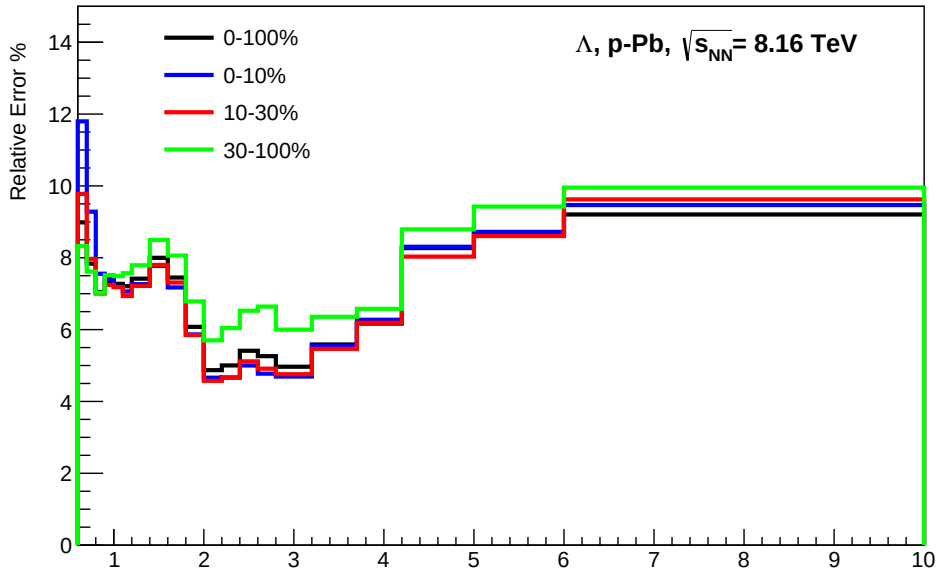
Figures 5.12 to 5.15, show the MB systematic deviations. Initially, it was attempted to calculate the total uncertainties in all 12 multiplicity classes by using the same method as in MB, but the errors obtained were too large due to large contributions from statistical fluctuations. Therefore a study was performed with larger samples, namely the 0-10%, 10-30% and 30-100% multiplicity classes, which roughly contained 1/3 of the total available statistics each. Figure 5.16 show how the systematics vary in these 3 classes. A similar study for systematic uncertainties were obtained in p-Pb 5.02 TeV. The final decision adopted was to assign the uncertainties that were obtained when using the integrated MB sample, and thus not to apply any multiplicity dependence on the uncertainties.

5.6.2.1 Uncorrelated Systematics

Having computed the errors in each multiplicity class, a study was performed to extract the component which is uncorrelated across multiplicity for each class. The error is largely correlated across multiplicity. However, a fraction of the total error could be dependent on the multiplicities considered. To attempt to isolate that part, particular deviations are compared to the ones in the minimum bias sample. The following ratio will determine the uncorrelated component of the total systematic



(a)



(b)

Figure 5.16: Total systematic uncertainties obtained for K_S^0 (a) and Λ (b) in three different multiplicity classes (0-10%, 10-30% and 30-100%) with roughly 1/3 of the MB statistics in each.

uncertainty for each selection cut in multiplicity class m :

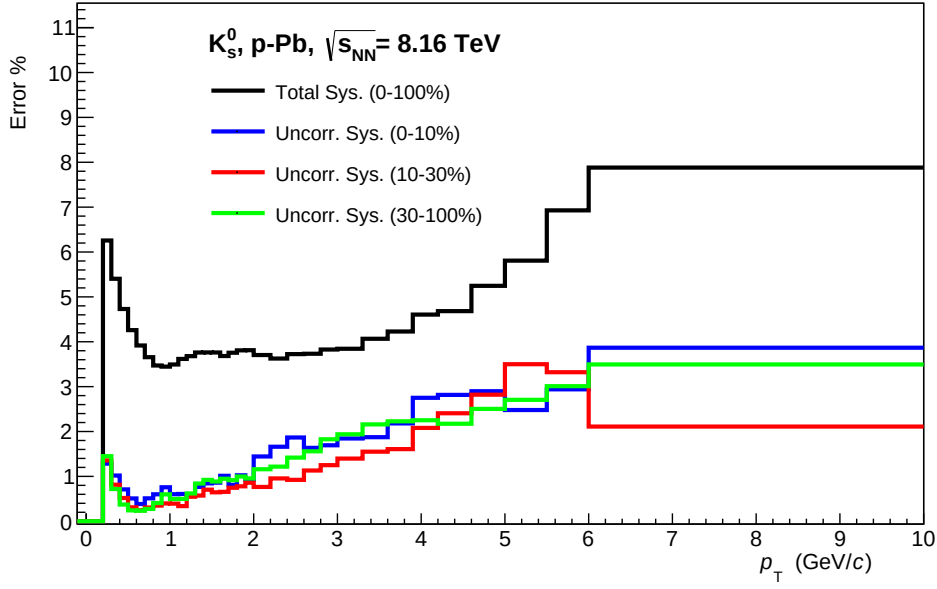
$$R = \frac{Y_m^{dev}/Y_m^{def}}{Y_{MB}^{dev}/Y_{MB}^{def}} \quad (5.23)$$

where Y_m^{dev} and Y_m^{def} are the measured deviated and default yield, respectively (in a given p_T interval). The deviated yield is the one obtained when varying a particular cut.

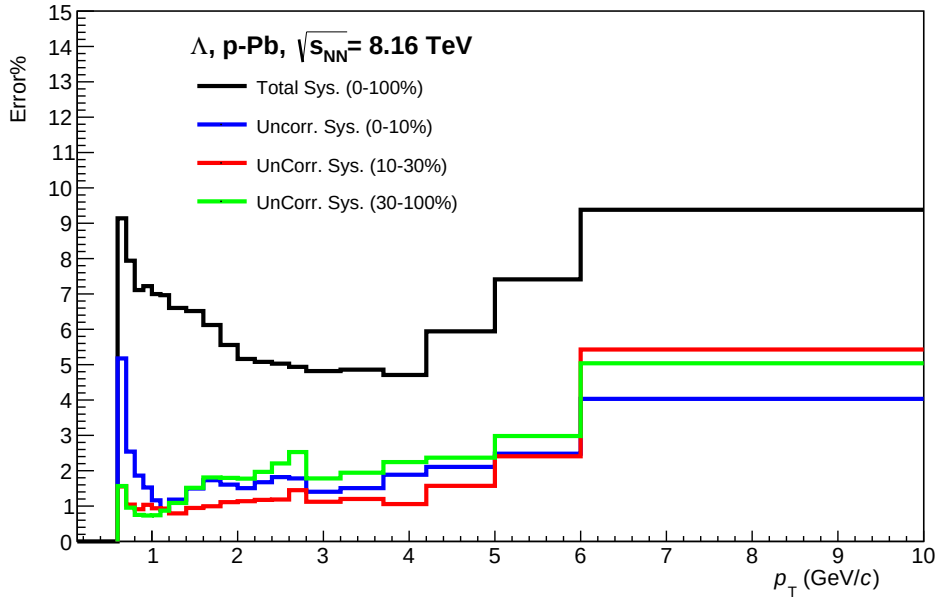
Just like for the multiplicity-dependence study of the total systematic uncertainties, it is crucial to have sufficient statistics in the samples used in order to obtain reliable results. In order to achieve that, the MB sample was sub-divided again into 3 sub-samples: the 0-10% , 10-30% and the 30-100% multiplicity classes. For the computation of the uncorrelated error, several contributions were cancelled. This is because they do not show a dependence on multiplicity, and consequently, the ratio in Equation 5.23 is taken to be 1. Only up-and-down fluctuations around unity are seen as a function of p_T for that ratio. These sources were therefore cancelled out of the sum for the total uncorrelated-in-multiplicity errors. The removed fluctuations are those produced by the variation of the following cuts:

- TPC Crossed Rows
- TPC Crossed Rows Over Findables
- Material Budget
- Proper Lifetime
- Cut on dE/dx
- Competing species rejection window

Figure 5.17, show the uncorrelated systematics error in three different multiplicity classes and its comparison to the total systematics error in minimum bias. It has



(a)



(b)

Figure 5.17: The multiplicity-uncorrelated uncertainties in comparison to the total (MB) systematic uncertainty for K_S^0 (a) and Λ (b)

been observed from this figure that the uncorrelated systematics error show same trend in all the three multiplicity classes and is always less than total systematics error in minimum bias.

Table 5.1: Selections applied to K_S^0 , Λ and $\bar{\Lambda}$

Topological Variable	K_S^0 low, mid, high pt	Λ ($\bar{\Lambda}$) low, mid, high pt
V0 transv. decay radius(cm)	1	1
DCA Negative Track to PV(cm)	0.1,0.04,0.02	0.04,0.05,0.03
DCA Positive Track to PV(cm)	0.1,0.04,0.02	0.03,0.02,0.02
V0 Cosine of Pointing Angle	0.990,0.999,0.999	0.985,0.995,0.950
DCA V0 Daughters(σ)	1.2,0.7,0.5	1.5,1.3,0.9
Selection	K_S^0 low, mid,high pt	Value for Λ ($\bar{\Lambda}$) low, mid, high pt
V0 Vertex Type	Offline Vertexer	Offline Vertexer
Rapidity Interval	$-0.5 < y < 0$	$-0.5 < y < 0$
Proper Lifetime (mL/p)(cm)	13	40
TPC dE/dx Selection(Real Data Only)	3	3
Primary Selection (MC Only)	IsPhysicalPrimary()	IsPhysicalPrimary()
MC Association (MC Only)	PDG code association	PDG code association
σ for signal extraction	4	4
Daughter Track Pseudorapidity Interval	$ \eta < 0.8$	$ \eta < 0.8$
Daughter Track $N_{crossed}$	70	70
Daughter Track $N_{crossed}/N_{findable}$	0.8	0.8
Competing species rejection window	$5MeV/c^2$	$10MeV/c^2$

Table 5.2: Values of the selection cuts applied for systematics estimation of K_S^0

Topological Variable	very loose	loose (low, mid, high p_T)	very tight (low, mid, high p_T)	tight (low, mid, high p_T)
V0 transv. decay radius	—	—	1.4, 1.4, 2.0	1.6, 1.8, 3.0
DCA Negative Track to PV	0.02	—	0.2, 0.1, 0.07	0.4, 0.17, 0.12
DCA Positive Track to PV	0.02	—	0.2, 0.1, 0.07	0.4, 0.17, 0.12
V0 Cosine of Pointing Angle	0.950	0.97	—	0.995, 0.999, 0.999
DCA V0 Daughters	2	1.8, 1.4, 0.6	1.2, 0.7, 0.3	0.85, 0.4, 0.15
Proper Lifetime (mL/p)	25	—	10.9.8	8, 7, 6
Daughter Track $N_{crossed}$	—	—	80	90
Daughter Track $N_{crossed}/N_{findable}$	—	—	0.9	—
TPC PID σ	—	4	2.5	—
σ for signal extraction	—	5	[2, 3, 3]	—
Competing species rejection window	None	—	—	—

Table 5.3: Values of the selection cuts applied for systematics estimation of Λ ($\bar{\Lambda}$)

Topological Variable	very loose	loose (low,mid,high p_T)	very tight (low,mid,high p_T)	tight (low,mid,high p_T)
V0 transv. decay radius	—	1.1,1.25,1.7	1.5,1.8,2.0	1.7,3,4.8
DCA Negative Track to PV	0.02	—	0.16(0.05,0.05,0.04)	0.26(0.09,0.07,0.05)
DCA Positive Track to PV	0.02	—	0.05,0.05,0.04(0.16)	0.09,0.07,0.05(0.26)
V0 Cosine of Pointing Angle	0.950	0.97	—	0—
DCA V0 Daughters	2	1.8,1.6,1	1.3,0.9,0.4	1,0.6,0.3
Proper Lifetime (mL/p)	100	50	—	—
Daughter Track $N_{crossed}$	—	—	85	90
Daughter Track $N_{crossed}/N_{findable}$	—	—	0.9	—
TPC PID σ	—	4	2.5	—
σ for signal extraction	—	5	[2,3,3]	—
Competing species rejection window	None	—	—	—

Chapter 6

In this chapter, the results for K_S^0 , Λ and $\bar{\Lambda}$ are included and their interpretations are discussed. For the better understanding, the results of the present analysis are compared to the already published results from the previous analyses.

Chapter 6

Results and discussions

6.1 Transverse Momentum Spectra

The aim of the analysis is to measure the production rates of strange baryons in p–Pb collisions as a function of the transverse momentum and integrated in p_T for a particular multiplicity class, in the central rapidity region. To achieve this goal many steps are required to be followed which are mentioned below:

- The signal should be extracted from the invariant mass distribution for each momentum interval following the procedure described in Section 5.4. The p_T distribution obtained is the so-called “raw” spectrum (shown in Appendix A).
- This spectrum is subject to all the inefficiencies and the acceptance limitation of the detector, and must therefore be corrected for these effects. Moreover, Λ ($\bar{\Lambda}$) require feed-down correction. These signal correction methods are explained in Section 5.5.
- The corrected spectrum must be normalized to the number of events, p_T interval and rapidity interval.
- Finally, any source of systematics must be studied and corresponding uncer-

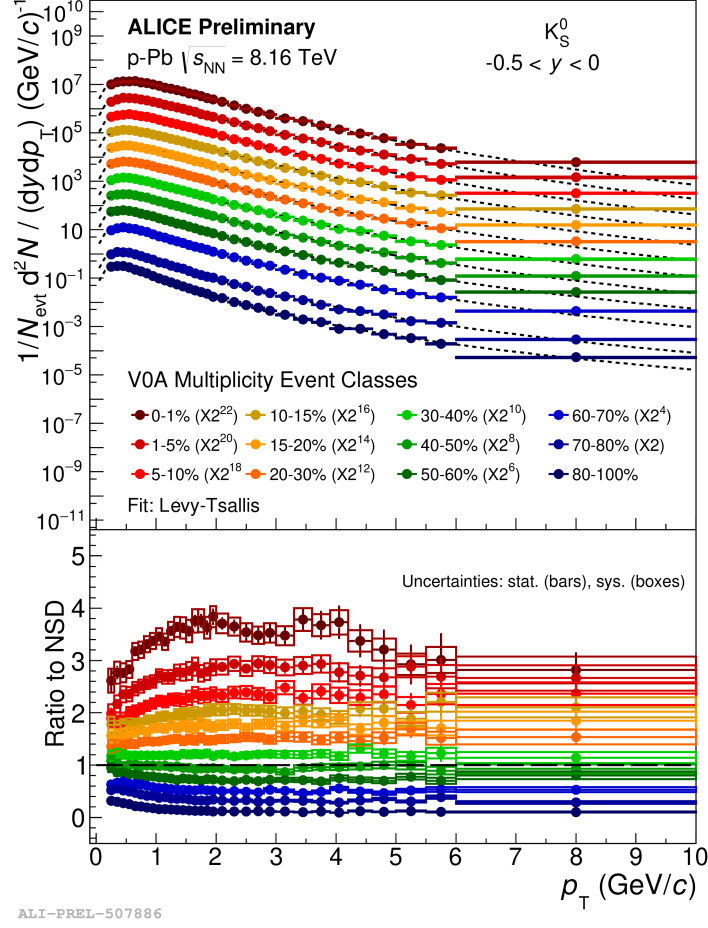


Figure 6.1: (Top panel): p_T spectra of K_S^0 in different V0A multiplicity classes. (Bottom panel): Ratio of the p_T spectra in different multiplicity classes to the minimum bias p_T spectrum.

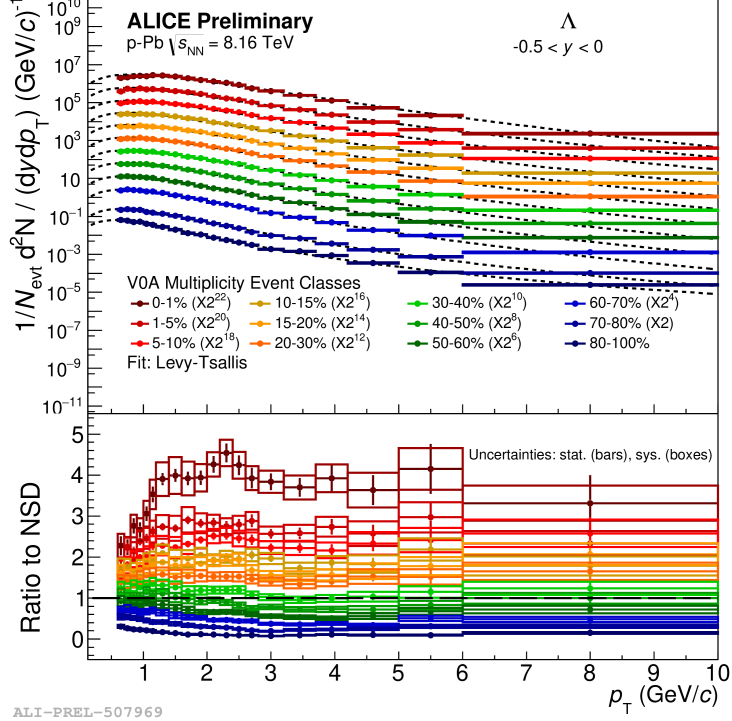
ainties have to be estimated (Section 5.6)

The corrected spectra can be expressed in the following way:

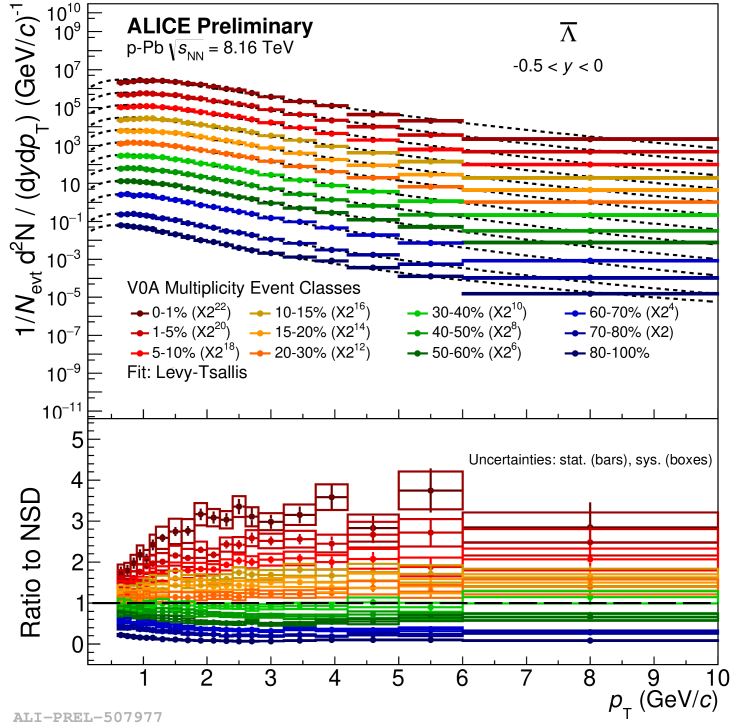
$$\frac{1}{N_{\text{events}}} \frac{d^2N}{dy dp_T}(p_T) = \frac{1}{N_{\text{events}}} \frac{1}{\Delta y} \left(\frac{S_{\text{raw count}}}{\epsilon(p_T)} \frac{1}{\Delta p_T} \right)_{i \text{ } p_T \text{ bin, } j \text{ multiplicity class}} \quad (6.1)$$

where $d^2N/dp_T dy$ is the differential yield in transverse momentum and rapidity for a particular multiplicity class, N_{events} is the number of events used in the analysis, $S_{\text{raw count}}$ is the raw signal of V0 extracted in a given p_T bin, $\epsilon(p_T)$ is the detector acceptance-efficiency correction in the same p_T bin, Δp_T is the width of the given p_T

6.1. TRANSVERSE MOMENTUM SPECTRA



(a)



(b)

Figure 6.2: (Top panel): p_T spectra of Λ (a) and $\bar{\Lambda}$ (b) in different V0A multiplicity classes. (Bottom panel): Ratio of the p_T spectra in different multiplicity classes to the minimum bias p_T spectrum.

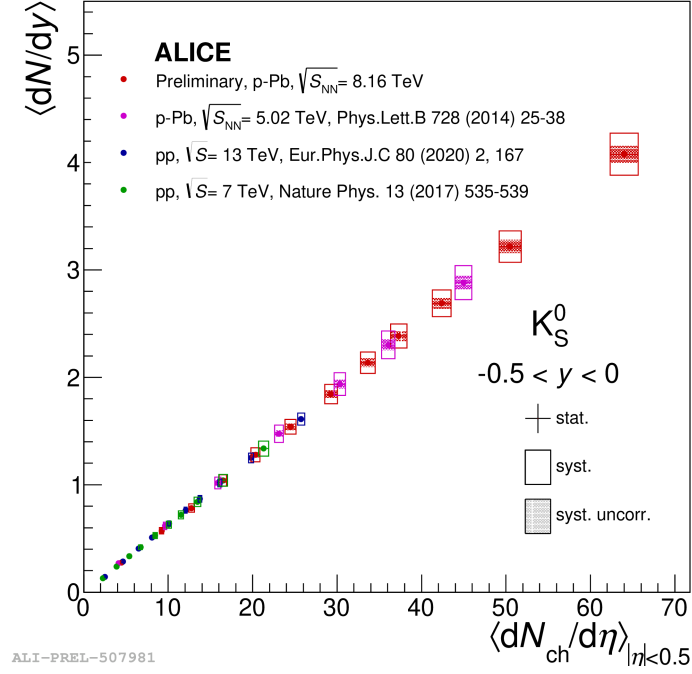
bin and Δy is the chosen rapidity interval.

The corrected spectra for the V0s in $-0.5 < y < 0$ are shown in Figures 6.1 and 6.2. Statistical and systematic errors have been separated and represented with lines and boxes, respectively. The effect of radial flow can be better appreciated in these figures where the spectra of the particles in a particular multiplicity class are divided by the spectra in multiplicity integrated class. The final p_T spectra are fitted with a Lévy-Tsallis function described in Appendix B, which is shown in the plots by the dashed lines.

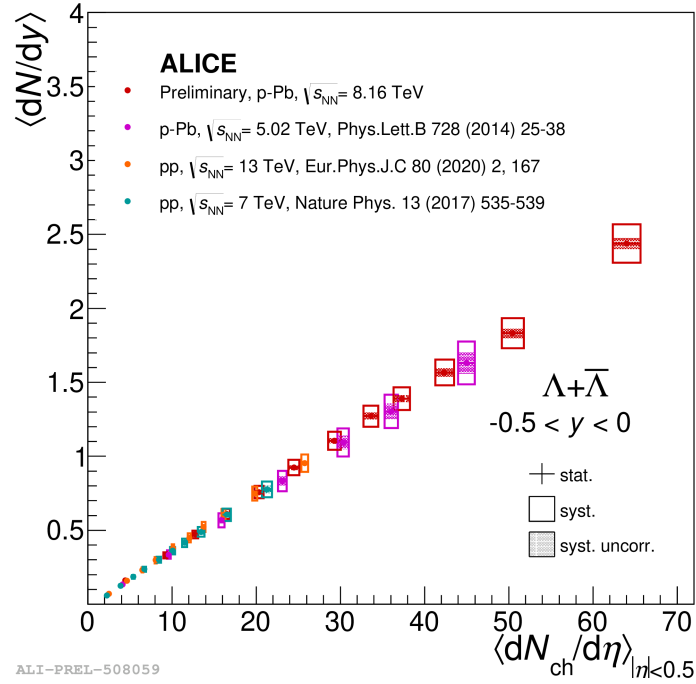
Similar results were previously obtained in Pb–Pb collisions where the spectra become harder with increasing centrality [314]. This hardening is found to be mass dependent which is consistent with the presence of a radial flow. The common expansion velocity of the fireball manifests itself as a mass dependent shift of peak of the spectra towards high p_T . At high p_T ($p_T > 5 \text{ GeV}/c$) the slopes do not depend on the particle species as the production of hard partons is dominated by fragmentation processes.

6.1.1 p_T -integrated yields and average p_T

The p_T -integrated production yields of K_S^0 and $(\Lambda + \bar{\Lambda})$ have been extracted from the measured spectra. The yields, $\langle dN/dy \rangle$, and mean transverse momentum, $\langle p_T \rangle$, are computed by extrapolating the spectra down to $p_T = 0 \text{ GeV}/c$. The extrapolation is performed in p–Pb collisions by fitting the single spectra using a Lévy-Tsallis function (described in Appendix B). The estimation of the systematic uncertainty on the yields was carried out by fitting the highest and lowest spectra obtained by shifting all data points within their systematic uncertainty. In addition, the systematic uncertainty on the extrapolation was estimated by using a set of fit functions (shown in Appendix A (Figure A.14a)) that differ in the functional slope at low p_T and by taking the maximum deviation from the Lévy-Tsallis function.

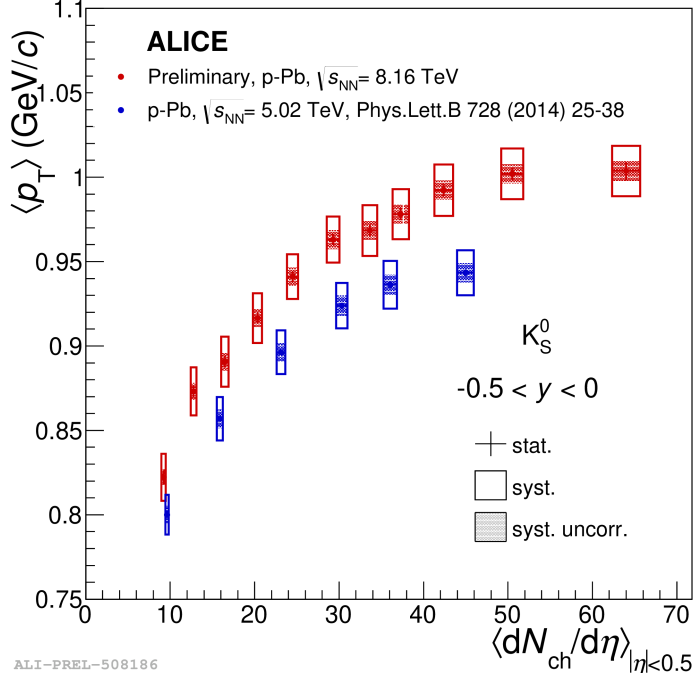


(a)

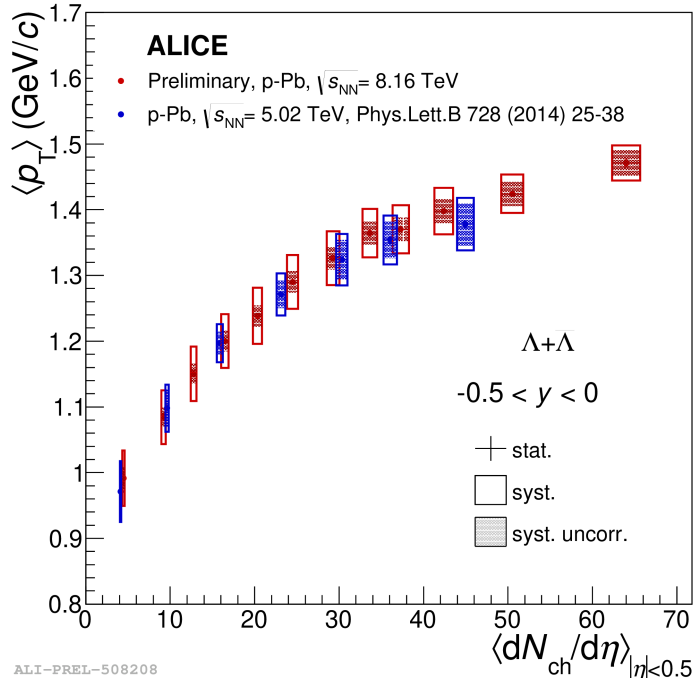


(b)

Figure 6.3: p_T integrated yield per unit rapidity, dN/dy as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$ for of K_S^0 (a) and $(\Lambda + \bar{\Lambda})$ (b)



(a)



(b)

Figure 6.4: Average transverse momentum, $\langle p_T \rangle$ as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta|<0.5}$ for of K_S^0 (a) and $(\Lambda + \bar{\Lambda})$ (b)

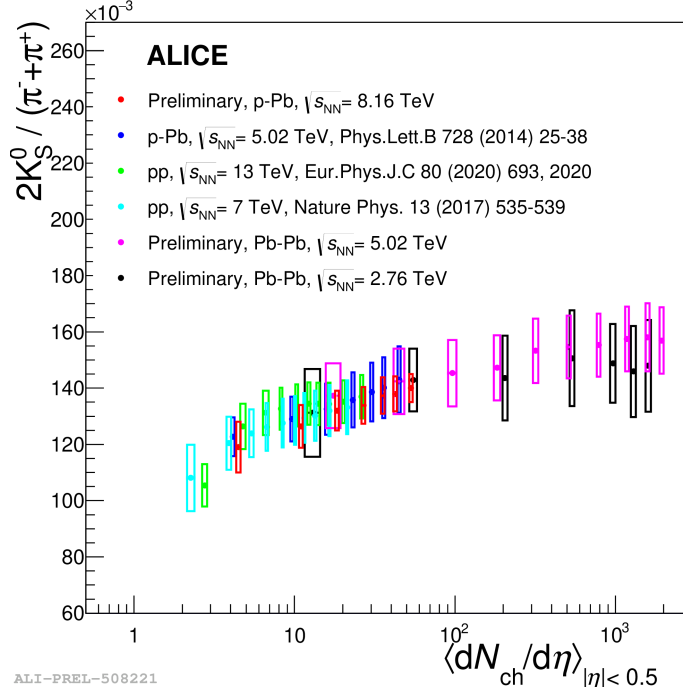
6.2. RATIO TO NON-STRANGE PARTICLE YIELDS: STRANGENESS ENHANCEMENT

The particle yields per unit of rapidity for K_S^0 and $(\Lambda + \bar{\Lambda})$ as a function of $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta|<0.5}$ for p-Pb collisions at $\sqrt{s_{\text{NN}}} = 8.16$ TeV are shown in Figures 6.3a and 6.3b respectively. The steep increase in the relative production of K_S^0 and $(\Lambda + \bar{\Lambda})$ as a function of $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta|<0.5}$ is observed. Comparison of the results at $\sqrt{s_{\text{NN}}}$ at 8.16 TeV with pp collisions at $\sqrt{s} = 7$ TeV [319], $\sqrt{s} = 13$ TeV [320], p-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [317] is made. It is also observed that there is only multiplicity dependence without any energy dependence and also independent of the initial colliding system.

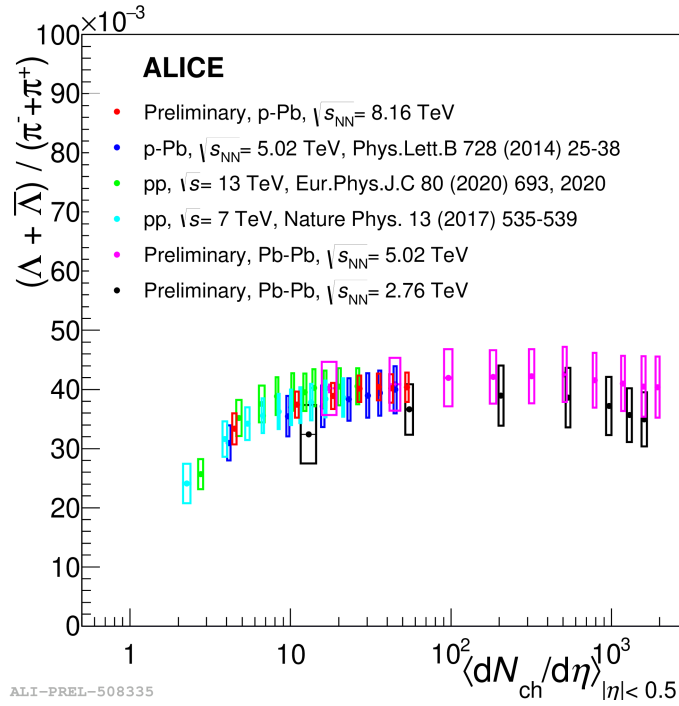
The $\langle p_T \rangle$, computed from the spectra of K_S^0 and $(\Lambda + \bar{\Lambda})$ measured in p-Pb collisions is reported as a function of the charged particle multiplicity density ($\langle dN_{\text{ch}}/d\eta \rangle_{|\eta|<0.5}$) in Figures 6.4a and 6.4b. The increase in the average transverse momentum is larger for particles with heavier masses, as expected in the presence of radial flow. For K_S^0 , the comparison with p-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [317] shows that the effect of radial flow is larger at $\sqrt{s_{\text{NN}}} = 8.16$ TeV while Λ does not show any energy dependence. These results illustrate the spectrum hardening, where there is a shift in the peak of the p_T distributions towards larger values of p_T at higher multiplicities supporting the evidence for collective behaviour in small systems that can be described by hydrodynamic models [318].

6.2 Ratio to non-strange particle yields: Strangeness Enhancement

The p_T -integrated yields of K_S^0 and $(\Lambda + \bar{\Lambda})$ scaled to the ones of pions are reported as measured in p-Pb collisions at $\sqrt{s_{\text{NN}}} = 8.16$ TeV in Figures 6.5a and 6.5b. The results of the analysis presented in this work are reported as a function of the charged particle multiplicity density and compared with the ones measured across several collision systems: pp collisions at $\sqrt{s} = 7$ TeV [319], $\sqrt{s} = 13$ TeV [320], p-Pb collisions



(a)



(b)

Figure 6.5: (a) $K_S^0/(\pi^+ + \pi^-)$ and (b) $(\Lambda + \bar{\Lambda})/(\pi^+ + \pi^-)$ yields as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$.

at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [317] and Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [321] and $\sqrt{s_{\text{NN}}} = 5.02$ TeV [322]. While the comparison of particle ratios measured in different systems highlights a continuous evolution as a function of the particle multiplicity, the new data confirm the trends observed at highest energy in p–Pb collisions. The observation of similar trends at $\sqrt{s_{\text{NN}}} = 2.76$ TeV and $\sqrt{s_{\text{NN}}} = 5.02$ TeV [321, 322] is consistent with the expectations from the thermal-statistical model, which predicts no significant evolution between the two energies.

6.3 Baryon Over Meson Ratio

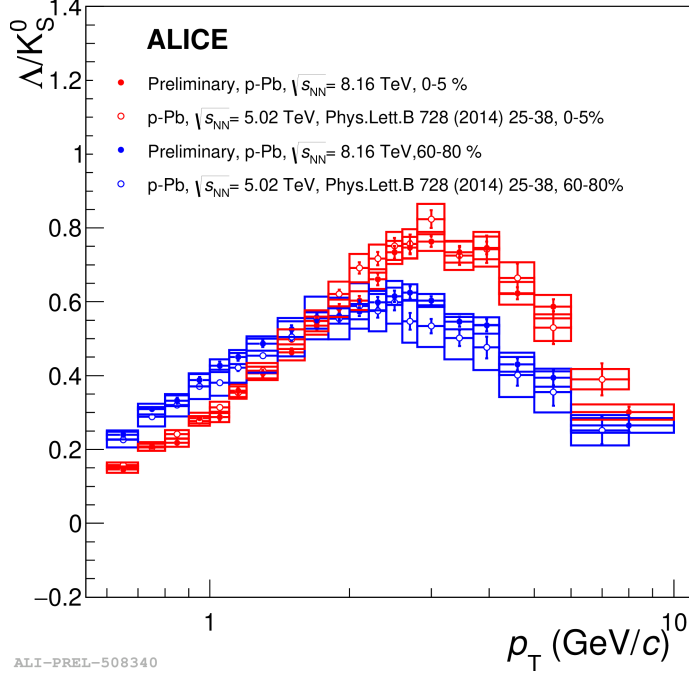
It is also one of the motivations of the thesis to investigate the strangeness baryon/meson ratio, Λ/K_S^0 for high multiplicity as well as low multiplicity regions of p–Pb collisions. The comparison of this ratio in p–Pb at $\sqrt{s_{\text{NN}}} = 8.16$ TeV and p–Pb at 5.02 TeV [317] is also made.

Similar to Pb–Pb collisions [323], we have observed that this ratio for p–Pb shows a rise at low p_T due to an increase in radial flow as well as the hard scattering effects at high p_T . It also shows a clear difference at intermediate p_T between particle yields ratio of the high multiplicity and the low multiplicity classes. Thus these results are consistent with the one seen in Pb–Pb collisions, where it is interpreted as the collective expansion of the QGP.

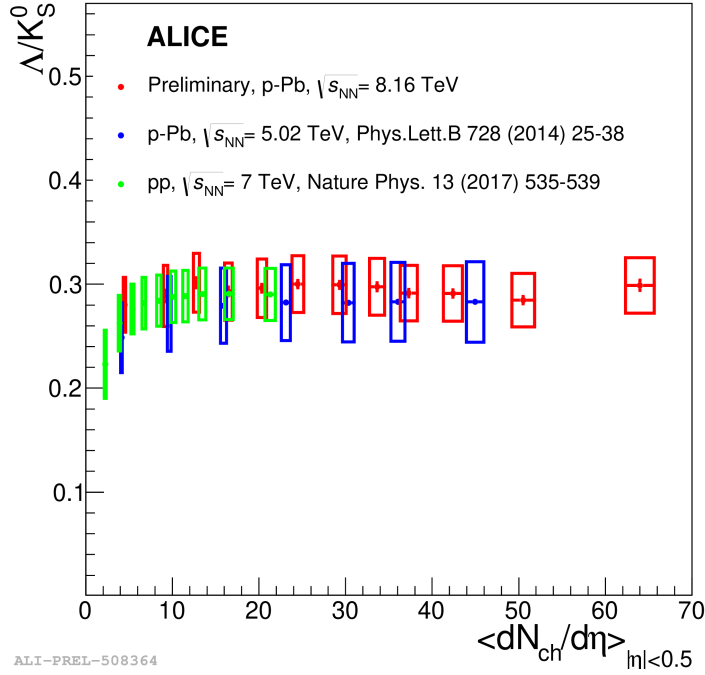
Two observations can be made from this study shown in Figure 6.6a:

- The ratio, Λ/K_S^0 does not appear to be energy dependent instead it is related to the multiplicity. The peak is more enhanced for high multiplicity case (0-5 %) when compared to the low multiplicity class (60-80 %).
- A shift in peak towards the high p_T region is observed while moving to the high multiplicity class (0-5 %) from the low multiplicity class (60-80 %).

Both these observations add evidence to the suggested QGP explanation i.e. an effect



(a)



(b)

Figure 6.6: (a) Λ/K_S^0 as a function of p_T in 0–5% and 60–80% multiplicity classes and two different energies, (b) Λ/K_S^0 as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$

of recombination and/or radial flow.

6.3.1 Redistribution and enhanced baryon production

The measured Λ/K_S^0 ratio suggests at first glance that there is enhanced baryon production at intermediate p_T . This can be explained by coalescence models [324, 325]. According to these models: In a medium, 3 quarks close to each other can easily exchange color to form a colorless state, while baryon production through fragmentation requires the production of six quarks, a baryon and an anti-baryon in order to conserve baryon number. Thus if coalescence is a dominant source of particle production, an enhancement in baryon production relative to meson should be observed.

An alternative model is that the enhancement at intermediate p_T is caused by a redistribution of particles in p_T ; baryons produced at low p_T may be pushed to higher p_T by collective flow [326]. Flow models can generate an excess of baryons at intermediate p_T without any increase in overall baryon production, the impact of the collective flow is greater on the heavier baryons than on the lighter mesons and so the baryons are pushed to higher p_T .

Coalescence models incorporate a similar effect, not only are baryons easier to produce through coalescence, their typical momenta is higher than mesons because they take in the total momentum of three quarks rather than two.

6.3.2 Baryon over meson ratio as a function of multiplicity

In Figure 6.6b, shows the Λ/K_S^0 as a function of multiplicity. The data is compared to pp collisions at $\sqrt{s} = 7$ TeV [319] and p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [317]. The following observations can be made from this plot:

- Baryon over meson ratio (Λ/K_S^0) as a function of multiplicity does not show any significant evolution with multiplicity. Moreover, there is no dependence on energy or system size for this ratio.

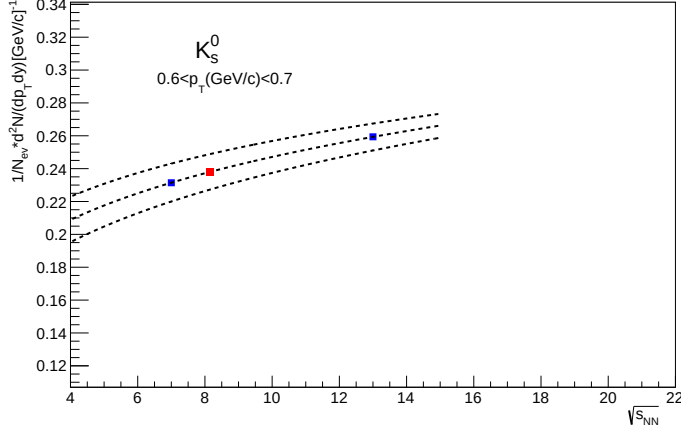


Figure 6.7: Fit to the dependence on $\sqrt{s_{\text{NN}}}$ of the pp yields for K_S^0 for a p_T bin. The fitted function is used to interpolate the value at 8.16 TeV (red point on the curve).

- Both the particles are neutral and strange while one of them is baryon and the other is meson and their ratio as a function of multiplicity is almost flat. This clearly means that the slope observed in Figure 6.5 is not because of any baryon/meson effect or mass effect but is strangeness enhancement effect.

6.4 Nuclear Modification Factor

The p–Pb collisions provide a good test to check whether the suppression observed in Pb–Pb collisions is due to the creation of QGP or due to cold nuclear matter (CNM) effects. This can be studied by the so-called nuclear modification factor defined as:

$$R_{pPb}(p_T) = \frac{d^2 N_{pPb}/dy dp_T}{\langle N_{\text{Coll}} \rangle \cdot d^2 N/dy dp_T} \quad (6.2)$$

where the numerator represents the measured spectral yield in p–Pb collisions, the denominator includes $\langle N_{\text{Coll}} \rangle = 7.118$ [327], and the measured spectral yields in inelastic pp collisions.

Due to the fact that the spectra of K_S^0 , Λ and $\bar{\Lambda}$ in inelastic pp collisions are not available at $\sqrt{s_{\text{NN}}} = 8.16$ TeV, the reference pp spectra needs to be interpolated using

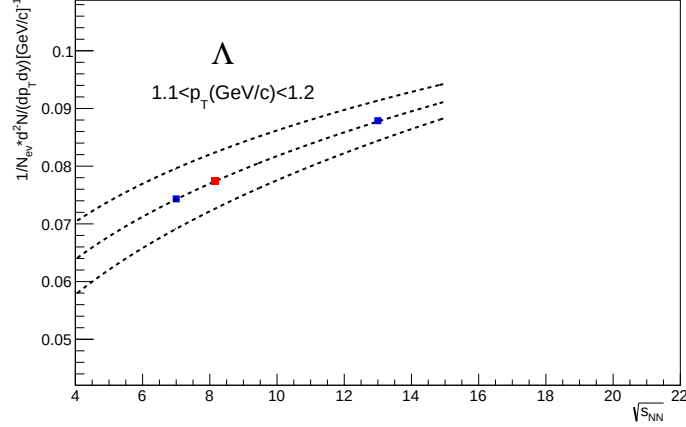


Figure 6.8: Fit to the dependence on $\sqrt{s_{\text{NN}}}$ of the pp yields for $\Lambda + \bar{\Lambda}$ for a p_{T} bin. The fitted function is used to interpolate the value at 8.16 TeV (red point on the curve).

the available data at 7 TeV and 13 TeV. The procedure used to obtain the reference spectral yields is described in [328]. First of all, for any bin in p_{T} the dependence on $\sqrt{s_{\text{NN}}}$ of the yields in the reference spectra is fitted to a power-law function $\beta \sqrt{s_{\text{NN}}}^{\alpha}$, as in Figures 6.7 and 6.8, for K_{S}^0 and Λ respectively, for a particular p_{T} bin. The obtained function is then used to interpolate the yield value at $\sqrt{s_{\text{NN}}} = 8.16$ TeV, represented by the red squares in the plot. The interpolated spectra for K_{S}^0 are shown, in Fig 6.9 (similarly obtained for Λ), together with the reference spectra at 7 and 13 TeV used for the interpolation. Systematics on the constructed reference spectra are then estimated.

To this end, the difference in the interpolated yields at 8.16 TeV obtained when fitting the central values of the 7 and 13 TeV spectra (central dotted black line in Figures 6.7 and 6.8) and the ones obtained when the central values are moved $\pm$ the relative systematic errors (up and down dotted lines in Figures 6.7 and 6.8) is considered.

The comparison of interpolated pp spectra at 8.16 TeV and p-Pb spectra at 8.16 TeV is shown in Figure 6.10. Since the p_{T} bins used in the reference pp spectra at 7

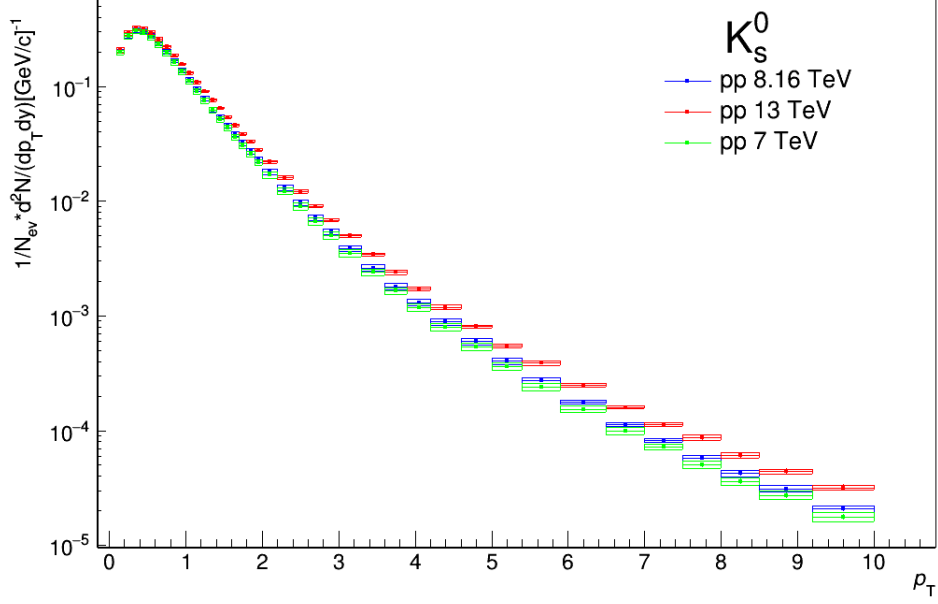


Figure 6.9: Reference pp spectra at 13 TeV (red marker) and 7 TeV (green marker) and interpolated spectra at 8.16 TeV (blue marker) for K_S^0 .

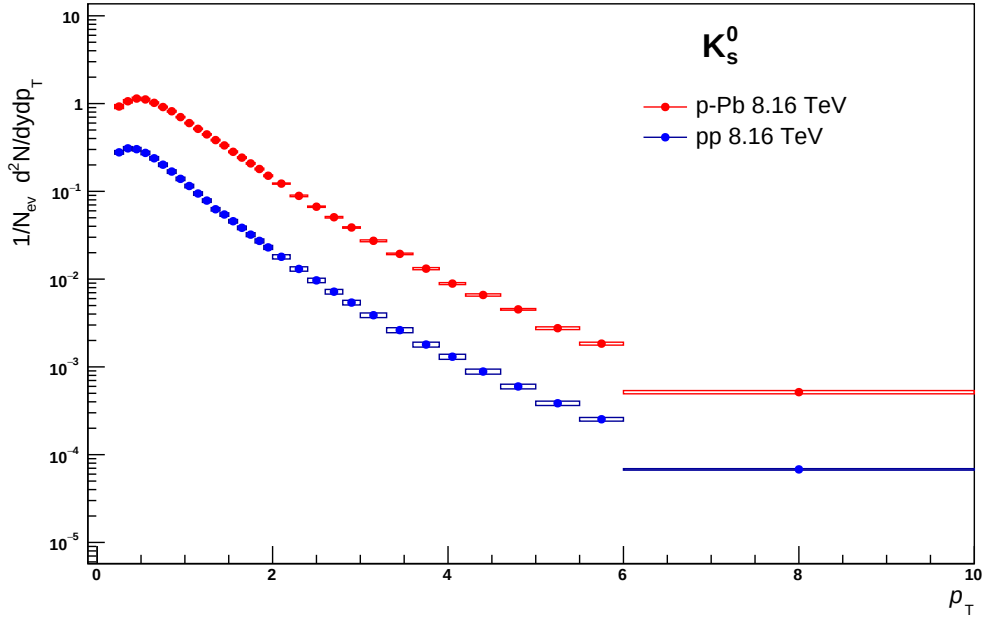


Figure 6.10: p-Pb spectra at 8.16 TeV (red marker) and re-binned interpolated pp spectra (blue marker) for K_S^0 .

6.5. RATIO OF RESONANCE PARTICLES TO GROUND-STATE PARTICLES

and 13 TeV were different from the p_T bins used in this analysis for p-Pb spectra so re-binning was done and the spectra for interpolated pp spectra at 8.16 TeV is shown in Figure 6.10 .

The nuclear modification factor is then calculated through Equation 6.2. The nuclear modification factor in p-Pb collisions for K_S^0 (red markers) and Λ (blue markers) can be found in Figure 6.11. In this figure, the bars and boxes represent the statistical and systematic uncertainties respectively.

The nuclear modification factor in p-Pb collisions is consistent with unity at high p_T , showing that the strong suppression observed in Pb-Pb is not due to Cold Nuclear Matter (CNM) effects but rather due to final state partonic energy loss in the hot and dense QGP created in Pb-Pb collisions [314]. This way we can disentangle the cold nuclear matter effects from final state effects.

The R_{pPb} measurements for K_S^0 and Λ shows a greater enhancement for Λ which is a baryon, at intermediate p_T . This enhancement can be described by Cronin effect. This effect is a generic term used for the observed broadening of transverse momentum spectra at intermediate p_T in p-Pb collisions as compared to pp. It is generally attributed to the influence of the multiple parton scattering through matter prior to the hard scattering that produces the observed high p_T hadron.

Therefore, the nuclear modification factor can be used to study the effect of matter on particle production. The differences between baryons and mesons is believed to be due hadrons production through quark coalescence at intermediate p_T .

6.5 Ratio of Resonance particles to Ground-state particles

Resonances are short lived particle having lifetimes $\sim 10^{-23}$ s. Due to their short lifetimes, they are useful tools to understand the mechanisms of particle production

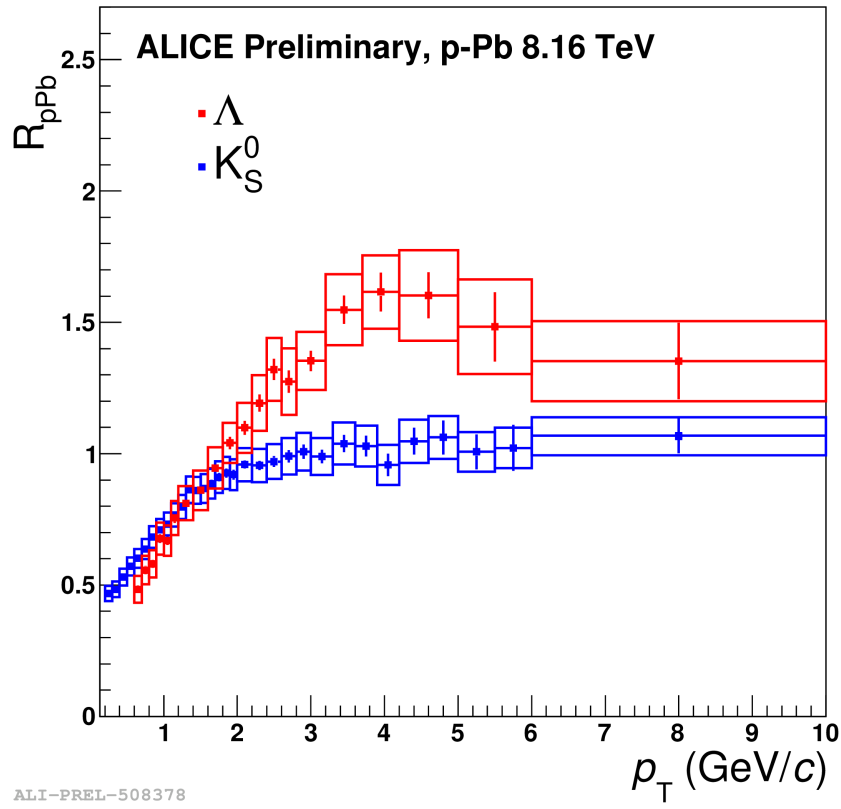


Figure 6.11: Nuclear Modification Factor, R_{pPb} for K_S^0 and Λ as a function of p_T

6.5. RATIO OF RESONANCE PARTICLES TO GROUND-STATE PARTICLES

and properties of the hadronic phase formed in heavy-ion collisions. If the time interval between the chemical and kinetic freezeouts is comparable with the lifetime of the resonance, the hadronic phase has enough time for scattering of the daughter particles. As a result, the yield of resonances might be modified due to in-medium effects [315]. So, measurements of various resonances with lifetimes from ~ 1 fm/ c to ~ 50 fm/ c enable us to investigate the role of the re-scattering and regeneration processes in the hadronic medium. Production properties of hadronic resonances and their comparison with those of stable hadrons are useful tools to understand the properties of the strongly interacting matter formed in relativistic nucleus-nucleus collisions. The rescattering process suppresses the observed yields with respect to what is produced at chemical freeze-out while regeneration enhances the yield of resonances [316]. Thus, the study of resonance particles can provide a hint for finite lifetime of the hadronic phase (resonance to stable hadron yield ratios).

Figure 6.12, shows the yield ratios for K^{*0} (892) to K_S^0 and Λ (1520) to Λ . This can be observed from the figure that K^{*0} (892) to K_S^0 yield ratio shows suppression but Λ (1520) to Λ yield ratio do not show any suppression. The Λ (1520) to Λ yield ratios in p-Pb 8.16 TeV are compared with the ratios obtained in p-Pb 5.02 TeV and similar trend is observed in both the energies. This provides hint for the formation of hadronic phase and its non-zero and finite lifetime in p-Pb collision system. This lifetime increases with the increase in the size of the colliding system.

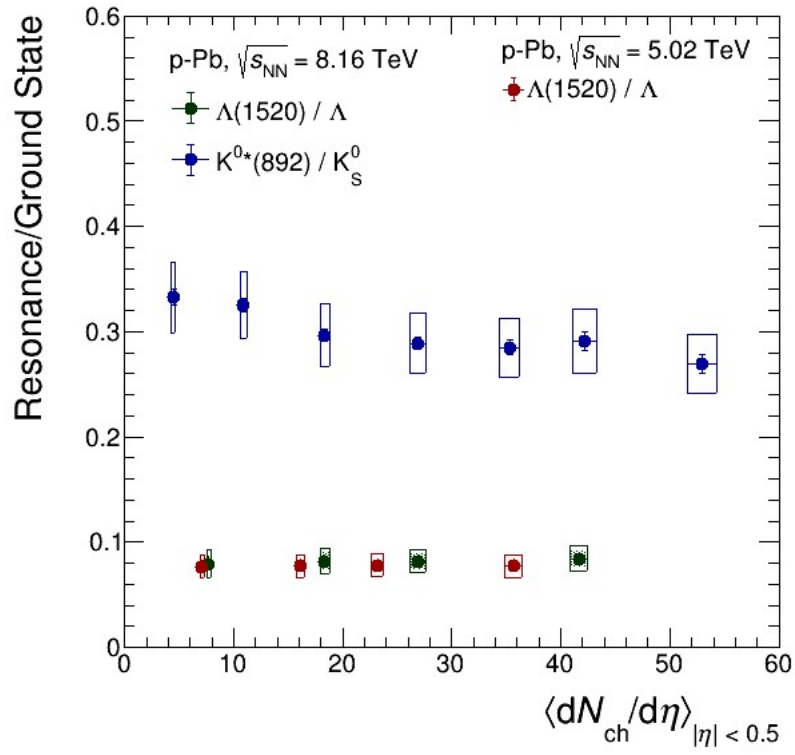


Figure 6.12: p_T integrated particle yield ratios for $K^{*0}(892)/K_S^0$ and $\Lambda(1520)/\Lambda$ as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta| < 0.5}$. Comparison of $\Lambda(1520)/\Lambda$ is also made with data points from p-Pb 5.02 TeV

Chapter 7

This chapter includes a brief summary of the results obtained for the present analysis and their physical interpretations. Moreover, the future perspectives of the current studies which are planned for Run3 and also for ALICE3 detector are discussed.

Chapter 7

Conclusions & future perspectives

In this work, the results on the production of strange and neutral particles (K_S^0 , Λ and $\bar{\Lambda}$) in p-Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV are presented. The measurement is carried out using the ALICE detector. Along with real data analysis, the Monte Carlo simulations are used to perform the efficiency, feed-down corrections and results comparisons.

The results obtained in the present work include the transverse momentum (p_T) spectra measurements in different multiplicity classes, the ratio of the p_T integrated yield of strange hadrons to pions, baryon to meson ratio (Λ/K_S^0) as a function of p_T and multiplicity and nuclear modification factor (R_{pPb}) calculations.

The main physics outcomes of the research work can be summarised as:

- We observed the hardening of p_T spectra of the strange particles from low multiplicity to high multiplicity, which can be quantised as the $\langle p_T \rangle$ increases as a function of multiplicity. The result is understood as an effect of the radial flow in the case of heavy-ion collisions.
- The ratios of the yields of the strange to non-strange hadrons across the different systems (pp, p-Pb, Pb-Pb and Xe-Xe) are found to have a smooth evolution

with charged particle multiplicity but there is no collision systems dependency found.

- The baryon over meson (Λ/K_S^0) ratio shows no significant change as a function of multiplicity suggesting that strangeness enhancement is not driven by mass nor it is a baryon/meson effect.
- Baryon enhancement: baryon over meson (Λ/K_S^0) ratio as a function of p_T shows a peak at intermediate p_T which is as an interplay of radial flow and parton recombination at intermediate p_T .
- Mass ordering of R_{pPb} : We observed mass ordering of the nuclear modification factor (R_{pPb}) at intermediate p_T which is quantitatively similar to that in Pb-Pb collisions. However, more studies are required to understand the exact phenomena of the results, since a clear baryon-meson splitting is also observed.
- Suppression is observed for $K^{*0} (892) / K_S^0$ but not for $\Lambda (1520) / \Lambda$ in the p-Pb collision system at 8.16 TeV, which provides hint for the formation of hadronic phase in p-Pb system and its finite lifetime.

7.0.1 Future perspectives

In recent studies in ALICE, it has been observed that small collision systems (pp, p-Pb) exhibit several striking signals of collective or QGP-like behaviour. The enhancement in the production of multi-strange hadrons, flow-like correlation results and baryon enhancement at intermediate p_T signaling the possible QGP-like property in the small systems has prompted the search of the phenomena which can mimic such effects in small systems. The next long-term goal of QGP research with the ALICE is to investigate and identifying the origin of these effects. By exploring this we can establish whether it arises from initial state effects, final state effects or a combination of the two.

The new decade of data starting with Run 3 (May 2022-2024), followed by Run 4 and Run 5 brings a much higher luminosity. The idea that a QGP may be formed in high multiplicity small systems (pp and p-Pb) is one of the most interesting and debatable topics in the field of particle physics. The ALICE experiment upgraded its two main central barrel detectors Inner Tracking System (ITS) and Time Projection Chamber (TPC) and forward detector Muon Forward Tracker (MFT) to cope up high statistics collisions data and precise tracking of particles down to the low p_T region to meet its physics goals.

The huge increase in statistics from Run 3, will allow us to explore to study multi-strange particle production in Pb-Pb down to lower multiplicities and thus extend the multiplicity range in pp collisions to overlap those of peripheral Pb-Pb data. Whether the multi-strange particle enhancements in pp collisions will continue to rise linearly past those from Pb-Pb data, will saturate at Pb-Pb multiplicities, or saturate below Pb-Pb multiplicities will give new insight into the mechanism of strangeness enhancement in small systems (more clear from Figure 7.1). In particular, one could study whether this phenomenon is due to QGP droplet forming in these systems or a completely new QCD process. With the new data, it can be possible to correlate the strangeness enhancement and elliptical flow property of the QGP behaviour

The study of charm particles using the direct decay channel via the strange particle will be extremely useful to investigate the property of charm in QGP. As charm quark is produced exclusively in initial hard scatterings such measurements will thus provide not only an insight into QGP medium but also the hadronisation from pp to A-A collisions. In particular, the Ξ_c^+ and Ω_c^0 measurements via the strange particles Ξ^+ and Ω^- studies respectively. The reconstruction of charmed hadrons decaying weakly to strange particles poses a significant challenge. The new Inner Tracking System ITS2 of ALICE can be able to track strange hadrons directly before

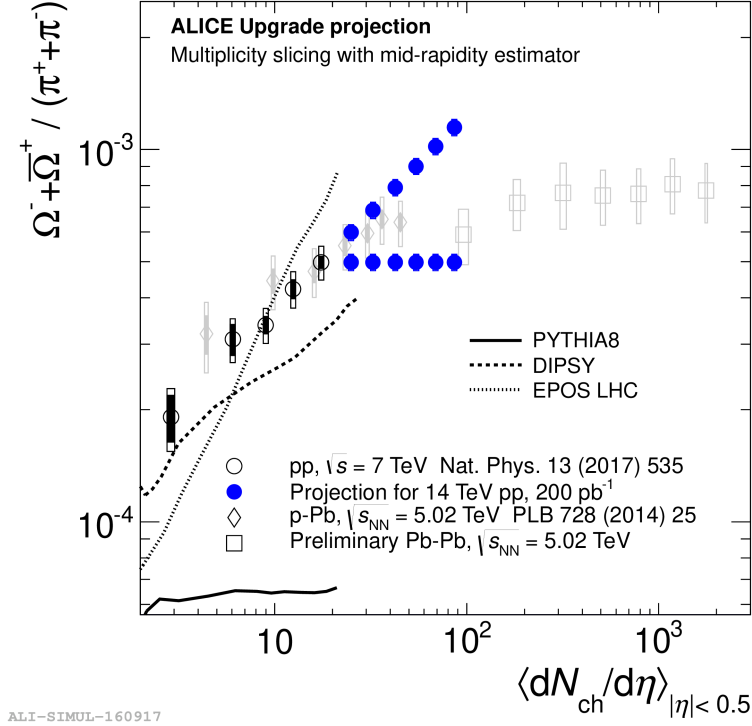


Figure 7.1: $(\Omega^- + \bar{\Omega}^+)/(\pi^+ + \pi^-)$ yields as a function of $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta|<0.5}$

the decay, using a novel experimental technique known as “Strangeness Tracking” [329, 330].

The large rapidity and transverse momentum acceptance of the proposed apparatus of ALICE3 experiment [331] will provide an ideal environment for the study of the hadronic physics. Using this experiment, measurements of multi-charm hadrons, such as the Ξ_{cc}^+ , Ξ_{cc}^{++} , Ω_{cc}^+ , Ω_{cc}^{++} , and exotic states such as the newly-discovered T_{cc}^+ [332, 333], would provide a direct window on hadron formation from a deconfined Quark-Gluon Plasma. These studies may shed light on the interplay between energy loss in QGP medium and onset of collective behaviour for charm quarks.

Appendix A

Additional Plots

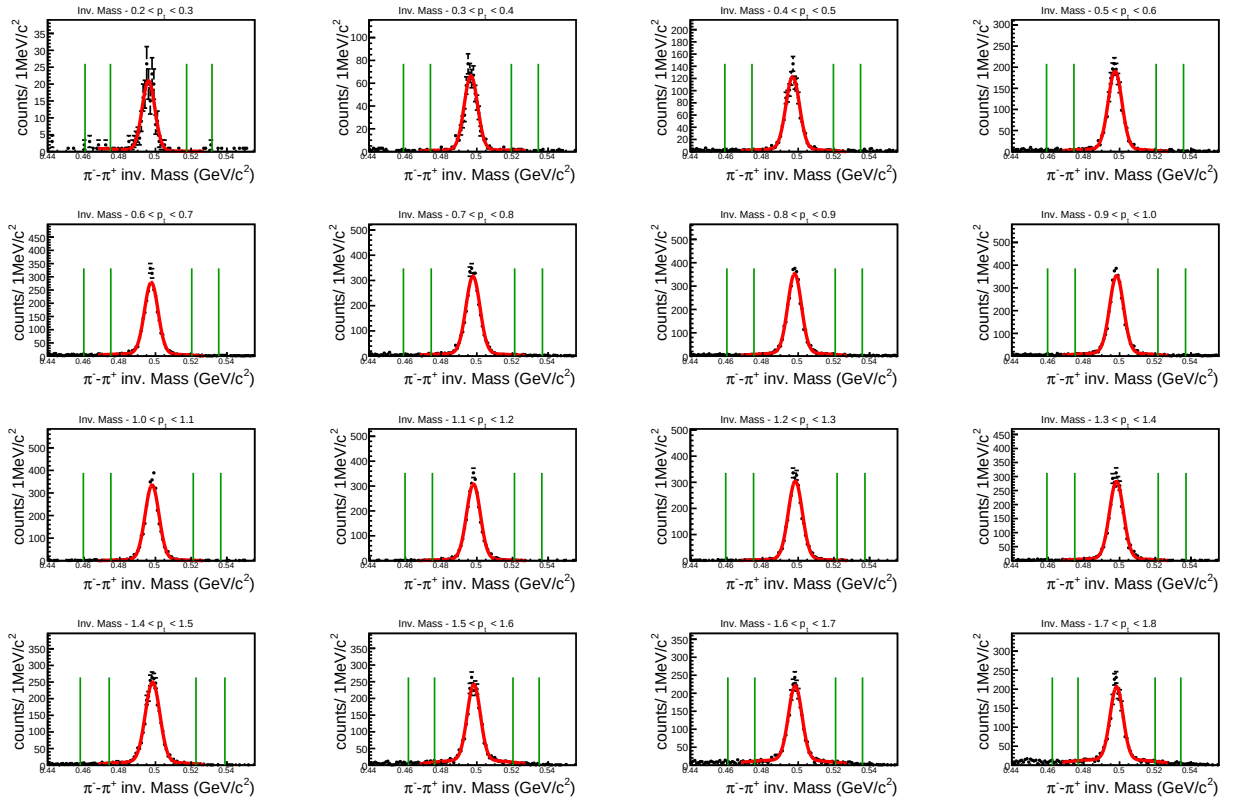


Figure A.1: invariant mass fits for K_S^0 in 0-1 % multiplicity bin in few p_T bins.

APPENDIX A. ADDITIONAL PLOTS

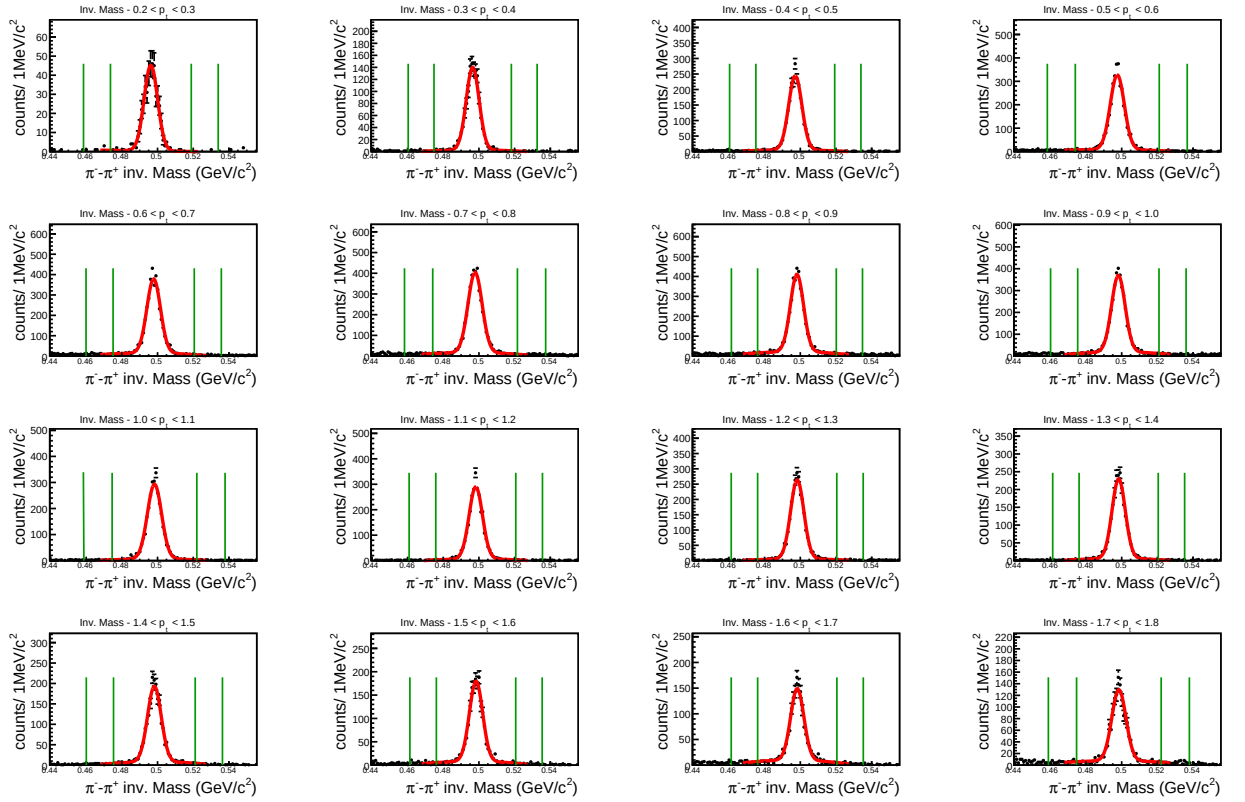


Figure A.2: invariant mass fits for K_s^0 in 80-100 % multiplicity bin in few p_T bins

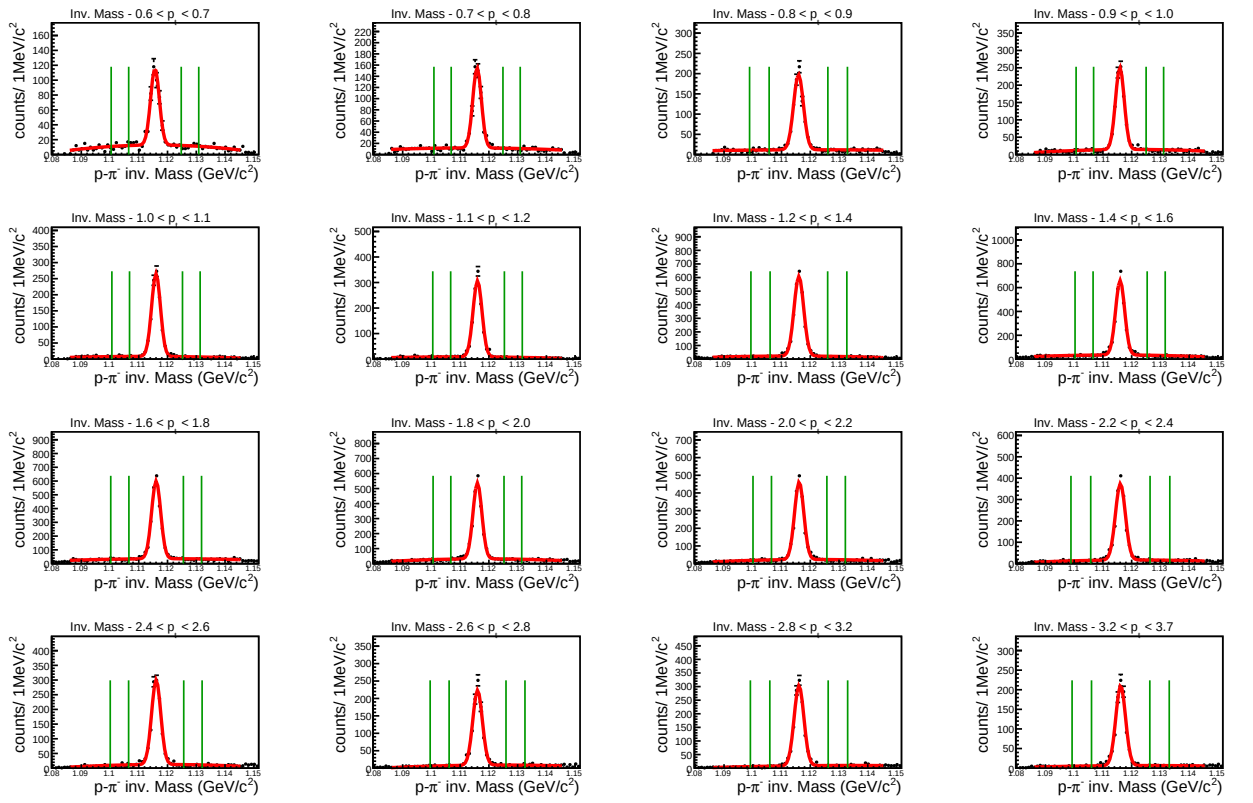


Figure A.3: invariant mass fits for Λ in 0-1 % multiplicity bin in few p_T bins.

APPENDIX A. ADDITIONAL PLOTS

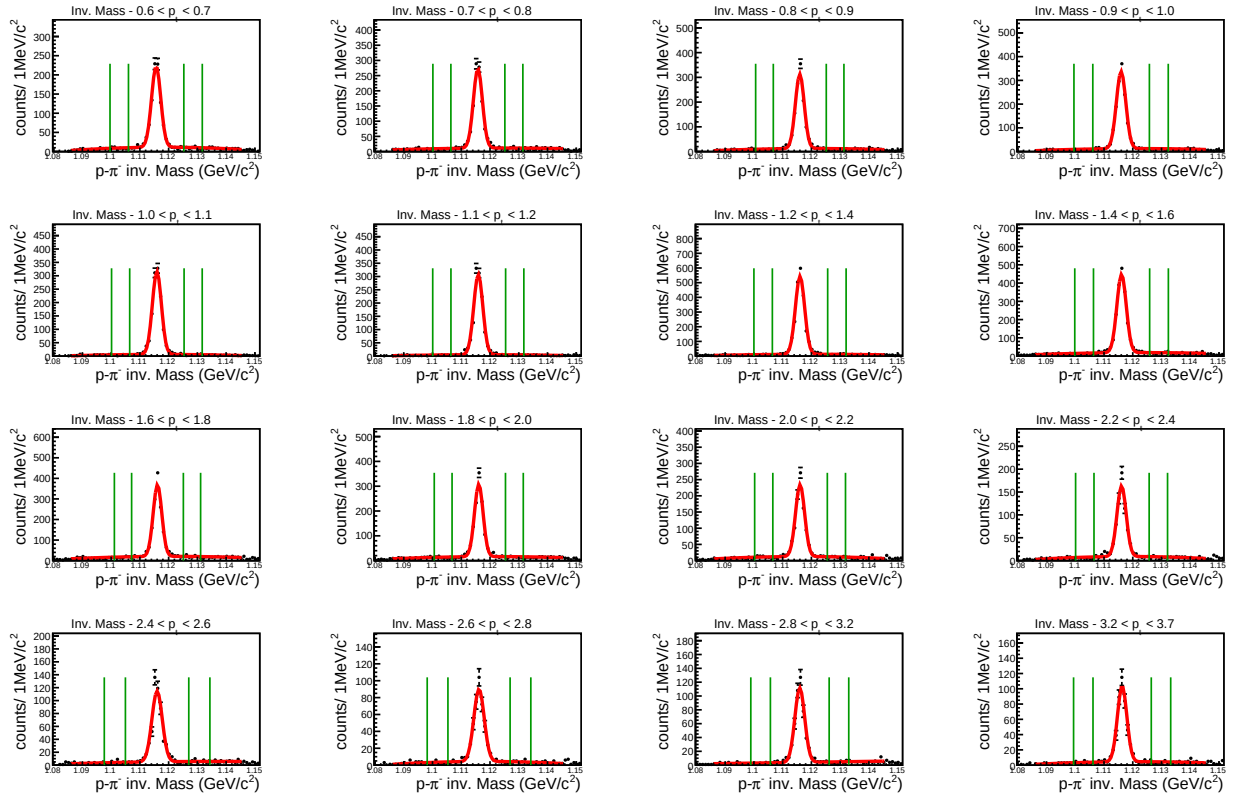


Figure A.4: invariant mass fits for Λ in 80-100 % multiplicity bin in few p_T bins

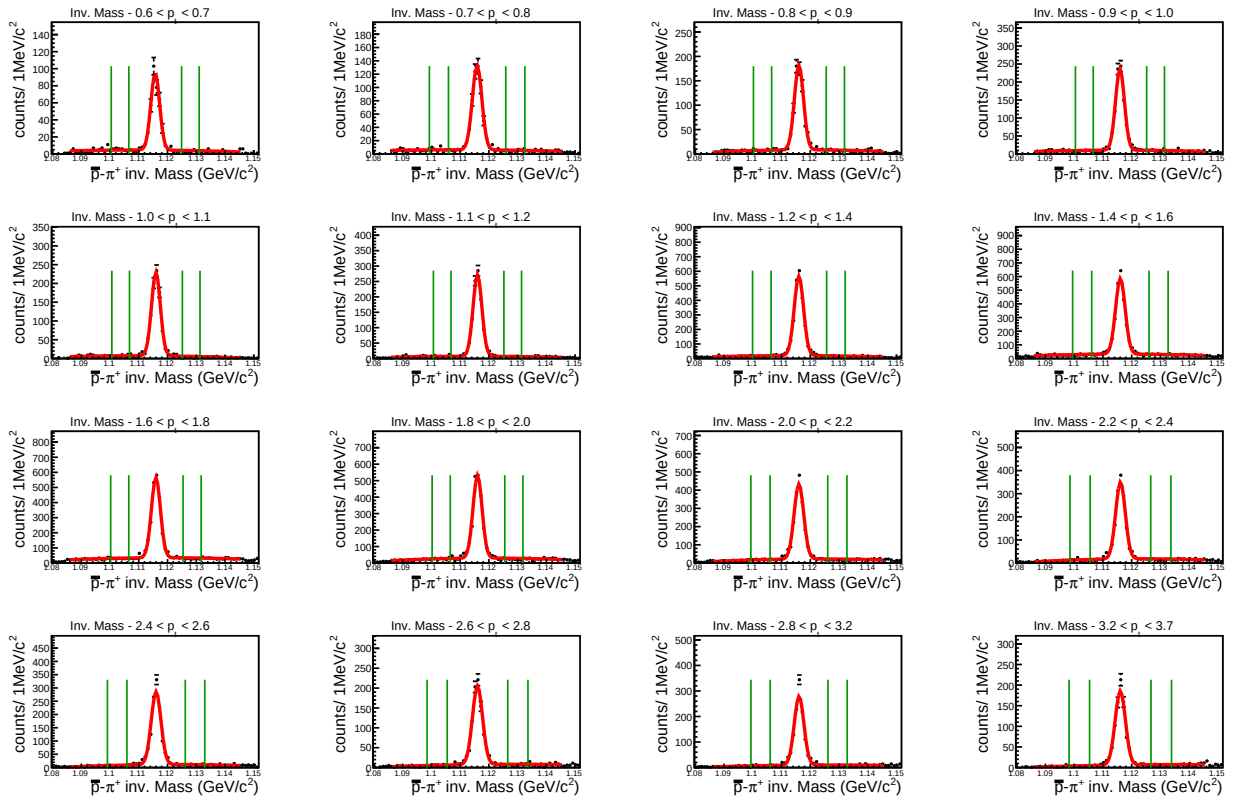


Figure A.5: invariant mass fits for $\bar{\Lambda}$ in 0-1 % multiplicity bin in few p_T bins.

APPENDIX A. ADDITIONAL PLOTS

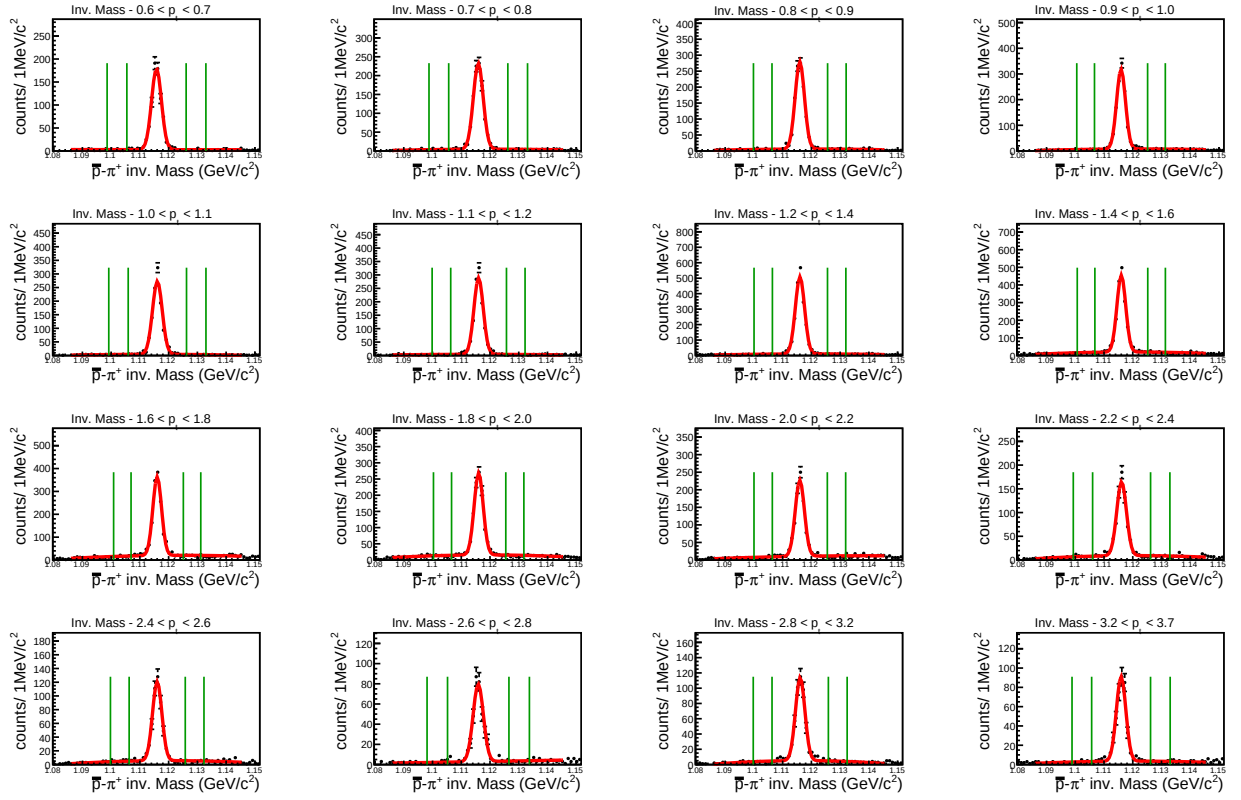


Figure A.6: invariant mass fits for $\bar{\Lambda}$ in 80-100 % multiplicity bin in few p_T bins

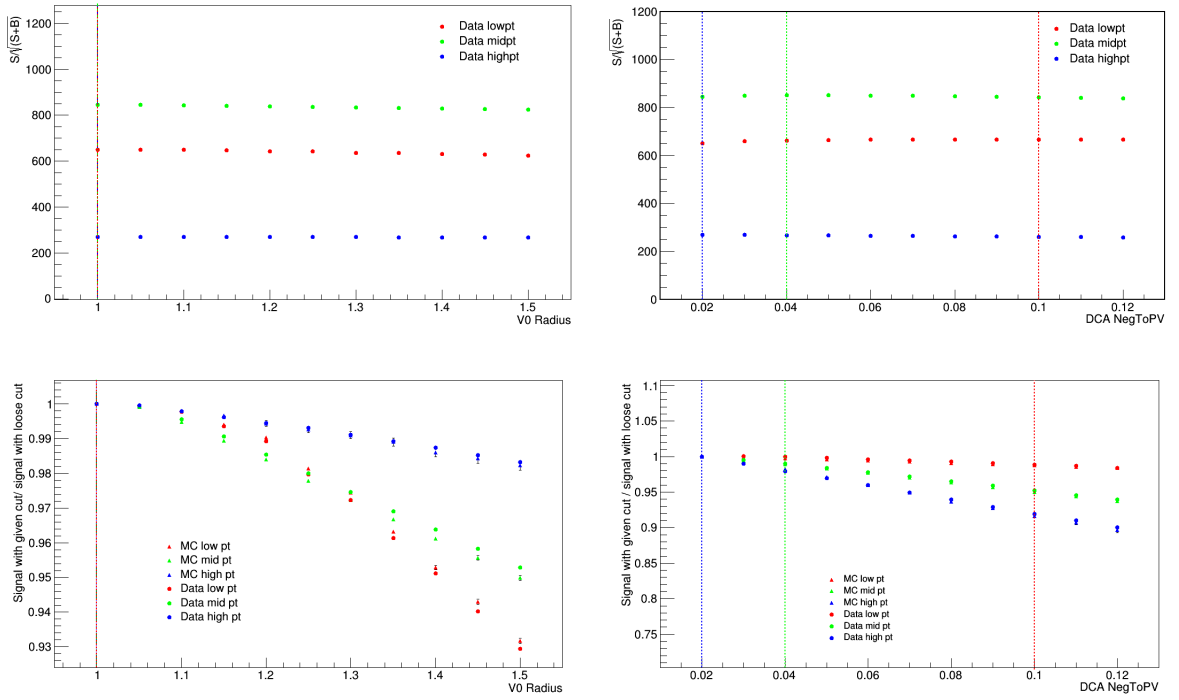


Figure A.7: Significance and signal losses for K_S^0 , with respect to loosest cut considered, in low, mid and high p_T bins for the variables V0Radius and DCA NegToPV.

APPENDIX A. ADDITIONAL PLOTS

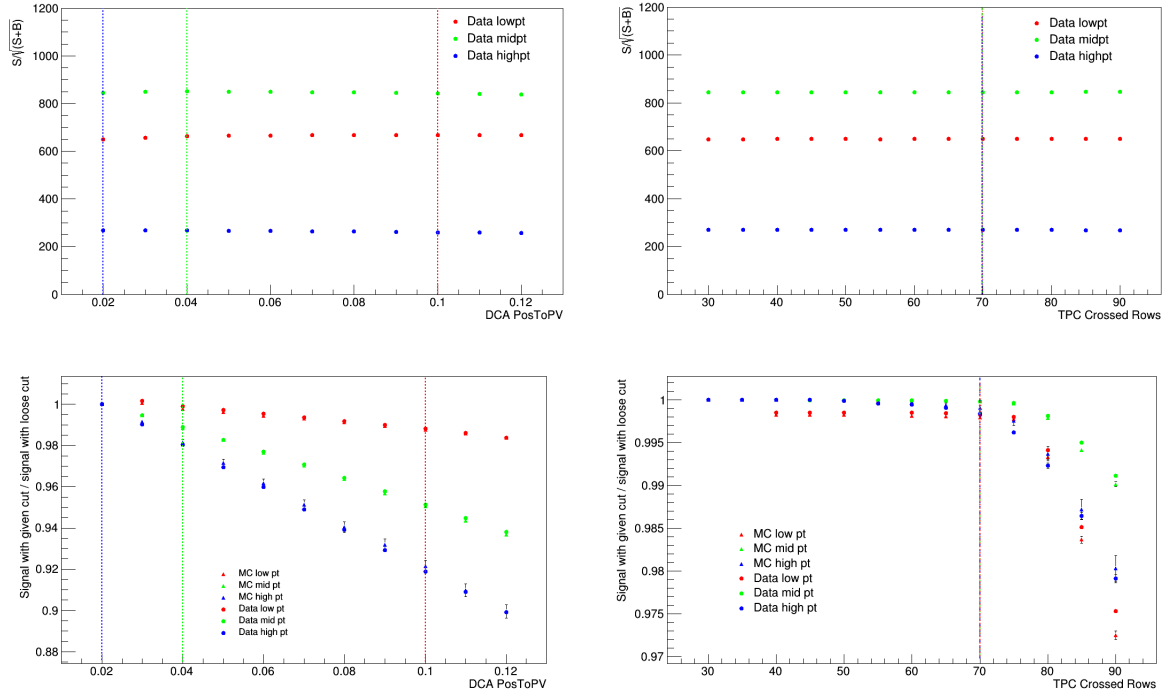


Figure A.8: Significance and signal losses for K_S^0 , with respect to loosest cut considered, in low, mid and high p_T bins for the variables DCA PosToPV and TPC Crossed Rows.

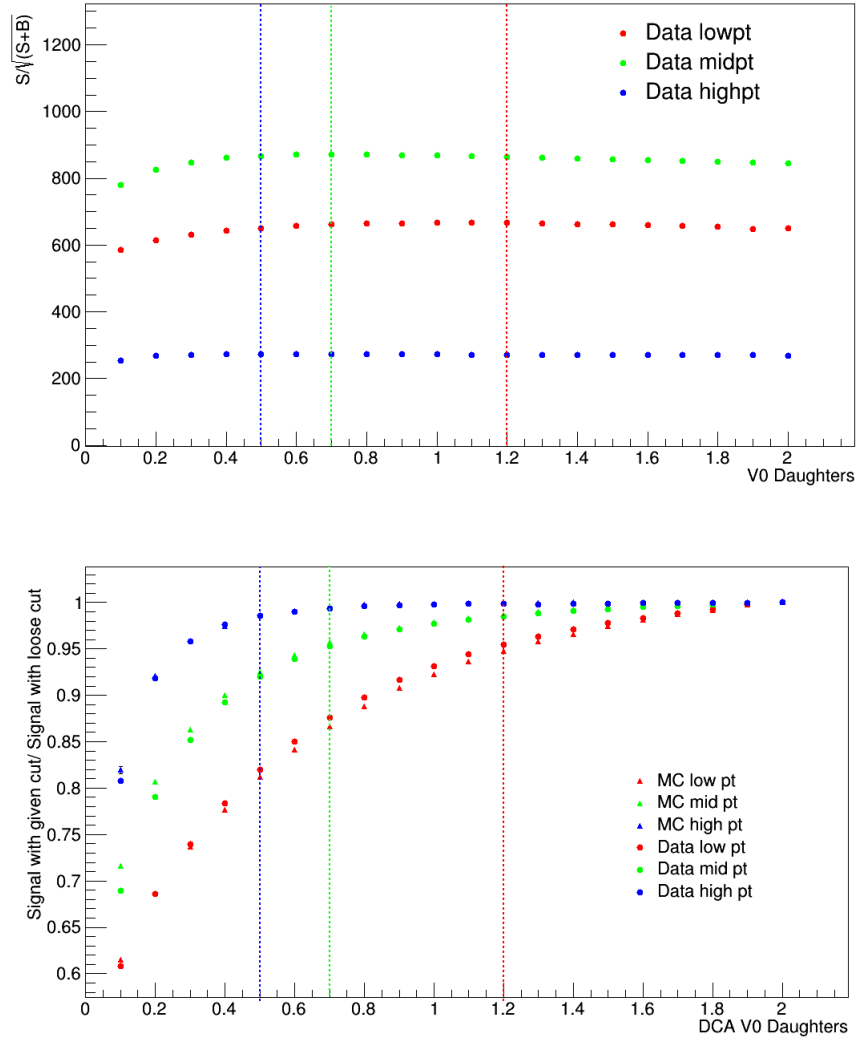


Figure A.9: Significance and signal losses for K_S^0 , with respect to loosest cut considered, in low, mid and high p_T bins for the variable DCA V0Daughters.

APPENDIX A. ADDITIONAL PLOTS

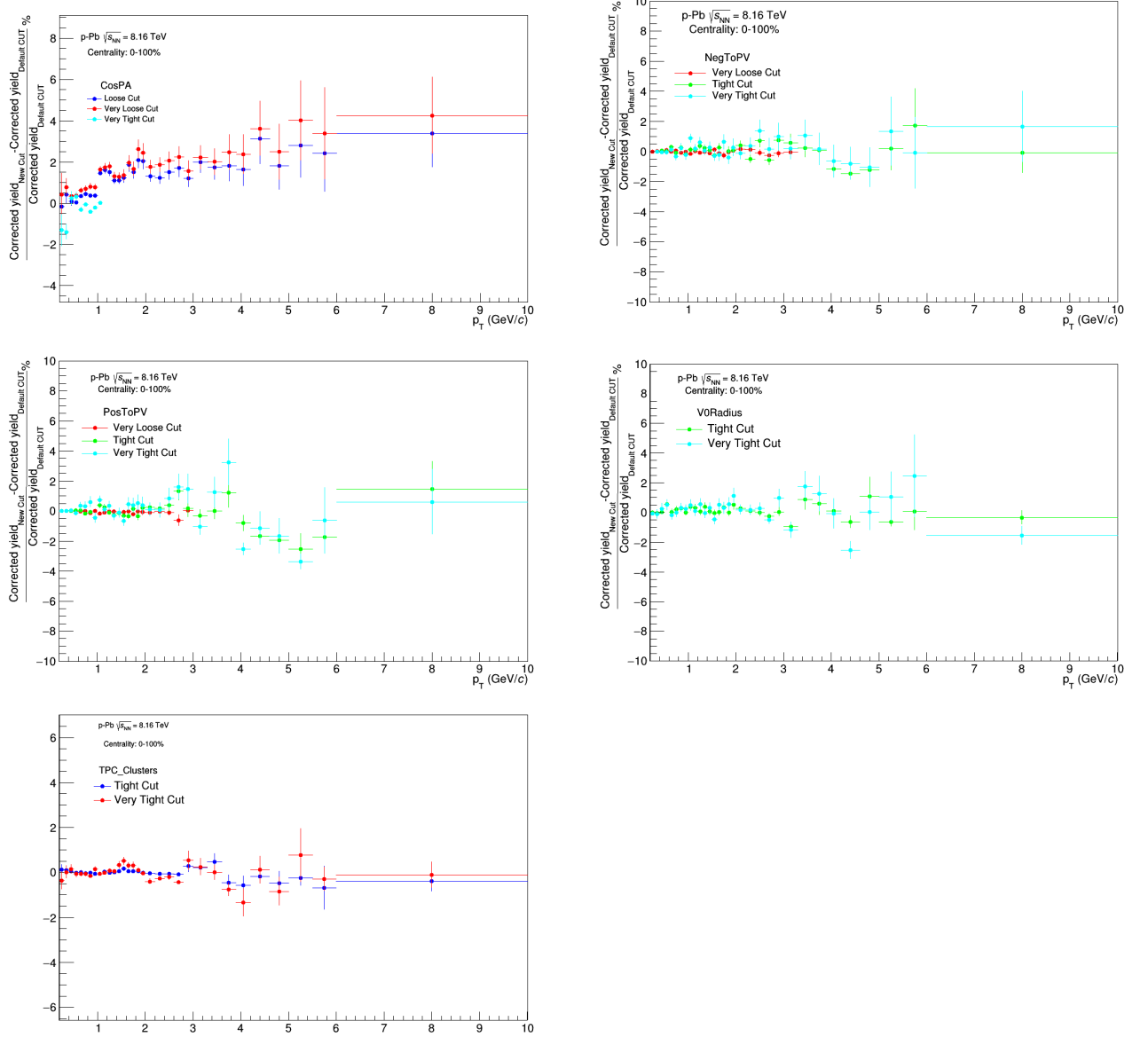


Figure A.10: Signal deviations seen in corrected K_S^0 spectra due to cut selection changes. Error bars are Roger-Barlow errors representing an acceptable deviation due to statistical fluctuations.

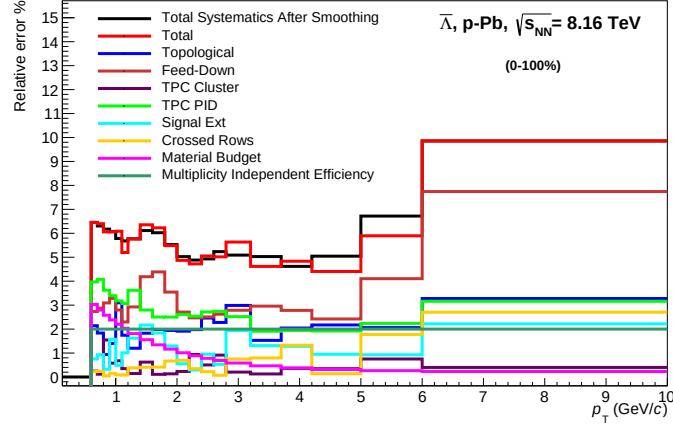


Figure A.11: Total Systematics obtained in 0-100% multiplicity class for $\bar{\Lambda}$

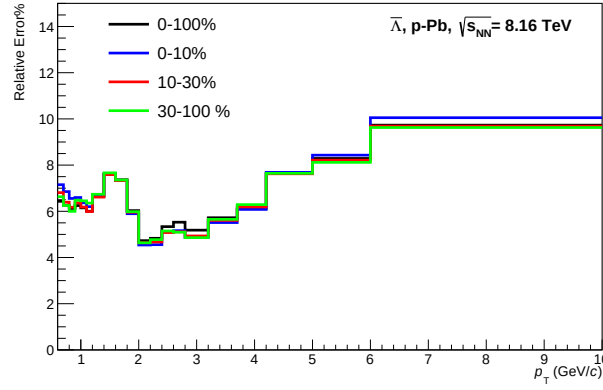


Figure A.12: Total systematic uncertainties obtained in three different multiplicity classes (0-10%, 10-30% and 30-100%) with roughly 1/3 of the MB statistics each for $\bar{\Lambda}$.

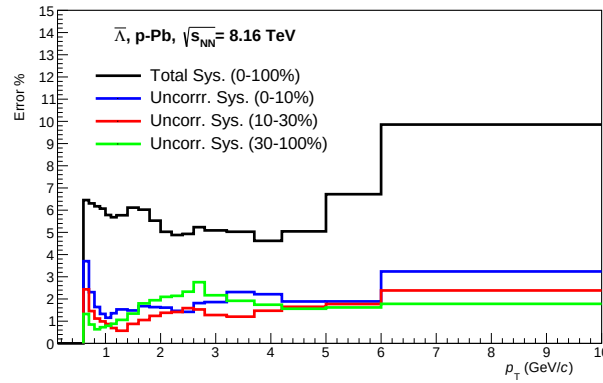
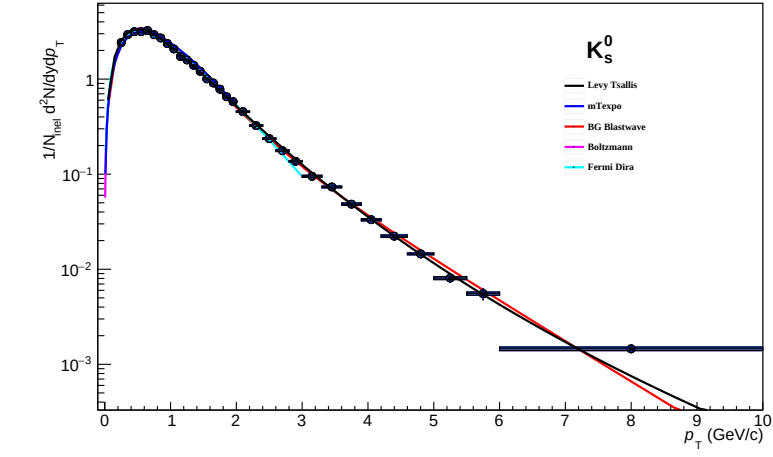
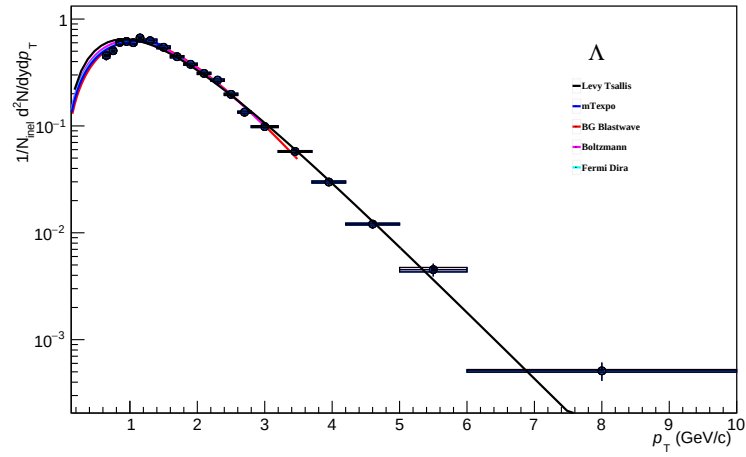


Figure A.13: The multiplicity-uncorrelated uncertainties in comparison to the total (MB) systematic uncertainty for $\bar{\Lambda}$.

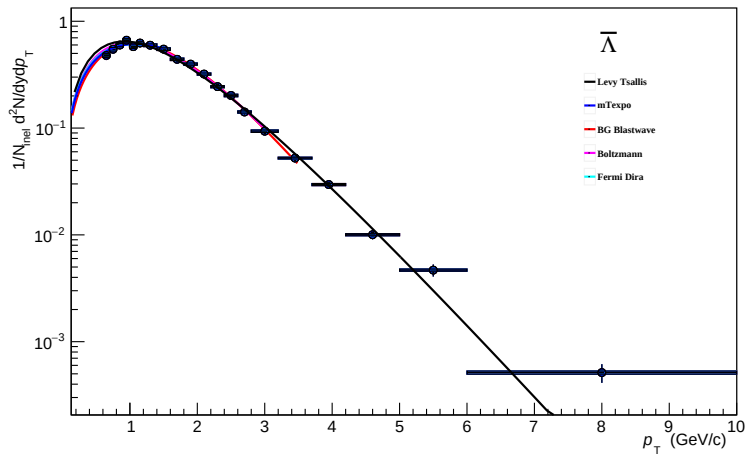
APPENDIX A. ADDITIONAL PLOTS



(a)

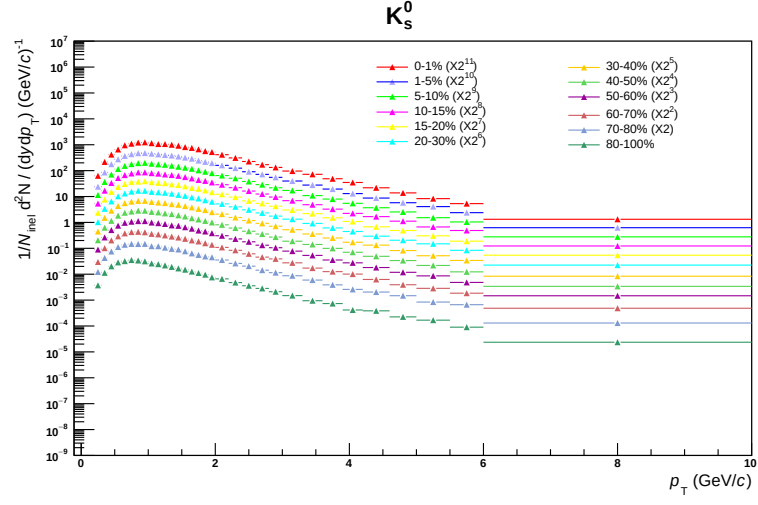


(b)

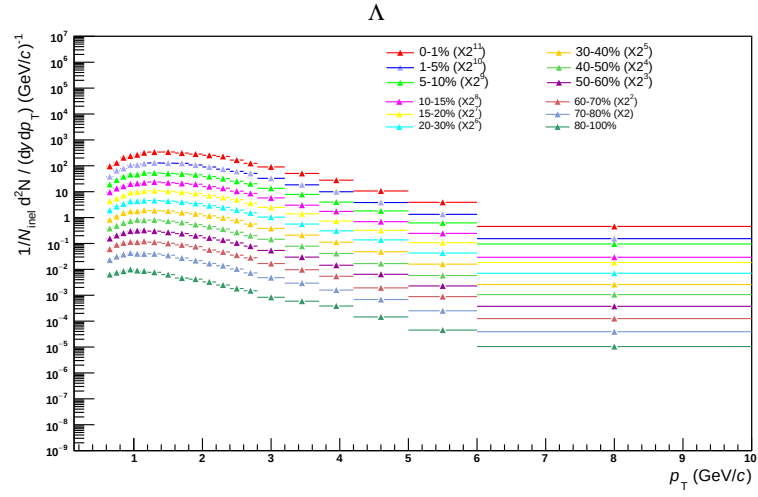


(c)

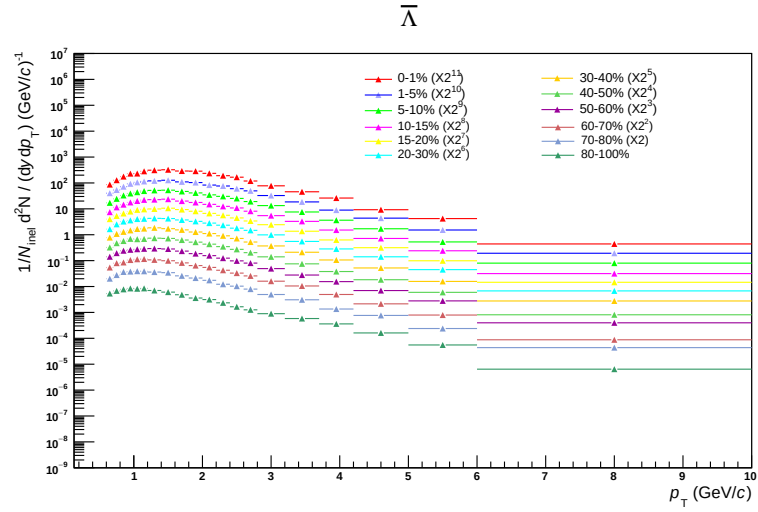
Figure A.14: (a) K_S^0 , (b) Λ and (c) $\bar{\Lambda}$ p_T spectra fitted with different functions to compute systematics due to extrapolation.



(a)



(b)



(c)

Figure A.15: (a) K_S^0 , (b) Λ and (c) $\bar{\Lambda}$ raw p_T spectra in different multiplicity classes

APPENDIX A. ADDITIONAL PLOTS

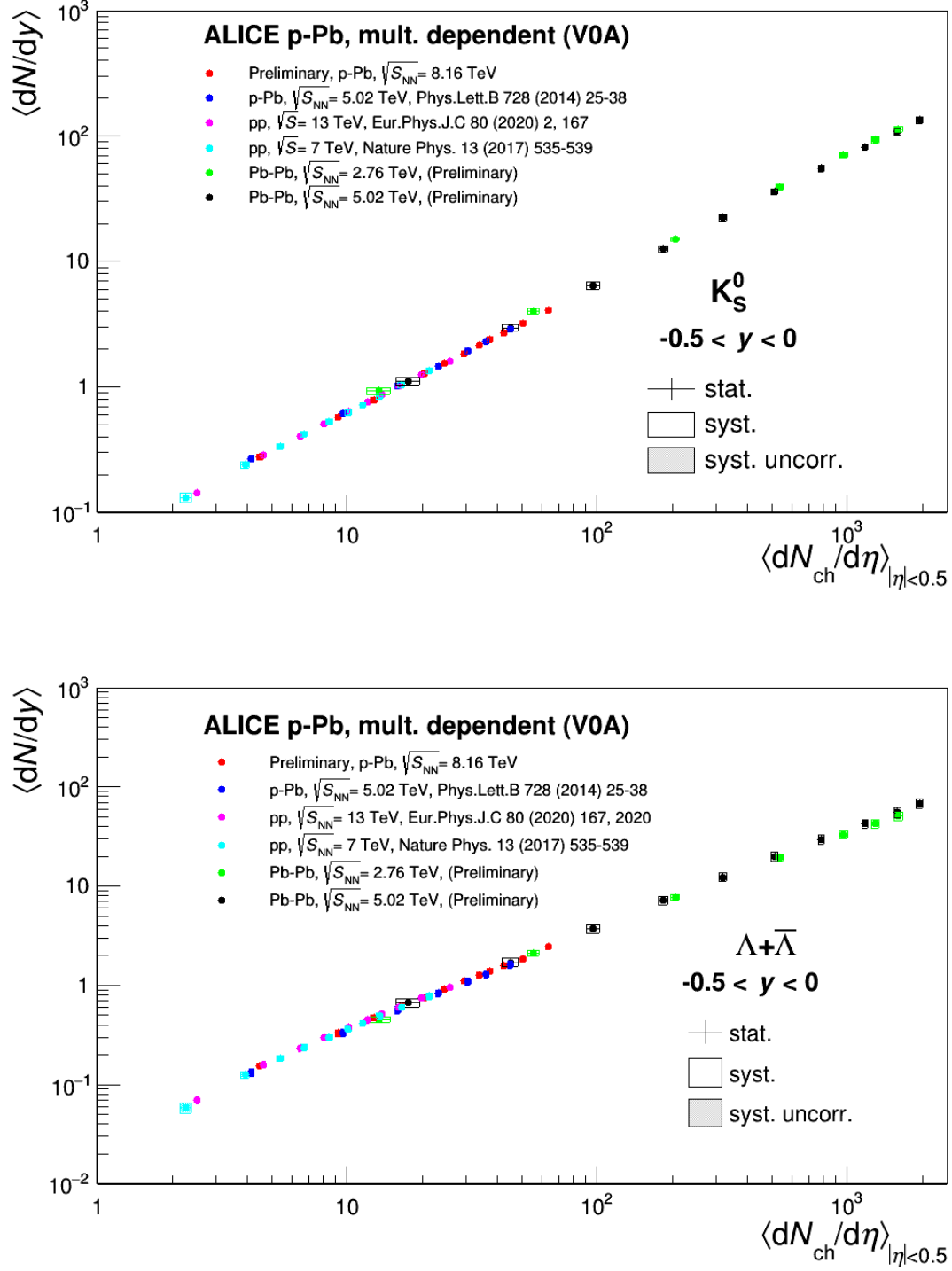


Figure A.16: p_T integrated yield per unit rapidity, dN/dy as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta|<0.5}$ for K_S^0 and $(\Lambda + \bar{\Lambda})$ including data points from Pb-Pb collisions

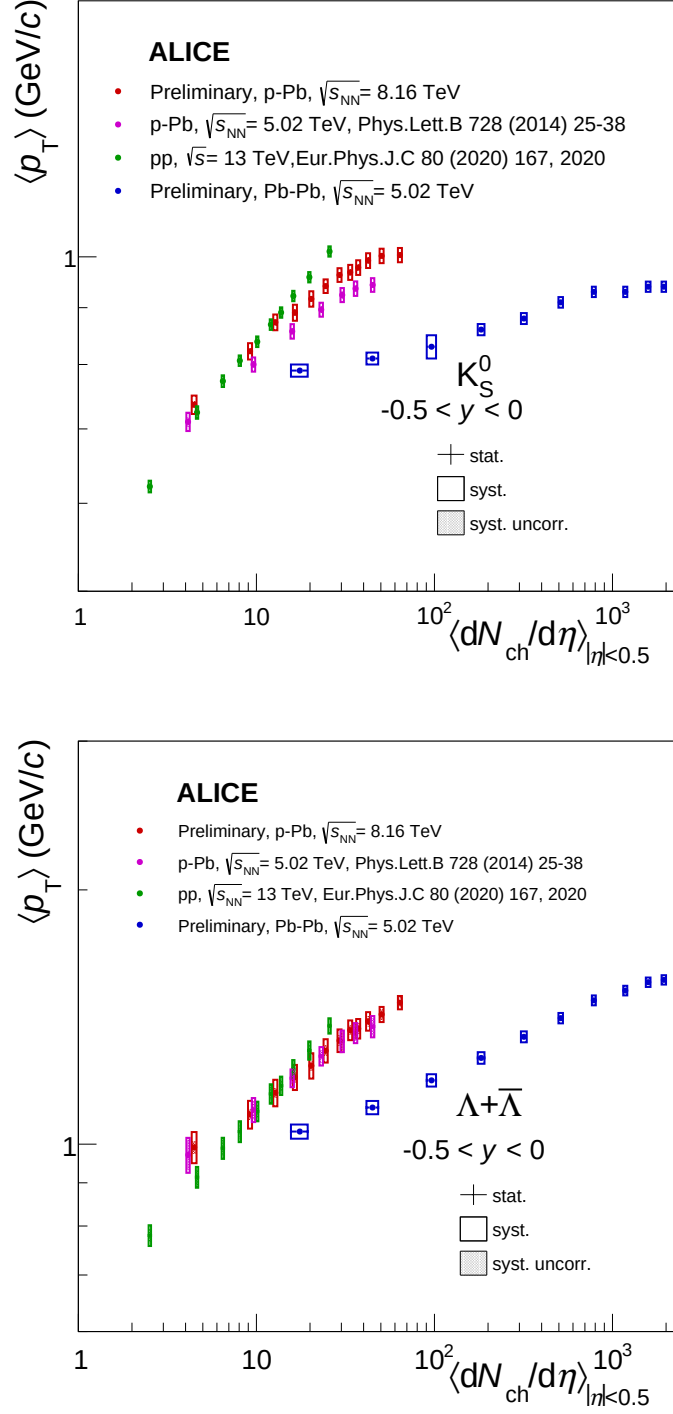


Figure A.17: Average transverse momentum, $\langle p_T \rangle$ as a function of $\langle dN_{ch}/d\eta \rangle_{|\eta|<0.5}$ for K_S^0 and $(\Lambda + \bar{\Lambda})$ including data points from Pb–Pb collisions

Appendix B

Lévy-Tsallis function

This function is used to describe the transverse momentum distribution of hadrons originating from a thermally equilibrated source. The Lévy-Tsallis function expresses the p_T distribution of hadrons as:

$$\frac{d^2N}{dydp_T} = \frac{(n-1)(n-2)}{nT_k[nT_k + m(n-2)]} \frac{dN}{dy} p_T \left(1 + \frac{m_T - m}{nT_k} \right)^{-n} \quad (\text{B.1})$$

where n is exponent, T_k is inverse slope parameter, m is particle mass. T_k is determined by the low p_T part, while n is by the tail. Also n and T_k can be highly correlated, especially if we don't have enough p_T range.

This equation assumes a Boltzmann-Gibbs-type exponential ($\propto e^{-p_T/T}$) at small p_T and a power-law at large p_T ($\propto T/(np_T)^n$), and thus can describe the soft, thermal, bulk of a hadronic system, and the particles from hard processes at large momenta.

Appendix C

Geometry and centrality in heavy-ion collisions

C.0.1 Geometry

Nuclei are very extended objects for all the scales of interest in high-energy physics. For this reason, the geometry of the collision plays a central role in the analysis and interpretation of the experimental results. In the centre-of-mass frame, owing to Lorentz contraction in the longitudinal direction, the two nuclei can be seen as two thin disks of transverse size $2R_A \simeq 2A^{1/3}$ fm. Some relevant quantities which are relevant to understand this concept are given below:

- The impact parameter which is the distance between the centres of the two colliding nuclei. The impact parameter characterizes the centrality of the collision. A central collision is one with small impact parameter in which the two nuclei collide almost head-on; a peripheral collision, on the contrary, is one with large impact parameter. A special case is that of ultraperipheral collisions in which the two nuclei interact only through the photons created by the large electromagnetic field of the ions [334].

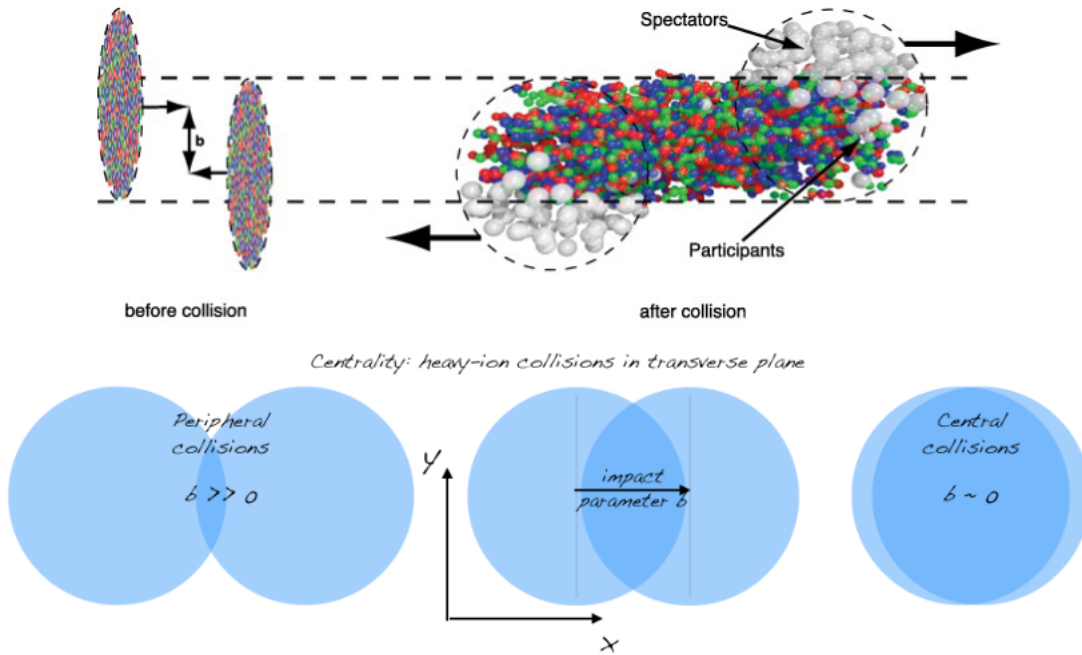


Figure C.1: Top Left: Two heavy ions before the collision with the impact parameter b . Top Right: The spectators remain unaffected while in the participant zone, particle production takes place - UrQMD calculation [335]. Bottom: An illustration of the often used nomenclature. Three collisions in the transverse plane are shown (from left to right): peripheral, semi-central, and central

- The number of participant nucleons, N_{part} , within the colliding nuclei which is the total number of protons and neutrons which take part in the collision. The rest are sometimes called spectators and continue travelling almost unaffected; they are normally measured in forward detectors such as the Zero-Degree Calorimeter (ZDC) to help in fixing the centrality of the collision.
- The number of collisions, N_{coll} , is the total number of incoherent nucleon–nucleon collisions.

Figure C.1 shows two heavy ions before and after a collision - a visualization of UrQMD calculation [335]. The number of collisions and the number of nucleons that collided depends on the impact parameter b of the collision. This figure also shows sketches of few selected collision “centralities” in the transverse plane (perpendicular

to beam axis) - peripheral collisions where impact parameter is large, semi-central, and central with b close to zero. Note, the impact parameter together with the axis along the beam define the so-called reaction plane. Centrality of the heavy-ion collisions is often expressed in the percentile of the total inelastic cross-section. Typically a peripheral collision is considered beyond 80% of the cross-section (for example in Pb-Pb collisions) where the impact parameter approaches two times the radius of the nuclei. While collisions referred to as “central” are at typically 0-10% (or 0-5%) of the cross-section. It is worthwhile to note, that particle production can take place for collisions with $b > 2R$ where it is dominated by electromagnetic processes. These so-called ultra-peripheral collisions (UPC) are on their own an interesting field of research and among other interesting aspect provide unique insight into the structure of nuclei. For a recent review of UPC see Ref.[336]

C.0.2 Centrality

In heavy-ion collisions, the event centrality is of utmost importance to understand the underlying physics of the collision. The event centrality is related to the impact parameter, defined as the distance between the centroids of the two colliding nuclei in a plane transverse to the beam axis, of the collision. The impact parameter tells about the overlap volume of the two nuclei. This overlap volume determines the geometrical parameters, like number of participant nucleons (N_{part}), number of spectator nucleons (N_{spec}), and the number of binary collisions (N_{coll}). The impact parameter can not be determined experimentally. However, the global observables, like total charged particles (N_{ch}), transverse energy (E_T), or energy deposited in ZDC (E_{ZDC}), and so forth, are related to this geometrical quantity. By combining the experimental observables with simulation, one can estimate the impact parameter and hence the centrality of the event class. The centrality is expressed as the percentile (C) of the total hadronic interaction crosssection corresponding to the charged

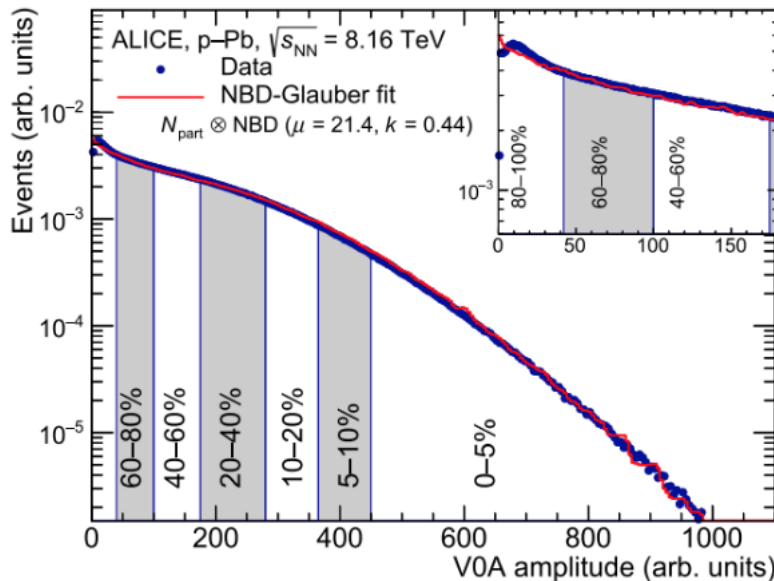


Figure C.2: Distribution of the sum of amplitudes in the V0A scintillators for p-Pb collisions at $\sqrt{s_{NN}}$ 8.16 TeV. The distribution is fitted with the NBD-Glauber fit shown as a line. The inset shows a zoom of the most peripheral region.

particle multiplicity above certain threshold and is given by

$$c = \frac{1}{\sigma_{AA}} \int_0^b \frac{d\sigma}{db'} db' \quad (C.1)$$

In Equation C.1, σ_{AA} is the total nuclear interaction cross section of A+A collision. Assuming constant luminosity, the cross section can be replaced by the number of observed events after the trigger efficiency correction. But at very high energy, when these two nuclei pass by each other, there is a large QED cross section because of the electromagnetic field [337, 338]. This QED cross section is much larger than the hadronic cross section and this contaminates the most peripheral events. That is why the centrality determination is restricted to some percentile where the QED contribution is negligible. The fraction of hadronic events excluded by such cut as well as the trigger efficiency can be estimated by using a Glauber model simulation. For a given impact parameter, the N_{part} and N_{coll} can be estimated by Glauber Monte

Carlo method. The parametrized Negative Binomial Distribution (NBD) can be used to describe the nucleon-nucleon collisions. For heavy-ion collisions, N_{part} and N_{coll} are used to generate the number of charged particles by incorporating two-component model in the following way:

$$N_{ancestors} = f \times N_{part} + (1 - f) \times N_{coll} \quad (C.2)$$

This $N_{ancestors}$ refers to the “independent emitting source”. The two-component model given in Equation C.1 incorporates the soft and hard interactions. Soft process is related to the N_{part} and hard process is related to N_{coll} . The functional form of NBD distribution is given by

$$P(\mu, k, n) = \frac{\Gamma(n + k)}{\Gamma(n + 1)\Gamma(k)} \cdot \frac{(\mu/k)^n}{(\mu/k + 1)^{n+k}} \quad (C.3)$$

Equation C.3 represents the probability of measuring n hits per ancestor. Here, μ represents the mean multiplicity per ancestor and k controls the width of the distribution. In $p + p(\bar{p})$ collision, a Negative Binomial Distribution (NBD) with a fixed value of μ and k well describes the charged multiplicity data for most of multiplicity range though requiring a second NBD ingredient to well describe the tail of the distribution [339–341]. The charged particle multiplicity for nucleus-nucleus collisions with a given impact parameter is generated by sampling $N_{ancestors}$ times the $p + p$ multiplicity, which is generated by using NBD. Finally, a χ^2 minimisation is done by fitting the Glauber Monte Carlo generated multiplicity and the charged particle multiplicity obtained from the collision data. The χ^2 minimization will give us the value of f , μ , and k . This gives a connection between an experimental observable and a Glauber Monte Carlo. From this one can have access to N_{part} and N_{coll} for a given class of centrality by NBD-Glauber fit. For example, the centrality determination in ALICE

APPENDIX C. GEOMETRY AND CENTRALITY IN HEAVY-ION COLLISIONS

using VZERO (V0) amplitude is given in Figure C.2. The two-component model is fitted with the V0 amplitude in Figure C.2 to find out the N_{part} and N_{coll} values for a corresponding centrality [337].

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List of Publications/Conferences

A. List of presentations/participation in conferences/schools:

1. Nov., 2017, SERC School on EHEP, NISER, Orissa, India
2. Sept., 2018, Asia-Europe-Pacific School for High Energy Physics, Quy Nhon, Vietnam
3. Dec., 2018, [Poster] DAE-BRNS Symposium for HEP, IIT Madras, Chennai, India. Title: “Strangeness production in p–Pb at 8.16 TeV.”
4. Aug., 2019, 14th CERN-Fermilab Hadron Collider Physics Summer School, Geneva, Switzerland.
5. Feb., 2019, [Talk] International Workshop on Forward and Jet Physics at LHC, Bose Institute, India. Title: “ V^0 production in p–Pb collisions with ALICE at the LHC.”
6. Nov., 2019, [Poster] QUARK MATTER 2019, Wuhan, China
Title: “ K_S^0 and Λ production in p–Pb collision system at 8.16 TeV.”
7. Dec., 2019, [Talk] JKScience Congress, University of Jammu.
Title: “Studying strangeness production from small to large systems with ALICE at the LHC.”.
8. Dec., 2020, [Talk] DAE-BRNS Symposium for HEP, NISER, Orissa, India (online). Title: “Energy dependence of strangeness production in small collision systems with ALICE.”
9. Feb., 2021, [Invited Talk] International Master Class (online).
Title: “Strangeness enhancement as a signature of QGP.”

10. July, 2021, [Talk] 19th International Conference on Hadron Spectroscopy and Structure (HADRON 2021), Universidad Nacional Autónoma de México (UNAM, Mexico) (online).
Title: “Recent results on strange and multi-strange particle production with ALICE at the LHC.”
11. April, 2022, [Poster] XXIXth International Conference on Ultra-relativistic Nucleus-Nucleus Collisions (Quark Matter 2022), Krakow, Poland (Hybrid). Title: “Strange particle production in p–Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV with ALICE at the LHC.”
12. Dec. 2022, [Talk] 66th DAE Symposium on Nuclear Physics, Cotton University, Guwahati, Assam, India. Title: “Multiplicity dependent studies for strangeness production with ALICE.”
13. Feb. 2023, [Poster] 8th International Conference on Physics and Astrophysics of Quark-Gluon Plasma, Puri, Orissa, India.
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B. Work published/preprints as conference proceedings:

1. M. Sharma [ALICE], Strangeness production in p–Pb collisions at 8.16 TeV, Springer Proc. Phys. **261** (2021), 1019-1022 doi:10.1007/978-981-33-4408-2_148 [arXiv:1911.04845 [hep-ex]].
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1. S. Acharya *et al.* [ALICE], “Production of light-flavor hadrons in pp collisions at $\sqrt{s} = 7$ and $\sqrt{s} = 13$ TeV,” Eur. Phys. J. C **81** (2021) no.3, 256 doi:10.1140/epjc/s10052-020-08690-5 [arXiv:2005.11120 [nucl-ex]].
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¹I am a co-author of more than 200 Journal Articles with the ALICE Collaboration and the full list of publications can be found at <https://inspirehep.net/authors/1627742>

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Multiplicity dependence of (multi-)strange hadron production in proton-proton collisions at $\sqrt{s} = 13$ TeV

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Abstract The production rates and the transverse momentum distribution of strange hadrons at mid-rapidity ($|y| < 0.5$) are measured in proton-proton collisions at $\sqrt{s} = 13$ TeV as a function of the charged particle multiplicity, using the ALICE detector at the LHC. The production rates of K_S^0 , Λ , Ξ , and Ω increase with the multiplicity faster than what is reported for inclusive charged particles. The increase is found to be more pronounced for hadrons with a larger strangeness content. Possible auto-correlations between the charged particles and the strange hadrons are evaluated by measuring the event-activity with charged particle multiplicity estimators covering different pseudorapidity regions. When comparing to lower energy results, the yields of strange hadrons are found to depend only on the mid-rapidity charged particle multiplicity. Several features of the data are reproduced qualitatively by general purpose QCD Monte Carlo models that take into account the effect of densely-packed QCD strings in high multiplicity collisions. However, none of the tested models reproduce the data quantitatively. This work corroborates and extends the ALICE findings on strangeness production in proton-proton collisions at 7 TeV.

1 Introduction

The production rates of strange and multi-strange hadrons in high-energy hadronic interactions are important observables for the study of the properties of Quantum Chromodynamics (QCD) in the non-perturbative regime. In the simplest case, strange quarks (s) in proton-proton (pp) collisions can be produced from the excitation of the sea partons. Indeed, in the past decades a significant effort has been dedicated to the study of the actual strangeness content of the nuclear wave function [1]. In QCD-inspired Monte Carlo generators based on Parton Showers (PS) [2] the hard (perturbative) interactions are typically described at the Leading Order (LO). In this picture the s quark can be produced in the hard par-

tonic scattering via flavour creation and flavour excitation processes as well as in the subsequent shower evolution via gluon splitting. At low transverse momentum, $s\bar{s}$ pairs can be produced via non-perturbative processes, as described for instance in string fragmentation models, where the production of strangeness is suppressed with respect to light quark production due to the larger strange quark mass [3]. However, these models fail to quantitatively describe strangeness production in hadronic collisions [4–6].

An enhanced production of strange hadrons in heavy-ion collisions was suggested as a signature for the creation of a Quark-Gluon Plasma (QGP) [7,8]. The main argument in these early studies was that the mass of the strange quark is of the order of the QCD deconfinement temperature, allowing for thermal production in the deconfined medium. The lifetime of the QGP was then estimated to be comparable to the strangeness relaxation time in the plasma, leading to full equilibration. Strangeness enhancement in heavy-ion collisions was indeed observed at the SPS [9] and higher energies [10,11]. However, strangeness enhancement is no longer considered an unambiguous signature for deconfinement (see e.g. [12]). Strange hadron production in heavy-ion collisions is currently usually described in the framework of statistical-hadronisation (or thermal) models [13,14]. In central heavy-ion collisions, the yields of strange hadrons turn out to be consistent with the expectation from a grand-canonical ensemble, i.e. the production of strange hadrons is compatible with thermal equilibrium, characterised by a common temperature. On the other hand, the strange hadron yields in elementary collisions are suppressed with respect to the predictions of the (grand-canonical) thermal models. The suppression of the relative abundance of strange hadrons with respect to lighter flavours was suggested to be, at least partially, a consequence of the finite volume, which makes the application of a grand-canonical ensemble not valid in hadron-hadron and hadron-nucleus interactions (canonical suppression) [15–17]. However, this approach cannot explain the observed particle abundances if the same volume is assumed for both strange and non-strange hadrons [18] and

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does not describe the system size dependence of the ϕ meson, a hidden-strange hadron [19, 20].

The ALICE Collaboration recently reported an enhancement in the relative production of (multi-) strange hadrons as a function of multiplicity in pp collisions at $\sqrt{s} = 7$ TeV [21] and in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [22, 23]. In the case of p–Pb collisions, the yields of strange hadrons relative to pions reach values close to those observed in Pb–Pb collisions at full equilibrium. These are surprising observations, because thermal strangeness production was considered to be a defining feature of heavy-ion collisions, and because none of the commonly-used pp Monte Carlo models reproduced the existing data [3, 21]. The mechanisms at the origin of this effect need to be understood, and then implemented in the state-of-the-art Monte Carlo generators [3].

In this paper, strangeness production in pp interactions is studied at the highest energy reached at the LHC, $\sqrt{s} = 13$ TeV. We present the measurement of the yields and transverse momentum (p_T) distributions of single-strange (K_S^0 , Λ , $\bar{\Lambda}$) and multi-strange (Ξ^- , Ξ^+ , Ω^- , Ω^+) particles at mid-rapidity, $|y| < 0.5$, with the ALICE detector [24]. The comparison of the present results with the former ones for pp and p–Pb interactions allows the investigation of the energy, multiplicity and system size dependence of strangeness production. Schematically, the multiplicity of a given pp event depends on (i) the number of Multiple Parton Interactions (MPI), (ii) the momentum transfer of those interactions, (iii) fluctuations in the fragmentation process. A systematic study of the biases induced by the choice of the multiplicity estimator along with the specific connections to the underlying MPI are also discussed in this paper.

The paper is organised as follows. In Sect. 2 we discuss the data set and detectors used for the measurement; in Sect. 3 we describe the analysis techniques; in Sect. 4 we cover the evaluation of the systematic uncertainties; in Sect. 5 we present and discuss the results; in Sect. 6 we report our conclusions.

2 Experimental setup and data selection

A detailed description of the ALICE detector and its performance can be found in [24, 25]. In this section, we briefly outline the main detectors used for the measurements presented in this paper. The ALICE apparatus comprises a central barrel used for vertex reconstruction, track reconstruction and charged-hadron identification, complemented by specialised forward detectors. The central barrel covers the pseudorapidity region $|\eta| < 0.9$. The main central-barrel tracking devices used for this analysis are the Inner Tracking System (ITS) and the Time-Projection Chamber (TPC), which are located inside a solenoidal magnet providing a 0.5 T magnetic field. The ITS is composed of six cylindrical layers of high-

resolution silicon tracking detectors. The innermost layers consist of two arrays of hybrid Silicon Pixel Detectors (SPD), located at an average radial distance r of 3.9 and 7.6 cm from the beam axis and covering $|\eta| < 2.0$ and $|\eta| < 1.4$, respectively. The SPD is also used to reconstruct tracklets, short two-point track segments covering the pseudorapidity region $|\eta| < 1.4$. The outer layers of the ITS are composed of silicon strips and drift detectors, with the outermost layer having a radius $r = 43$ cm. The TPC is a large cylindrical drift detector of radial and longitudinal sizes of about $85 < r < 250$ cm and $-250 < z < 250$ cm, respectively. It is segmented in radial “pad rows”, providing up to 159 tracking points. It also provides charged-hadron identification information via specific ionisation energy loss in the gas filling the detector volume. The measurement of strange hadrons is based on “global tracks”, reconstructed using information from the TPC as well as from the ITS, if the latter is available. Further outwards in radial direction from the beam-pipe and located at a radius of approximately 4 m, the Time of Flight (TOF) detector measures the time-of-flight of the particles. It is a large-area array of multigap resistive plate chambers with an intrinsic time resolution of 50 ps. The V0 detectors are two forward scintillator hodoscopes employed for triggering, background suppression and event-class determination. They are placed on either side of the interaction region at $z = -0.9$ m and $z = 3.3$ m, covering the pseudorapidity regions $-3.7 < \eta < -1.7$ and $2.8 < \eta < 5.1$, respectively.

The data considered in the analysis presented in this paper were collected in 2015, at the beginning of Run 2 operations of the LHC at $\sqrt{s} = 13$ TeV. The sample consists of 50 M events collected with a minimum bias trigger requiring a hit in both V0 scintillators in coincidence with the arrival of proton bunches from both directions. The interaction probability per single bunch crossing ranges between 2 and 14%.

The contamination from beam-induced background is removed offline by using the timing information in the V0 detectors and taking into account the correlation between tracklets and clusters in the SPD detector, as discussed in detail in [25]. The contamination from in-bunch pile-up events is removed offline excluding events with multiple vertices reconstructed in the SPD. Part of the data used in this paper were collected in periods in which the LHC collided “trains” of bunches each separated by 50 ns from its neighbours. In these beam conditions most of the ALICE detectors have a readout window wider than a single bunch spacing and are therefore sensitive to events produced in bunch crossings different from those triggering the collision. In particular, the SPD has a readout window of 300 ns. The drift speed in the TPC is about 2.5 cm/ μ s, which implies that events produced less than about 0.5 μ s apart cannot be resolved. Pile-up events produced in different bunch crossings are removed exploiting multiplicity correlations in detectors having different readout windows.

3 Analysis details

The results are presented for primary strange hadrons.¹ The measurements reported here have been obtained for events having at least one charged particle produced with $p_T > 0$ in the pseudorapidity interval $|\eta| < 1$ (INEL > 0), corresponding to about 75% of the total inelastic cross-section. In order to study the multiplicity dependence of strange and multi-strange hadrons, for each multiplicity estimator the sample is divided into event classes based on the total charge deposited in the V0 detectors (V0M amplitude) or on the number of tracklets in two pseudorapidity regions: $|\eta| < 0.8$ and $0.8 < |\eta| < 1.5$. The event classes are summarised in Table 1. Since the measurement of strange hadrons is performed in the region $|y| < 0.5$, the usage of these three estimators allows one to evaluate potential biases on particle production, arising from measuring the multiplicity in a pseudorapidity region partially overlapping with the one of the reconstructed strange hadrons ($N_{\text{tracklets}}^{|\eta| < 0.8}$), or disjoint from it (V0M and $N_{\text{tracklets}}^{0.8 < |\eta| < 1.5}$). In particular, the effect of fluctuations can be expected to be stronger if the multiplicity estimator and the observable of interest are measured in the same pseudorapidity region. The usage of two different non-overlapping estimators allows the study of the effect of a rapidity gap between the multiplicity estimator and the measurement of interest.

The events used for the data analysis are required to have a reconstructed vertex in the fiducial region $|z| < 10$ cm. As mentioned in the previous section, events containing more than one distinct vertex are tagged as pile-up and discarded. For each event class and each multiplicity estimator, the average pseudorapidity density of primary charged-particles $\langle dN_{\text{ch}}/d\eta \rangle$ is measured at mid-rapidity ($|\eta| < 0.5$) using the technique described in [27]. The $\langle dN_{\text{ch}}/d\eta \rangle$ values, corrected for acceptance and efficiency as well as for contamination from secondary particles and combinatorial background, are listed in Table 1. When multiplicity event classes are selected outside the $|\eta| < 0.8$ region, the corresponding charged particle multiplicity at mid-rapidity is characterized by a large variance. In the case of the V0M estimator, the variance ranges between 30 and 70% of the mean $dN_{\text{ch}}/d\eta$ for the highest and lowest multiplicity classes, respectively.

The strange hadrons K_S^0 , Λ , $\bar{\Lambda}$, Ξ^- , $\bar{\Xi}^+$, Ω^- and $\bar{\Omega}^+$ are reconstructed at mid-rapidity ($|y| < 0.5$) with an invariant mass analysis, exploiting their specific weak decay topology. The following decay channels are studied [28]:

¹ A primary particle [26] is defined as a particle with a mean proper decay length $c\tau$ larger than 1 cm, which is either (a) produced directly in the interaction, or (b) from decays of particles with $c\tau$ smaller than 1 cm, excluding particles produced in interactions with material.

Table 1 Event classes selected according to different multiplicity estimators (see text for details). For each estimator, the second column summarises the relative multiplicity w.r.t. the INEL > 0 event class and the third column represents the corresponding fraction of the INEL > 0 cross-section. The average charged particle multiplicity density is reported in the last column for all event classes. For all the multiplicity estimators, the charged particle multiplicity density is quoted in $|\eta| < 0.5$

Event class	$\frac{\langle N_{\text{tracklets}}^{ \eta < 0.8} \rangle_{\text{mult}}}{\langle N_{\text{tracklets}}^{ \eta < 0.8} \rangle_{\text{INEL} > 0}}$	$\frac{\langle N_{\text{tracklets}}^{0.8 < \eta < 1.5} \rangle_{\text{mult}}}{\langle N_{\text{tracklets}}^{0.8 < \eta < 1.5} \rangle_{\text{INEL} > 0}}$	$\frac{\langle dN_{\text{ch}}/d\eta \rangle}{\langle dN_{\text{ch}}/d\eta \rangle_{\text{INEL} > 0}}$	$\sigma/\sigma_{\text{INEL} > 0}$	$N_{\text{tracklets}}^{ \eta < 0.8}$	$N_{\text{tracklets}}^{0.8 < \eta < 1.5}$	$\langle dN_{\text{ch}}/d\eta \rangle$	$N_{\text{tracklets}}^{ \eta < 0.8}$	$N_{\text{tracklets}}^{0.8 < \eta < 1.5}$
INEL > 0	Not applicable	Not applicable	Not applicable	Not applicable	V0M	V0M	V0M	V0M	V0M
I	3.7	4.7	3.8	0.0–0.90%	0.0–0.91%	0.0–0.91%	6.89 $\pm$ 0.11	25.75 $\pm$ 0.40	26.32 $\pm$ 0.40
II	2.9	3.4	2.8	0.90–4.5%	0.91–4.6%	0.91–4.6%	19.83 $\pm$ 0.30	23.42 $\pm$ 0.35	19.51 $\pm$ 0.29
III	2.3	2.7	2.2	4.5–8.9%	4.6–9.1%	4.6–9.1%	16.12 $\pm$ 0.24	18.29 $\pm$ 0.28	15.45 $\pm$ 0.23
IV	2.0	2.2	1.9	8.9–13.5%	9.1–13.7%	9.1–13.7%	13.76 $\pm$ 0.21	14.90 $\pm$ 0.23	13.14 $\pm$ 0.20
V	1.8	1.9	1.7	13.5–18.0%	13.7–18.3%	13.7–18.2%	12.06 $\pm$ 0.18	12.90 $\pm$ 0.19	11.63 $\pm$ 0.17
VI	1.5	1.6	1.4	18.0–27.0%	18.3–27.4%	18.2–27.4%	10.11 $\pm$ 0.15	10.72 $\pm$ 0.16	9.50 $\pm$ 0.14
VII	1.2	1.2	1.1	27.0–36.1%	27.4–36.6%	27.4–36.6%	8.07 $\pm$ 0.12	8.14 $\pm$ 0.12	7.68 $\pm$ 0.11
VIII	0.95	0.85	0.91	36.1–45.3%	36.6–45.9%	36.6–45.9%	6.48 $\pm$ 0.10	5.95 $\pm$ 0.09	6.35 $\pm$ 0.10
IX	0.68	0.55	0.62	45.3–64.5%	45.9–65.0%	45.9–65.0%	4.64 $\pm$ 0.07	3.82 $\pm$ 0.06	4.36 $\pm$ 0.06
X	0.37	0.25	0.38	64.5–100.0%	65.0–100.0%	65.0–100.0%	2.52 $\pm$ 0.04	1.76 $\pm$ 0.03	2.67 $\pm$ 0.04

Table 2 Track, topological and candidate selection criteria applied to K_S^0 , Λ and $\bar{\Lambda}$ candidates. DCA stands for “distance of closest approach”, PV represents the “primary event vertex” and θ is the angle between the momentum vector of the reconstructed V^0 and the displacement vector between the decay and primary vertices. The selection on DCA between V^0 daughter tracks takes into account the corresponding experimental resolution

Topological variable	K_S^0 (Λ and $\bar{\Lambda}$) selection criteria
V^0 transverse decay radius	> 0.50 cm
DCA (Negative/positive track-PV)	> 0.06 cm
Cosine of V^0 pointing angle (θ_{V^0})	> 0.97 (0.995)
DCA between V^0 daughter tracks	< 1.0 standard deviations
Track selection	K_S^0 (Λ and $\bar{\Lambda}$) selection criteria
Daughter track pseudorapidity interval	$ \eta < 0.8$
Daughter track $N_{crossedrows}$	≥ 70
Daughter track $N_{crossed}/N_{findable}$	≥ 0.8
p inner wall TPC (proton only)	> 0.3 GeV/c
TPC dE/dx	$< 5\sigma$
Out-of-bunch pile-up track rejection	Requested for at least one daughter
Candidate selection	K_S^0 (Λ and $\bar{\Lambda}$) selection criteria
Rapidity interval	$ y < 0.5$
Proper lifetime (mL/p)	< 20 cm (30 cm)
Competing V^0 rejection	$5 \text{ MeV}/c^2$ ($10 \text{ MeV}/c^2$)
MC association (MC only)	Identity assumption for V^0 and for daughter tracks

$$\begin{aligned}
 K_S^0 &\rightarrow \pi^+ + \pi^- & \text{B.R.} &= (69.20 \pm 0.05)\% \\
 \Lambda(\bar{\Lambda}) &\rightarrow p(\bar{p}) + \pi^-(\pi^+) & \text{B.R.} &= (63.9 \pm 0.5)\% \\
 \Xi^-(\bar{\Xi}^+) &\rightarrow \Lambda(\bar{\Lambda}) + \pi^-(\pi^+) & \text{B.R.} &= (99.887 \pm 0.035)\% \\
 \Omega^-(\bar{\Omega}^+) &\rightarrow \Lambda(\bar{\Lambda}) + K^-(K^+) & \text{B.R.} &= (67.8 \pm 0.7)\%
 \end{aligned}$$

In the following, we refer to the sum of particles and anti-particles, $\Lambda + \bar{\Lambda}$, $\Xi^- + \bar{\Xi}^+$ and $\Omega^- + \bar{\Omega}^+$, simply as Λ , Ξ and Ω .

The details of the analysis have been discussed in earlier ALICE publications [5,6,18,22]. The tracks retained in the analysis are required to cross at least 70 TPC readout pads out of a maximum of 159. Tracks are also required not to have large gaps in the number of expected tracking points in the radial direction. This is achieved by checking that the number of clusters expected based on the reconstructed trajectory and the measurements in neighbouring TPC pad rows do not differ by more than 20%.

Each decay product arising from V^0 (K_S^0 , Λ , $\bar{\Lambda}$) and cascade (Ξ^- , $\bar{\Xi}^+$, Ω^- , $\bar{\Omega}^+$) candidates is verified to lie within the fiducial tracking region $|\eta| < 0.8$. The specific energy loss (dE/dx) measured in the TPC, used for the particle identification (PID) of the decay products, is also requested to be compatible within 5σ with the one expected for the corresponding particle species’ hypothesis. The dE/dx is evaluated as a truncated mean using the lowest 60% of the values out of a possible total of 159. This leads to a resolution of about 6%. A set of “geometrical” selections is applied in order to identify specific decay topologies (topological selection), improving the signal/background ratio. The topological variables used for V^0 s and cascades are described in detail

in [6]. In addition, in order to reject the residual out-of-bunch pile-up background on the measured yields, it is requested that at least one of the tracks from the decay products of the (multi-)strange hadron under study is matched in either the ITS or the TOF detector. The selections used in this paper are summarised in Table 2 for the V^0 s and in Table 3 for the cascades.

Strange hadron candidates are required to be in the rapidity window $|y| < 0.5$. K_S^0 (Λ) candidates compatible with the alternative V^0 hypothesis are rejected if they lie within $\pm 5 \text{ MeV}/c^2$ ($\pm 10 \text{ MeV}/c^2$) of the nominal Λ (K_S^0) mass. A similar selection is applied to the Ω , where candidates compatible within $\pm 8 \text{ MeV}/c^2$ of the nominal Ξ mass are rejected. The width of the rejected region was determined according to the invariant mass resolution of the corresponding competing signal. Furthermore, candidates whose proper lifetimes are unusually large for their expected species are also rejected to avoid combinatorial background from interactions with the detector material. The signal extraction is performed as a function of p_T . A preliminary fit is performed on the invariant mass distribution using a Gaussian plus a linear function describing the background. This allows for the extraction of the mean (μ) and width (σ_G) of the peak. A “peak” region is defined within $\pm 6(4)\sigma_G$ for V^0 s (cascades) with respect to μ for any measured p_T bin. Adjacent background bands, covering a combined mass interval as wide as the peak region, are defined on both sides of that central region. The signal is then extracted with a bin counting procedure, subtracting counts in the background region from those of the signal region. Alternatively, the signal is extracted by

Table 3 Track, topological and candidate selection criteria applied to charged Ξ^- , Ξ^+ , Ω^- and Ω^+ candidates. DCA stands for “distance of closest approach”, PV represents the “primary event vertex” and θ is the angle between the momentum vector of the reconstructed V^0 or cascade and the displacement vector between the decay and primary vertices. The selection on DCA between V^0 daughter tracks takes into account the corresponding experimental resolution

Topological variable	Ξ (Ω) selection criteria
Cascade transverse decay radius R_{2D}	> 0.6 (0.5) cm
V^0 transverse decay radius	> 1.2 (1.1) cm
DCA (Bachelor-PV)	> 0.04 cm
DCA (V^0 -PV)	> 0.06 cm
DCA (meson V^0 track-PV)	> 0.04 cm
DCA (baryon V^0 track-PV)	> 0.03 cm
DCA between V^0 daughter tracks	< 1.5 standard deviations
DCA (bachelor- V^0)	< 1.3 cm
Cosine of cascade pointing angle(θ_{casc})	> 0.97
Cosine of V^0 pointing angle (θ_{V^0})	> 0.97
V^0 invariant mass window	± 0.008 GeV/ c^2
Track selection	Ξ (Ω) selection criteria
Daughter track pseudorapidity interval	$ \eta < 0.8$
Daughter track $N_{TPCclusters}$	≥ 70
TPC dE/dx	$< 5\sigma$
Out-of-bunch pile-up track rejection	Requested for at least one daughter
Candidate selection	Ξ (Ω) selection criteria
Rapidity interval	$ y < 0.5$
Proper lifetime (mL/p)	$< 3 \times c\tau$
Competing cascade rejection (only Ω)	$ M(\Lambda\pi) - 1321 > 8$ MeV/ c^2
MC association (MC Only)	Identity assumption for cascades and for daughter tracks

fitting the background with a linear function extrapolated under the signal region. This procedure is used to compute the systematic uncertainty due to the signal extraction. Examples of the invariant mass peaks for all particles are shown in Fig. 1.

Only the Λ turns out to be affected by a significant contamination from secondary particles, coming from the decay of charged and neutral Ξ . In order to estimate this contribution we use the measured Ξ^- and Ξ^+ spectra, folded with a p_T -binned 2D matrix describing the decay kinematics and secondary Λ reconstruction efficiencies. The $\Xi \rightarrow \Lambda\pi$ decay matrix is extracted from Monte Carlo (see below for the details on the generator settings). The fraction of secondary Λ particles in the measured spectrum varies between 10 and 20%, depending on p_T and multiplicity. Further details on the uncertainties characterising the feed-down contributions are provided in the next section.

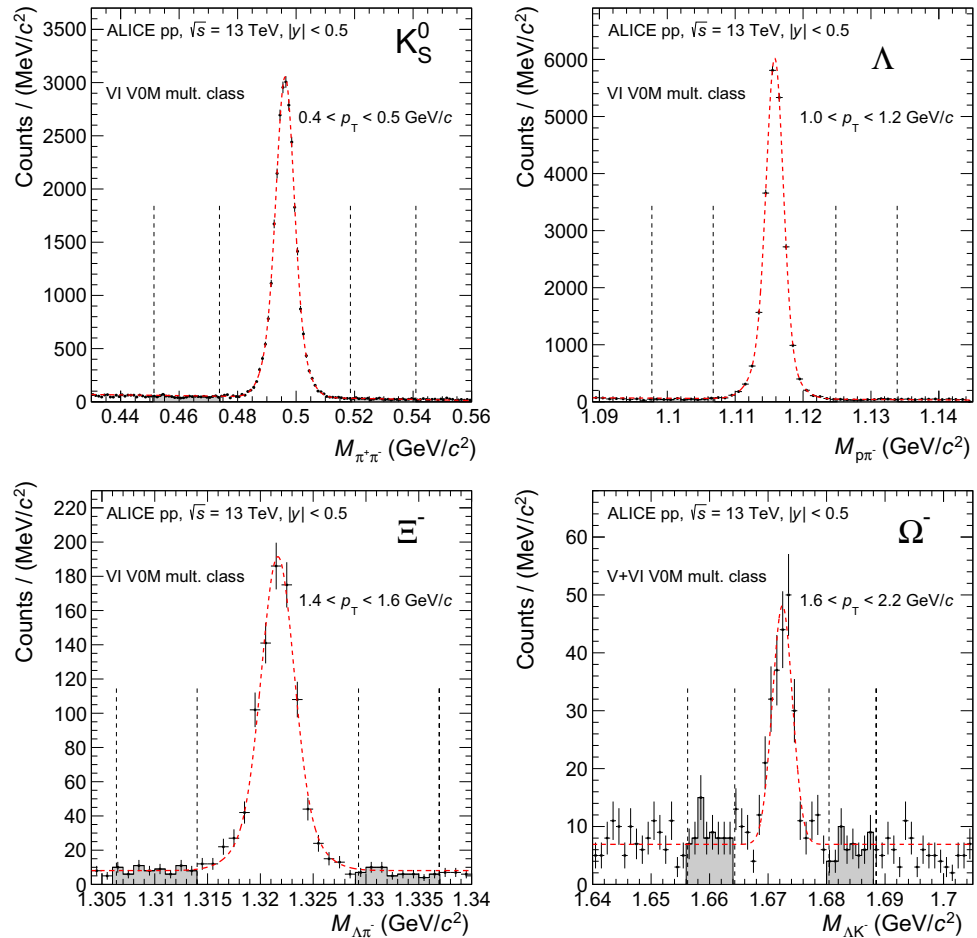
The raw p_T distributions are corrected for acceptance and efficiency using Monte Carlo simulated data. Events are generated using the PYTHIA 6.425, (Tune Perugia 2011) [29,30] event generator, and transported through a GEANT 3 [31] (v2-01-1) model of the detector. With respect to previous GEANT 3 versions, the adopted one contains a more realistic description of (anti)proton interactions. The quality of this description was cross-checked comparing to the

results obtained with the state-of-the-art transport codes FLUKA [32,33] and GEANT 4.9.5 [34]. It was found that a correction factor $< 5\%$ is needed for the efficiency of Λ , Ξ^+ , Ξ^- for $p_T < 1$ GeV/ c , while the effect is negligible at higher p_T . Events generated using PYTHIA 8.210 (tune Monash 2013) [35,36] and EPOS-LHC (CRMC package 1.5.4) [37] and transported in the same way are used for systematic studies, namely to compute the systematic uncertainties arising from the normalisation and from the closure of the correction procedure (details provided in the next sections).

The acceptance-times-efficiency changes with p_T , saturating at a value of about 40%, 30%, 30% and 20% at $p_T \simeq 2, 3, 3$ and 4 GeV/ c for K_S^0 , Λ , Ξ and Ω , respectively. These values include the losses due to the branching ratio. They are found to be independent of the multiplicity class within 2%, limited by the available Monte Carlo simulated data. The dependence of the efficiency on the generated p_T distributions was checked for all particle species. It is found to be relevant only in the case of the Ω , where large p_T bins are used. This effect is removed by reweighting the Monte Carlo p_T distribution with the measured one using an iterative procedure.

In order to compute $\langle p_T \rangle$ and the p_T -integrated production yields, the spectra are fitted with a Tsallis–Lévy [38] distribution to extrapolate in the unmeasured p_T region. The

Fig. 1 Invariant mass distributions for K_S^0 , Λ , Ξ^- , Ω^- in different V0M multiplicity and p_T intervals. The candidates are reconstructed in $|y| < 0.5$. The grey areas delimited by the short-dashed lines are used for signal extraction in the bin counting procedure. The red dashed lines represent the fit to the invariant mass distributions, shown for drawing purpose only



systematic uncertainties on this extrapolation procedure are evaluated using other fit functions, as discussed in Sec. 4.

4 Systematic uncertainties

Several sources of systematic effects on the evaluation of the p_T distributions were investigated. The main contributions for three representative p_T values are summarised in Table 4 for the INEL > 0 data sample.

The stability of the signal extraction method was checked by varying the widths used to define the “signal” and “background” regions, expressed in terms of number of σ_G as defined in Sect. 3. The raw counts were also extracted with a fitting procedure and compared to the standard ones computed by the bin counting technique. An uncertainty ranging between 0.2 and 3.5% depending on p_T is assigned to the signal extraction of the V^0 s and cascades based on these checks.

The stability of the acceptance and efficiency corrections was verified by varying all track, candidate and topological selection criteria within ranges leading to a maximum variation of $\pm 10\%$ in the raw signal yield. The results were compared to those obtained with the default selection criteria

(Sect. 3). Variations not compatible with statistical fluctuations (following the prescription in [39] with a 2σ threshold) are added to the systematic uncertainty.

The resulting uncertainty from topological and track selections (except TPC dE/dx) depends on p_T and amounts at most to 4%, 5%, 4% and 6% for K_S^0 , Λ , Ξ and Ω , respectively.

The TPC dE/dx selection is used to reduce the combinatorial background in the strange baryon invariant mass distribution. The uncertainty was evaluated varying the TPC dE/dx selection requirements between 4 and 7σ and was found to be at most 1% (3%) for Λ (Ξ and Ω). For the K_S^0 the uncertainty due to the TPC dE/dx usage was evaluated by comparing results obtained adopting the default loose PID requirements (5σ) with those obtained without applying any PID selection. The difference is found to be negligible ($< 1\%$).

The systematic uncertainty for the competing decay rejection was investigated by removing entirely this condition for Λ and Ω . It resulted in a deviation on the p_T spectra of at most 4% and 6%, respectively. For the K_S^0 the systematic uncertainty was evaluated by changing the width of the competing rejected mass window between 3 and 5.5 MeV/c^2 and the corresponding deviation was found to be at most 1%.

Table 4 Main sources and values of the relative systematic uncertainties (expressed in %) of the p_T -differential yields. These values are reported for low, intermediate and high p_T . The values for the INEL > 0 data sample are shown in the table. Results as a function of multiplicity

are further affected by an uncertainty originating from the multiplicity dependence of the efficiency (2%) and, in the case of the Λ , of the feed-down contributions (2%)

Hadron p_T (GeV/c)	K_S^0			$\Lambda + \bar{\Lambda}$			$\Xi^- + \bar{\Xi}^+$			$\Omega^- + \bar{\Omega}^+$		
	$\simeq 0.95$	$\simeq 4.8$	$\simeq 9.0$	$\simeq 0.50$	$\simeq 4.5$	$\simeq 7.3$	$\simeq 0.80$	$\simeq 3.2$	$\simeq 5.8$	$\simeq 1.3$	$\simeq 2.8$	$\simeq 4.7$
Signal extraction	0.6	0.6	1.6	1.1	2.4	1.1	0.6	0.4	negl.	2.1	2.1	3.3
Topological and track selection (but TPC dE/dx)	0.7	2.7	2.5	2.9	2.6	2.3	4.8	2.6	2.0	5.6	6.0	5.5
TPC dE/dx selection	0.1	negl.	negl.	0.7	0.3	1.7	0.4	0.3	0.3	1.7	2.2	1.6
Competing decay rejection	0.1	0.2	1.0	negl.	0.9	5.7	Not applied			1.2	3.1	5.6
Proper lifetime	0.1	negl.	negl.	1.1	negl.	negl.	0.5	0.8	0.8	3.1	1.8	1.0
Transport code (for anti-particles)	Not applied			1.8	negl.	negl.	1.1	negl.	negl.	0.6	negl.	negl.
Material budget	1.1	0.5	0.5	8.3	0.8	0.8	5.1	1.2	0.6	3.3	1.5	1.5
Feed-down correction	Not applied			2.2	1.3	2.1	Not applied			Not applied		
Out-of-bunch pile-up track rejection	1.1	1.5	1.5	1.2	2.4	2.4	1.0	2.5	2.5	3.0	3.0	3.0
Residual in-bunch pile-up	1.6	2.5	2.5	2.0	2.9	2.9	2.0	2.0	2.9	2.0	2.0	2.0
Total	2.4	4.1	4.3	9.6	5.5	7.8	7.5	4.4	4.4	8.6	8.6	9.6
Common (N_{ch} -independent)	2.3	4.0	3.8	9.2	5.3	7.0	6.8	4.2	4.2	7.7	8.4	7.5

For the strange baryons, the systematic uncertainty related to the proper lifetime was computed by varying the selection requirements between 2.5 and 5 $c\tau$. The variation range for the K_S^0 was set to 5–15 $c\tau$. The statistically significant deviations were found to be at most 3% for the Λ and negligible (< 1%) for all other particles.

An uncertainty related to the absorption in the detector material was assigned to the anti-baryons, mostly due to the interactions of anti-proton daughters. It was estimated on the basis of the comparison of the different transport codes mentioned in Sect. 3. The uncertainty on the absorption cross section for baryons and K_S^0 was found to be negligible.

Furthermore an additional 2% uncertainty is added to account for possible variations of the tracking efficiency with multiplicity (Sect. 3).

The uncertainty due to approximations in the description of the detector material was estimated with a Monte Carlo simulation where the material budget was varied within its uncertainty [25]. The assigned systematic uncertainty ranges between 8% at low p_T to about 1% at high p_T .

The Λ p_T spectrum is affected by an uncertainty coming from the feed-down correction, due to the uncertainties on the measured Ξ spectrum and on the multiplicity dependence of the feed-down fraction. Furthermore the contribution from neutral Ξ^0 was taken into account by assuming $\Xi^\pm/\Xi^0 = 1$ or using the ratio provided by the Monte Carlo (using the reference PYTHIA 6 sample described in the previous section). The difference between these two estimates was taken

into account in the calculation of the total uncertainty due to the feed-down correction, which ranges from 2% to 4% depending on p_T and multiplicity.

The systematic uncertainty due to the out-of-bunch pile-up rejection was evaluated by changing the matching scheme with the relevant detectors, considering the following configurations: matching of at least one decay track with the ITS (TOF) detector below (above) 2 GeV/c of the reconstructed (multi-) strange hadron; ITS matching of at least one decay track in the full p_T range. Half of the maximum difference between these configurations and the standard selection was taken as the systematic uncertainty, which was found to increase with transverse momentum and to saturate at high p_T , reaching a maximum value of 2.4% (3%) for V^0 s (cascades).

The effect of a possible residual contamination from in-bunch pile-up events was estimated varying the pile-up rejection criteria and dividing the data sample in three groups with an average interaction probability per bunch crossing of 3%, 6%, 13%, respectively. The resulting uncertainty is larger at low multiplicity, and ranges between 1% (3%) for the K_S^0 to 2% (3%) for the baryons in high-(low-)multiplicity events.

Several additional consistency checks have been performed which are described in the following. The analysis has been repeated separately for events with positive and negative z -vertex position, as well as considering candidates reconstructed in positive and negative rapidity windows. The

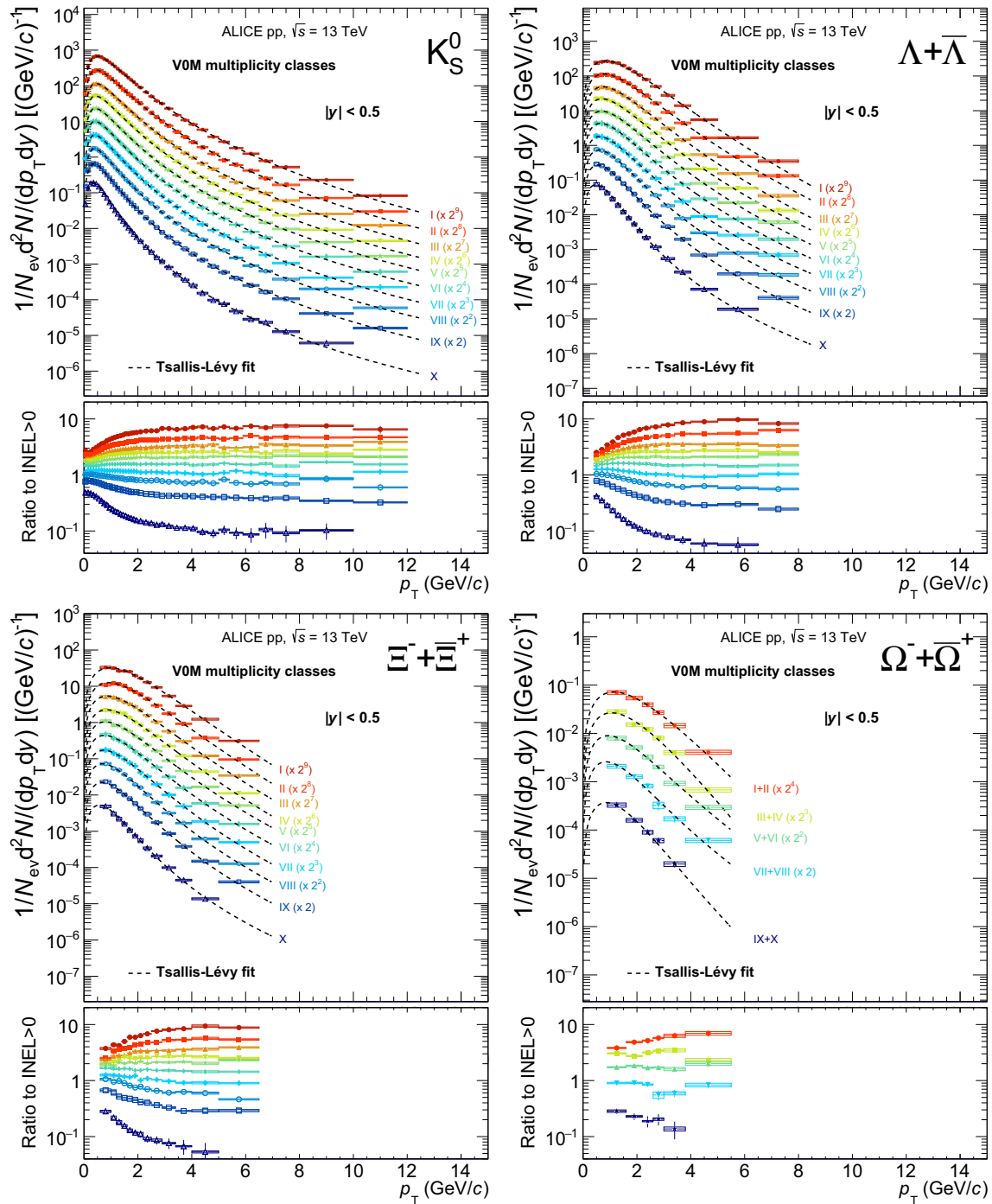


Fig. 2 Transverse momentum distribution of K_S^0 , Λ , Ξ , and Ω , for multiplicity classes selected using the V0 detector. Statistical and total systematic uncertainties are shown by error bars and boxes, respectively. In the bottom panels ratios of multiplicity dependent spectra to $\text{INEL} > 0$ are shown. The systematic uncertainties on the ratios are

obtained by considering only contributions uncorrelated across multiplicity. The spectra are scaled by different factors to improve the visibility. The dashed curves represent Tsallis–Lévy fits to the measured spectra

resulting p_T -spectra were found to be statistically compatible with the standard analysis.

In order to ensure that the estimated systematic uncertainties are not affected by statistical fluctuations, two checks

were performed. First of all, the threshold used to consider a variation statistically significant was varied between one and three standard deviations. Then, the analysis was repeated with wider p_T bins. These checks showed that statistical fluctuations

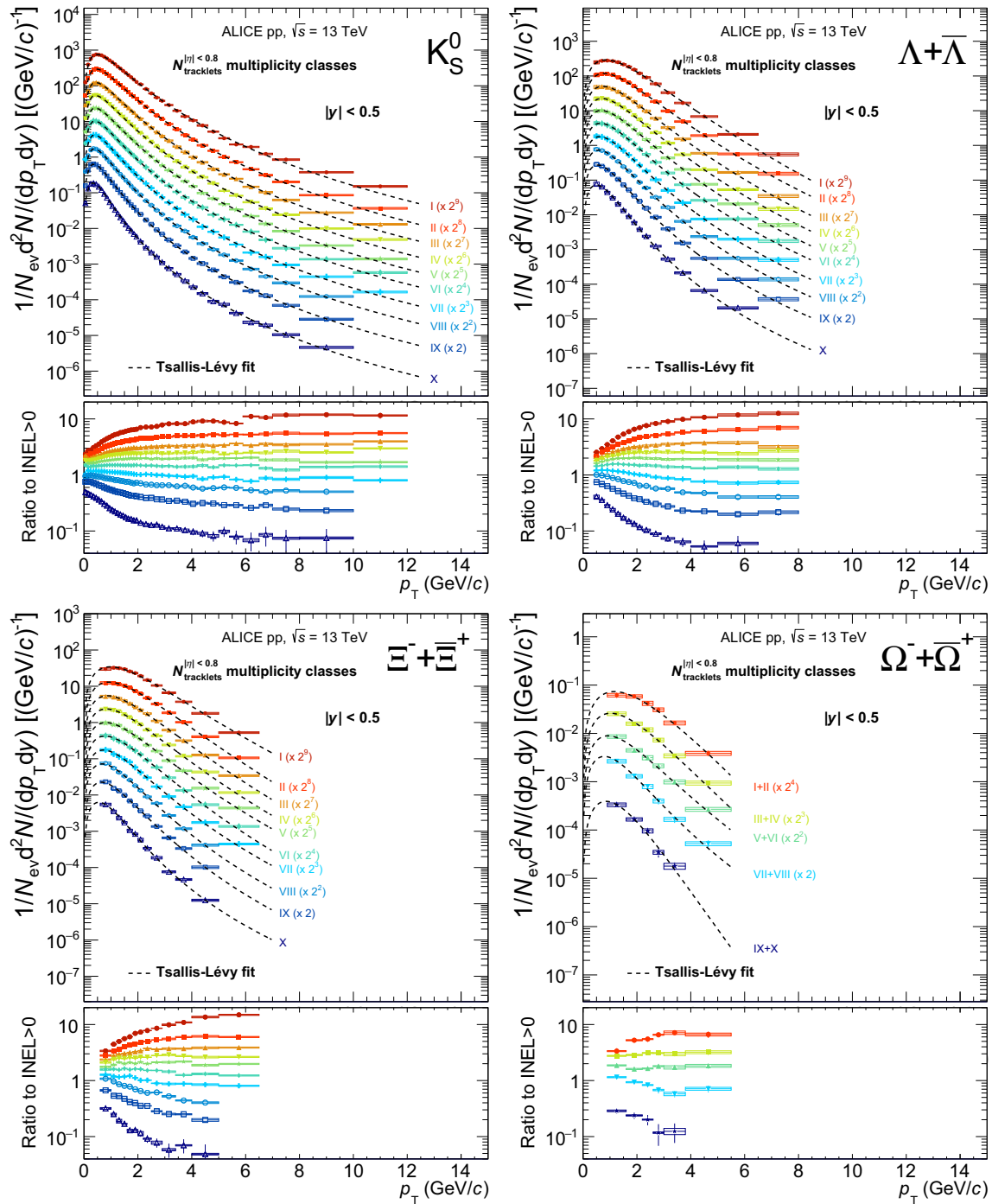


Fig. 3 Transverse momentum distribution of K_S^0 , Λ , Ξ , and Ω , for multiplicity classes selected using the tracklets in $|\eta| < 0.8$. Statistical and total systematic uncertainties are shown by error bars and boxes, respectively. In the bottom panels ratios of multiplicity dependent spectra to INEL > 0 are shown. The systematic uncertainties on the ratios

are obtained by considering only contributions uncorrelated across multiplicity. The spectra are scaled by different factors to improve the visibility. The dashed curves represent Tsallis–Lévy fits to the measured spectra

tuations in the systematic uncertainty analysis are well under control.

The results related to the $\langle p_T \rangle$ and p_T -integrated yields for all particles but the K_S^0 are further affected by an uncer-

tainty coming from the extrapolation to zero p_T . The default extrapolation is performed using a Lévy–Tsallis distribution. As confirmed by a χ^2 analysis, this function describes the

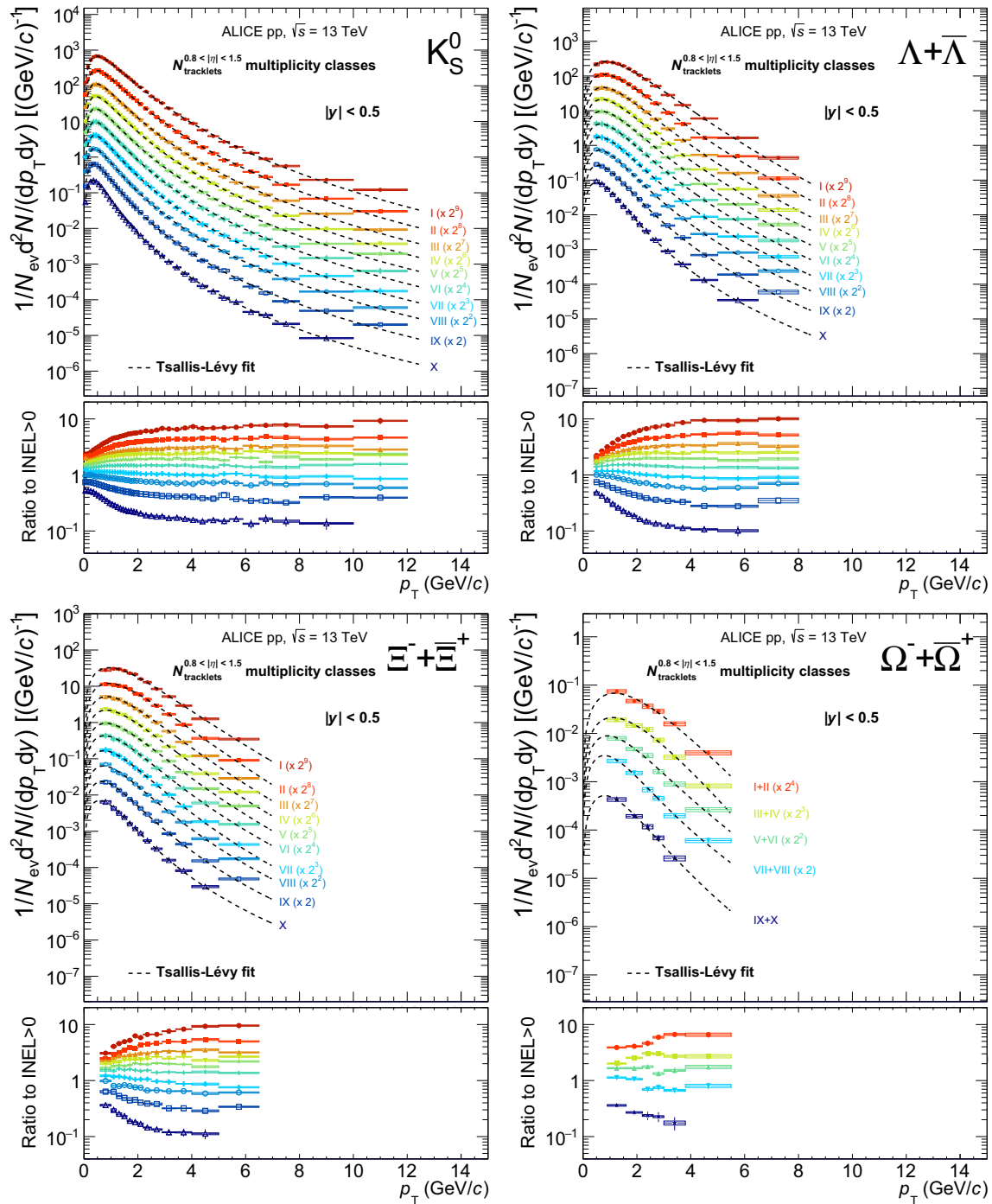


Fig. 4 Transverse momentum distribution of K_S^0 , Λ , Ξ , and Ω , for multiplicity classes selected using the tracklets in $0.8 < |\eta| < 1.5$. Statistical and total systematic uncertainties are shown by error bars and boxes, respectively. In the bottom panels ratios of multiplicity dependent spectra to INEL > 0 are shown. The systematic uncertainties on

the ratios are obtained by considering only contributions uncorrelated across multiplicity. The spectra are scaled by different factors to improve the visibility. The dashed curves represent Tsallis–Lévy fits to the measured spectra

p_T spectra well for all the examined strange hadrons over the measured p_T range.

The uncertainty on the extrapolated fraction was estimated repeating the fit to the spectra with five alternative functions

(Blast-Wave, Boltzmann, Bose–Einstein, m_T -exponential, Fermi–Dirac). Since these alternative functions do not, in general, describe the full p_T -distribution, the fit range was limited to a small number of data points in order to obtain a

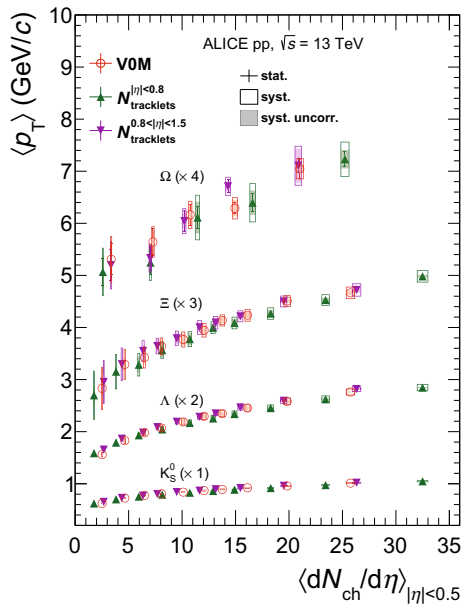


Fig. 5 $\langle p_T \rangle$ of K_S^0 , Λ , Ξ , and Ω in multiplicity event classes selected according to different estimators (see text for details). Statistical and systematic uncertainties are shown by error bars and empty boxes, respectively. Shadowed boxes represent uncertainties uncorrelated across multiplicity

good description of the fitted part of the spectrum. This procedure was repeated separately for the low and for the high p_T extrapolation, with the final uncertainties being dominated by the low p_T contribution. The reliability of the extrapolation uncertainty estimate was checked using a simple linear extrapolation to $p_T = 0$ as an extreme case and in a full Monte Carlo closure test where the EPOS model was used with data and PYTHIA was used to estimate the corrections. The resulting uncertainty on the integrated yields is around 2.5% for the Λ and ranges between 3% (4%) at high multiplicity to 19% (12%) at low multiplicity for the Ξ (Ω). The extrapolation uncertainty on $\langle p_T \rangle$ is $\sim 2\%$ for the Λ and ranges between 2% (3%) at high multiplicity to 12% (7%) at low multiplicity for the Ξ (Ω).

The main focus of this paper is the study of the multiplicity dependence of strangeness production. In this light, the different systematic uncertainties can be categorised in the following way:

1. Fully uncorrelated uncertainties: the change in the data is completely uncorrelated across multiplicity classes. These sources are rare, as most systematic effects have a smooth evolution with multiplicity.
2. Fully correlated uncertainties: lead to a correlated shift of the data in the same direction, independently of the multiplicity class being studied. The common part of this shift has to be considered separately, since it does not affect the shape of the multiplicity-dependent observable.

3. Anti-correlated uncertainties: the effect is opposite in low and high multiplicity events.

In the following we quote separately uncertainties that affect the trends as a function of multiplicity and those that lead to a constant fractional shift across all multiplicity classes. The effect of every systematic variation was evaluated simultaneously in each multiplicity class and in minimum bias events to separate the fully correlated part of the uncertainty.

Sources leading to a global, fully-correlated shift in the spectra were subtracted from the total uncertainty while the remaining contribution could in principle belong to any of the three aforementioned categories. However, since different (independent) sources are combined in the final uncertainties, it is a reasonable assumption to consider them as uncorrelated. For the figures shown in Sect. 5, total uncertainties are shown as boxes, while shadowed boxes represent uncertainties uncorrelated across different multiplicity classes. Statistical uncertainties are depicted as error bars.

5 Results and discussion

Particles and anti-particles turn out to have compatible transverse momentum distributions within uncertainties, consistently with previous results at the LHC. In the following, unless specified otherwise, we present results for their sum. The p_T distributions of strange hadrons, measured in $|y| < 0.5$, are shown in Figs. 2, 3, 4 for the different multiplicity classes, selected using the estimators VOM, tracklets in $|\eta| < 0.8$ and tracklets in $0.8 < |\eta| < 1.5$, as summarised in Table 1. The Lévy–Tsallis fit to the p_T distributions, used for the extrapolation, are also displayed. The bottom panels depict the ratio to the minimum bias (INEL > 0) p_T distribution. The p_T spectra become harder as the multiplicity increases, as also shown in Fig. 5, which shows the average p_T ($\langle p_T \rangle$) as a function of the mid-rapidity charged particle multiplicity. While the VOM and $N_{tracklets}^{0.8 < |\eta| < 1.5}$ results show the same trend, the spectra obtained with the $N_{tracklets}^{|\eta| < 0.8}$ estimator are systematically softer for comparable $dN_{ch}/d\eta$ values (though still compatible within uncertainty in the case of the strange baryons).

The increase of the $\langle p_T \rangle$ as a function of the charged-particle multiplicity (Fig. 5) is compatible, within uncertainties, for all particle species. The hardening of p_T spectra with the charged-particle multiplicity was already reported for pp [40] and p–Pb collisions [22] at lower energies.

It is interesting to notice that the ratio of the measured p_T spectra to the minimum bias one, shown in the bottom panel of Figs. 2, 3, 4, reaches a plateau for $p_T \gtrsim 4$ GeV/c. This applies to all particle species and to all multiplicity

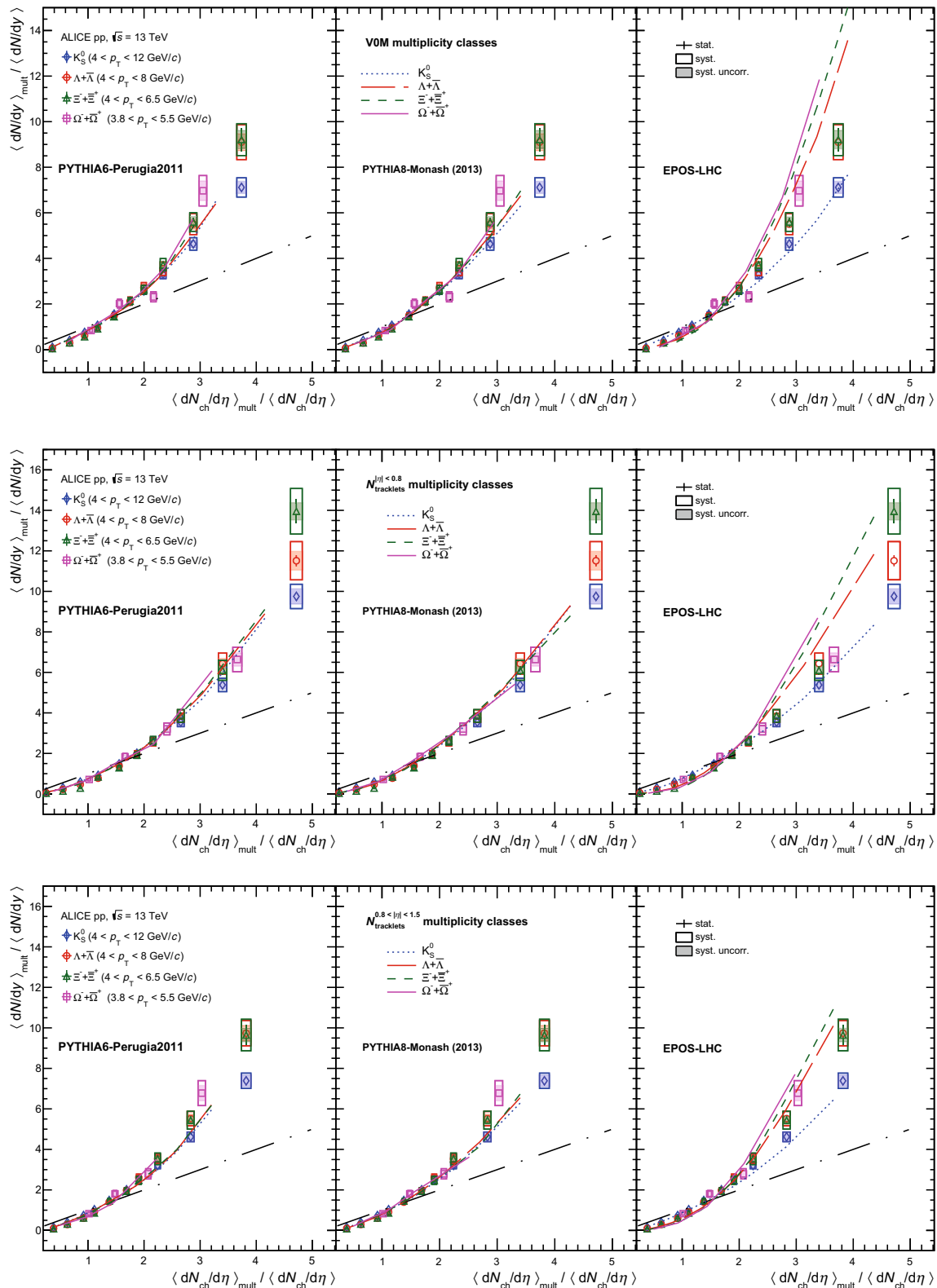


Fig. 6 Self-normalised yields of strange hadrons integrated for $p_T > 4$ GeV/c (3.8 GeV/c for the Ω) versus the self-normalised charged particle multiplicity quoted using different multiplicity estimators (see text for details). Statistical and systematic uncertainties are

shown by error bars and empty boxes, respectively. Shaded boxes represent uncertainties uncorrelated across multiplicity. Strange particle yields are compared to expectations from Monte Carlo models

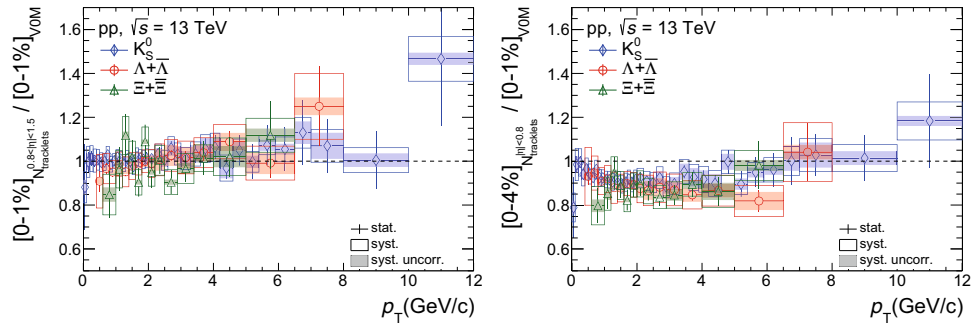


Fig. 7 Ratio of p_T spectra of K_S^0 , Λ and Ξ obtained with two different estimators for multiplicity classes with comparable average $dN_{ch}/d\eta$ value ($\langle dN_{ch}/d\eta \rangle \simeq 26$); the numerator does refer to the $N_{tracklets}^{0.8<|\eta|<1.5}$ ($N_{tracklets}^{|\eta|<0.8}$) estimator while the denominator refers to the VOM estima-

tor. Statistical and systematic uncertainties are shown by error bars and empty boxes, respectively. Shaded boxes represent uncertainties uncorrelated across multiplicity

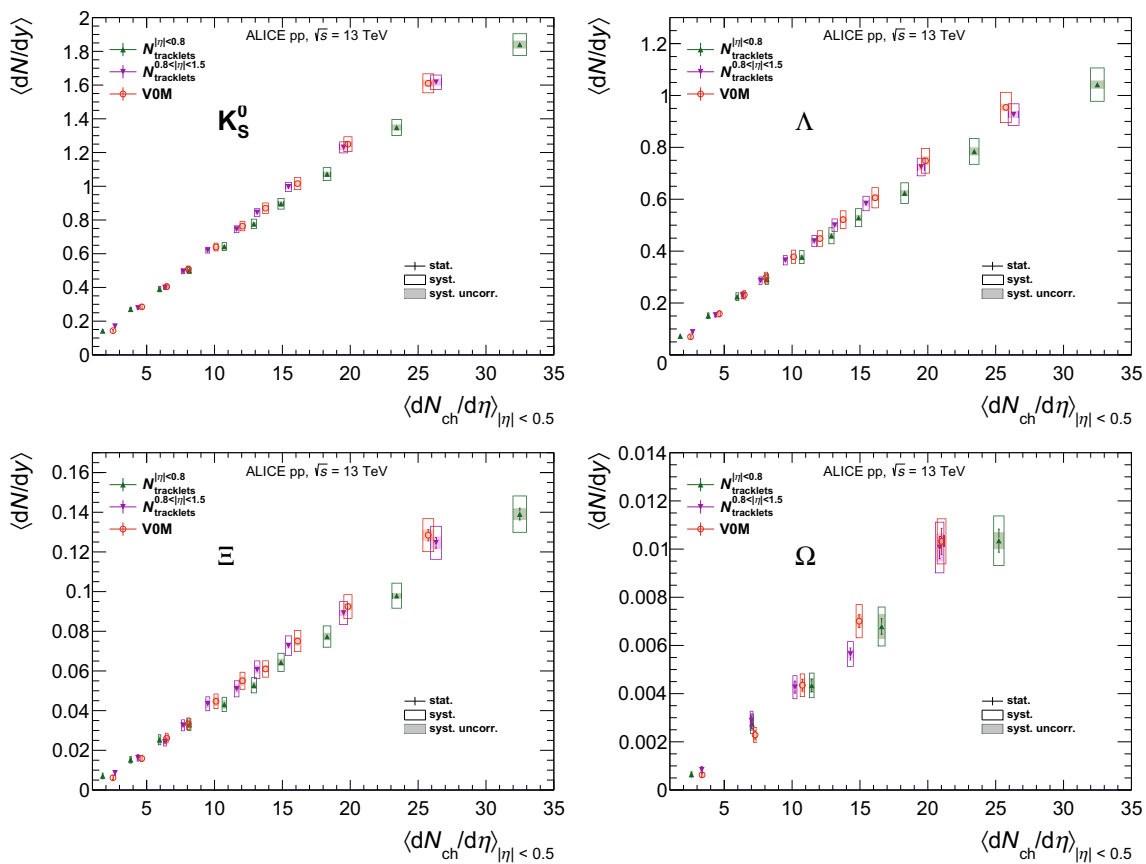


Fig. 8 dN/dy (integrated over the full p_T range) as a function of multiplicity for different strange particle species reported for different multiplicity classes (see text for details). Error bars and boxes represent

statistical and total systematic uncertainties, respectively. The bin-to-bin systematic uncertainties are shown by shadowed boxes

estimators. The trend at high- p_T is highlighted in Fig. 6, which shows the integrated yields for $p_T > 4$ GeV/c as a function of the mid-rapidity multiplicity. Both the yields of strange hadrons and the charged particle multiplicity are self-normalised, i.e. they are divided by their average quoted on the INEL > 0 sample. The high- p_T yields of strange hadrons

increase faster than the charged particle multiplicity. Despite the large uncertainties, the data also hint at the increase being non-linear. The self-normalised yields of strange hadrons reach, at high multiplicity, larger values for the $N_{tracklets}^{|\eta|<0.8}$ multiplicity selection as compared to the other estimators. For all multiplicity selections the self-normalised yields of baryons

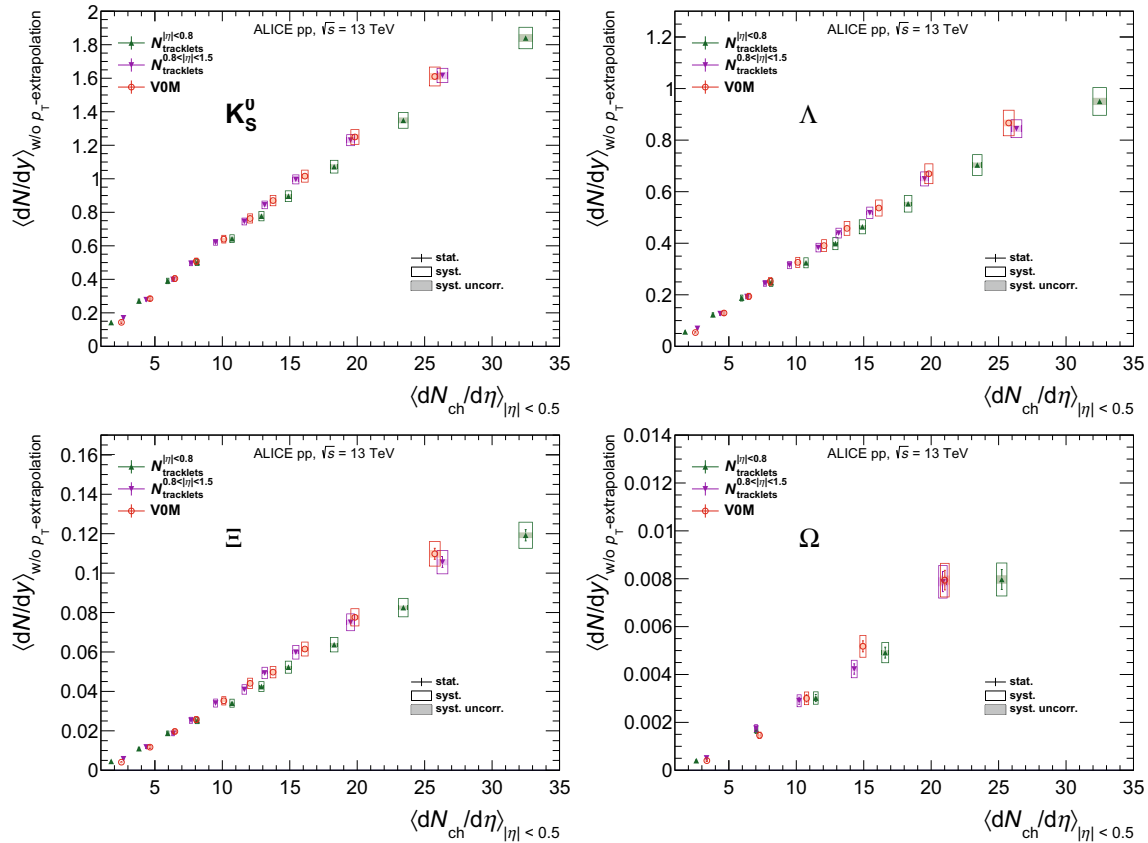


Fig. 9 dN/dy (integrated over the measured p_T ranges 0–12, 0.4–8, 0.6–6.5 and 0.9–5.5 GeV/ c for K_S^0 , Λ , Ξ^- and Ω^- , respectively) as a function of multiplicity for different strange particle species reported for

different multiplicity classes (see text for details). Error bars and boxes represent statistical and total systematic uncertainties, respectively. The bin-to-bin systematic uncertainties are shown by shadowed boxes

are higher than those of K_S^0 mesons. The Monte Carlo models EPOS-LHC, PYTHIA 8 and PYTHIA 6, also shown in Fig. 6, reproduce the overall trend of strange hadrons seen in the data, with EPOS-LHC showing a clear difference between the K_S^0 and the baryons, as observed in the data. Indeed, a non-linear increase can be explained by the combined effect of multiplicity fluctuations and interactions between different MPIs, which produces a collective boost [41]. Both color reconnection effects implemented in PYTHIA and a collective hydrodynamic expansion can account for the non-linear increase pattern.

The left (right) panel of Fig. 7 shows the ratio of the p_T spectra obtained with different estimators for multiplicity classes with comparable average $dN_{\text{ch}}/d\eta$ value ($\langle dN_{\text{ch}}/d\eta \rangle \simeq 26$); the numerator refers to the $N_{\text{tracklets}}^{0.8 < |\eta| < 1.5}$ ($N_{\text{tracklets}}^{|\eta| < 0.8}$) estimator while the denominator refers to the VOM estimator. These selections correspond to the highest multiplicity class studied for the VOM and $N_{\text{tracklets}}^{0.8 < |\eta| < 1.5}$ estimators. The $N_{\text{tracklets}}^{|\eta| < 0.8}$ estimator, on the other hand, was re-adjusted to lead to the same $dN_{\text{ch}}/d\eta$ value. These classes correspond to approximately 1%, 1% and 4% of the INEL > 0 cross sec-

tion for VOM, $N_{\text{tracklets}}^{|\eta| < 0.8}$ and $N_{\text{tracklets}}^{0.8 < |\eta| < 1.5}$, respectively. It is seen that the spectra are identical within uncertainties for the VOM and $N_{\text{tracklets}}^{0.8 < |\eta| < 1.5}$ estimators. The comparison between the VOM and $N_{\text{tracklets}}^{|\eta| < 0.8}$ estimators shows that the bias introduced by the latter estimator does not depend on the particle species and is more pronounced at low and intermediate p_T , although the uncertainties are large.

The p_T -integrated yields of strange hadrons are shown in Fig. 8 for the three estimators considered in this paper. For reference, Fig. 9 shows the results integrated in the measured p_T range with no extrapolation. These are characterised by a smaller uncertainty and can be useful for Monte Carlo tuning. In the rest of this section, we focus on the discussion of the fully extrapolated yields. The results obtained with the VOM and $N_{\text{tracklets}}^{0.8 < |\eta| < 1.5}$ estimators follow a similar linear trend, while the results obtained with $N_{\text{tracklets}}^{|\eta| < 0.8}$ tend to saturate, showing a lower dN/dy for a similar high-multiplicity class. In order to gain insight on this difference, a generator-level PYTHIA (tune Perugia 2011) simulation study was performed. The abundance of strange hadrons in $|\eta| < 0.5$ was studied as a function of the estimated mid-rapidity charged

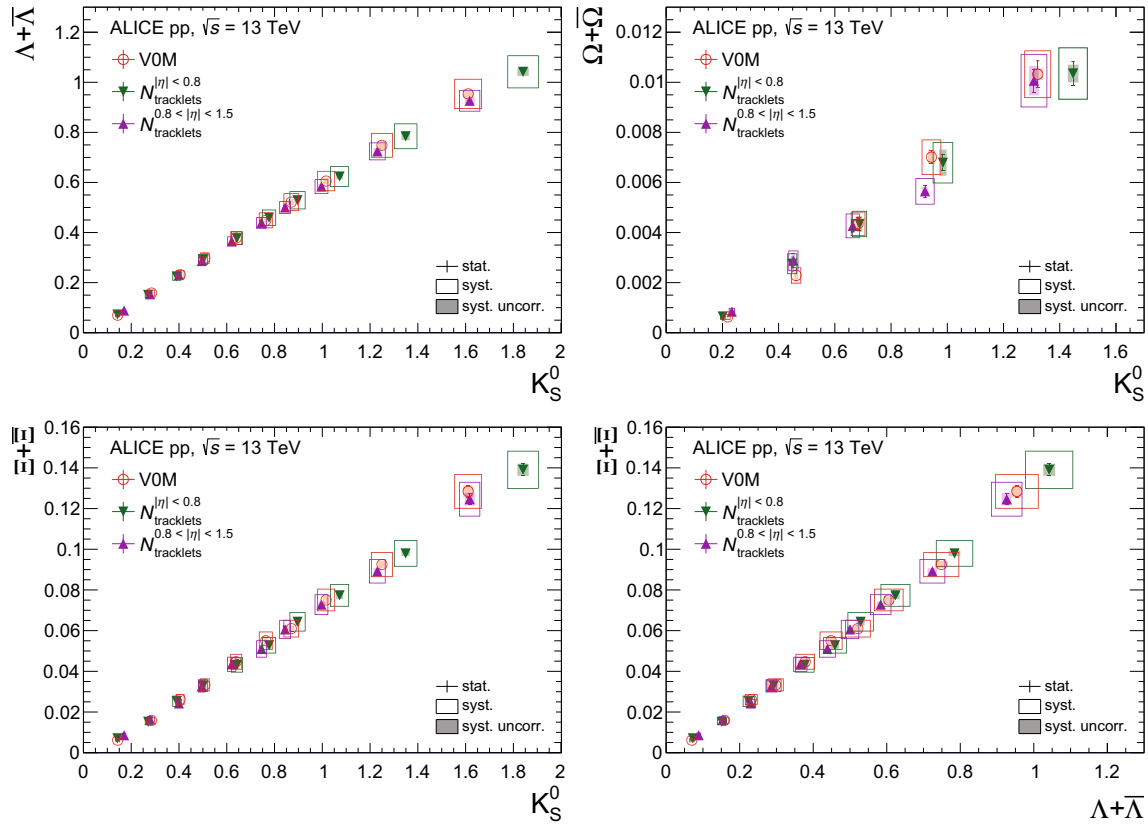


Fig. 10 Correlations between integrated yield of different strange hadrons in multiplicity classes selected according to different estimators (see text for details). Statistical and systematic uncertainties are

shown by error bars and empty boxes, respectively. Shaded boxes represent uncertainties uncorrelated across multiplicity

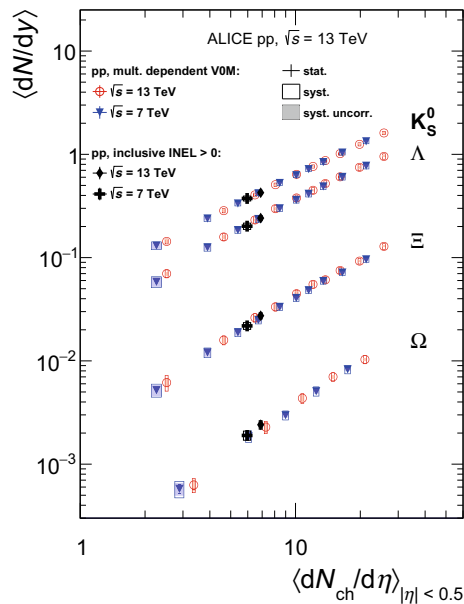


Fig. 11 Integrated yields of K_S^0 , Λ , Ξ , and Ω as a function of $dN_{ch}/d\eta$ in VOM multiplicity event classes at $\sqrt{s} = 7$ and 13 TeV. Statistical and systematic uncertainties are shown by error bars and empty boxes, respectively. Shaded boxes represent uncertainties uncorrelated across multiplicity. The corresponding results obtained for INEL > 0 event class are also shown

particle multiplicity for several event classes, selected using charged primary particles measured in the η ranges corresponding to the experimental estimators presented in this paper: this generator level study confirms the trend observed in the data, which can be understood in terms of a selection bias sensitive to fluctuations in particles yields. Indeed, an estimator based on charged tracks enhances charged primary particles over neutral particles or secondaries. If the multiplicity estimator and the observable under study are measured in the same η region one expects primary charged-particles to be enhanced with respect to strange hadrons, which explains the saturation of yields observed in Fig. 8 for the $N_{tracklets}^{|\eta| < 0.8}$ estimator. As a further check of this interpretation, Fig. 10 shows the correlations between the yields of different strange hadrons for the multiplicity estimators studied in this paper. The trend is linear, and similar for all estimators, confirming that the selection bias on primary charged-particles is stronger than that on strange hadrons.

The energy dependence of the strangeness yields and $\langle p_T \rangle$ versus the charged particle multiplicity at mid-rapidity is studied in Figs. 11 and 12, where our results are compared to the previous pp measurements at $\sqrt{s} = 7$ TeV [21]. The

minimum bias results in the $\text{INEL} > 0$ event class at $\sqrt{s} = 7$ and 13 TeV [5] are also shown.

As can be seen in Fig. 11, the yields of strange hadrons increase with the charged particle multiplicity following a power law behaviour, and the trend is the same at $\sqrt{s} = 7$ and 13 TeV. The $\text{INEL} > 0$ results also follow the same trend at all the tested centre-of-mass energies. This result indicates that the abundance of strange hadrons depends on the local charged particle density and turns out to be invariant with the collision energy, i.e. an energy scaling property applies for the multiplicity-dependent yields of strange hadrons. It should also be noted that the yields of particles with larger strange quark content increase faster as a function of multiplicity as already reported in [21]. Figure 13 shows the Λ/K_S^0 (no $\bar{\Lambda}$ contribution considered here), Ξ/K_S^0 and Ω/K_S^0 ratios compared to calculations from grand-canonical thermal models [13, 14], which were found to satisfactorily describe central Pb-Pb data at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. In the context of a canonical thermal model for a gas of hadrons, an increase of the relative strangeness abundance depending on the strange quark content can be interpreted as a consequence of a change in the system volume, called canonical suppression. Indeed, it was recently shown in [18]

that the existing ALICE data can be described within this framework, introducing an additional parameter to quantify the rapidity window over which strangeness is effectively correlated. It was also suggested that strangeness follows a universal scaling behaviour in all colliding systems, when the transverse energy density [42] or the multiplicity per transverse area [43] are used as a scaling variable. While there are some caveats in the observation reported in these papers (due to the uncertainties on the transverse size of the system) this is a very intriguing observation, that could help to clarify the origin of strangeness enhancement.

The $\langle p_T \rangle$ is seen in Fig. 12 to be harder at 13 TeV than at 7 TeV for event classes with a similar $dN_{\text{ch}}/d\eta$. All the tested Monte Carlo models describe in a qualitative way the observed smooth rise of the $\langle p_T \rangle$ with $dN_{\text{ch}}/d\eta$; from a quantitative point of view, EPOS-LHC provides a slightly better description of the $\langle p_T \rangle$ multiplicity evolution, especially for what concerns the strange baryons.

In Fig. 14 the results on the strange hadron yields as a function of the charged particle multiplicity at mid-rapidity are compared with some commonly-used general purpose QCD inspired models, focusing on the multiplicity classes defined by the VOM estimator. The yields of all measured strange

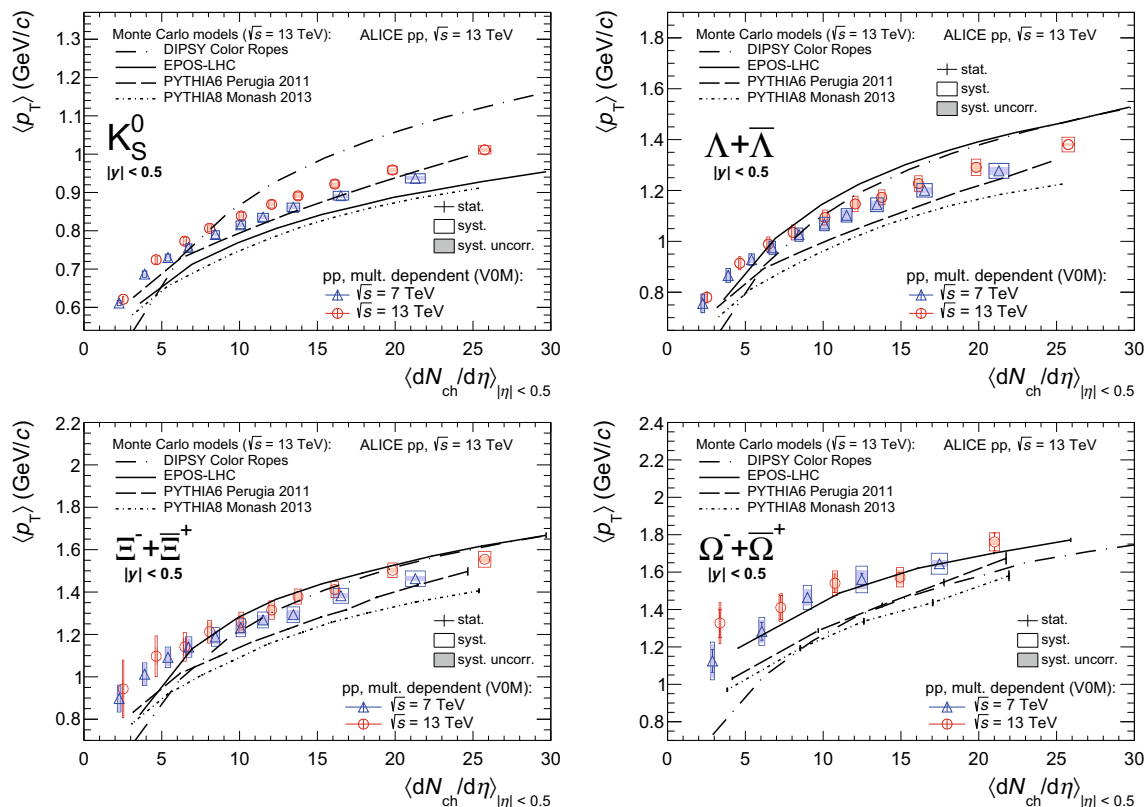


Fig. 12 $\langle p_T \rangle$ of K_S^0 , Λ , Ξ , and Ω as a function of $dN_{\text{ch}}/d\eta$ in VOM multiplicity event classes at $\sqrt{s} = 7$ and 13 TeV. Statistical and systematic uncertainties are shown by error bars and empty boxes, respectively.

Shaded boxes represent uncertainties uncorrelated across multiplicity. The results are compared to predictions from several Monte Carlo models

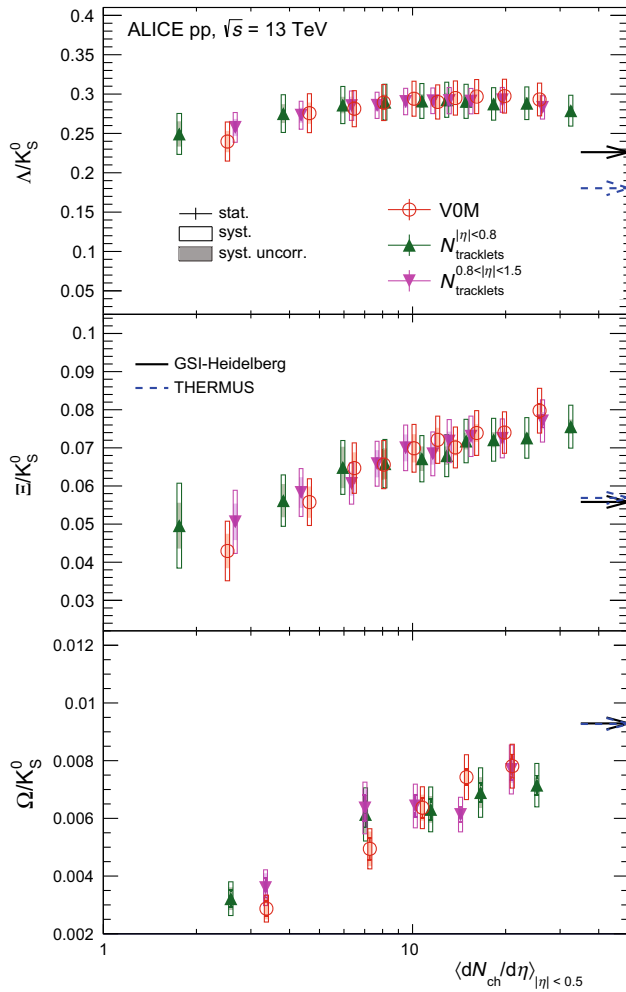


Fig. 13 Ratios of integrated yields of Λ (no $\bar{\Lambda}$ contribution considered here), Ξ , and Ω to K_S^0 as a function of $dN_{\text{ch}}/d\eta$ for different multiplicity estimators (see text for details) in pp collisions at $\sqrt{s} = 13$ TeV. Statistical and systematic uncertainties are shown by error bars and empty boxes, respectively. Shaded boxes represent uncertainties uncorrelated across multiplicity. The corresponding calculations from grand-canonical thermal models, which refer to most central Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV, are shown

hadrons show an almost linear increase with multiplicity: a linear fit to the baryon trends would intercept the $dN_{\text{ch}}/d\eta$ axis at positive values. Alternatively one may regard such qualitative observation as the trend of the points to have an enhanced (more than linear) increase for the hadrons with larger strangeness content.

The PYTHIA 8.210 event generator (tune Monash 2013) gives a reasonable description of the K_S^0 results, however it shows a less pronounced increase of the strange baryons yields versus the charged particle multiplicity than what is observed in the data. The description worsens with increasing strangeness content. Recent attempts to improve the PYTHIA color reconnection scheme to account for the observed strangeness enhancement are discussed in [3], how-

ever they cannot explain the reported experimental observations.

The DIPSY Monte Carlo [44,45] is based on a BFKL inspired initial state dipole evolution model [46] interfaced to the Ariadne model [47] for final state dipole evolution and to the PYTHIA hadronisation scheme, where the latter has been modified to take into account the high density of strings that occurs in events with several MPI. More specifically, depending on their transverse density, the strings are allowed to form “colour-ropes” leading to increased string tension, which leads to an increase of strangeness production stronger than in default PYTHIA [48]. However, despite a qualitative improvement in the description of the data, the discrepancy for the Ω distributions remains large. Moreover, as discussed in [21], DIPSY also predicts an increase of protons normalised to pions as a function of multiplicity, which is not observed in our $\sqrt{s} = 7$ TeV data.

The Monte Carlo generators based on EPOS [49] rely on the Parton Based Gribov Regge Theory (PBGRT) [50]. A common feature of all the EPOS versions is a collective evolution of matter in the secondary scattering stage in all reactions, from pp to AA, with a core-corona separation mechanism [51] which defines the initial conditions of the secondary interactions. In EPOS, the initial parton scatterings create “flux tubes” that either escape the medium and hadronise as jets or contribute to the core, described in terms of hydrodynamics. The core is then hadronised in terms of a grand-canonical statistical model. The relative amount of multi-hadron production arising from the core grows with the number of MPIs. The EPOS-LHC model adopted in this analysis (distributed with the CRMC package 1.5.4) shows a pronounced increase of the strangeness yields as a function of the charged particle multiplicity at mid-rapidity. Despite the differences in the underlying physics, the comparison to data leads to a similar conclusion for EPOS-LHC and DIPSY.

The comparison of the Monte Carlo model predictions to the data indicates that the origin of the strangeness enhancement in hadronic collisions and its relation to the QCD deconfinement phase transition are still open problems for models. While the current versions of these models do not reproduce the data quantitatively, it is possible that the agreement can be improved with further tuning or with new implementations. A quantitative description of strangeness production and enhancement in a microscopic Monte Carlo model would represent a major step to understand flavour production in hadronic collisions at high energy.

6 Summary

We studied the production of primary strange and multi-strange hadrons at mid-rapidity in pp collisions at $\sqrt{s} = 13$ TeV, focusing on the multiplicity dependence. The

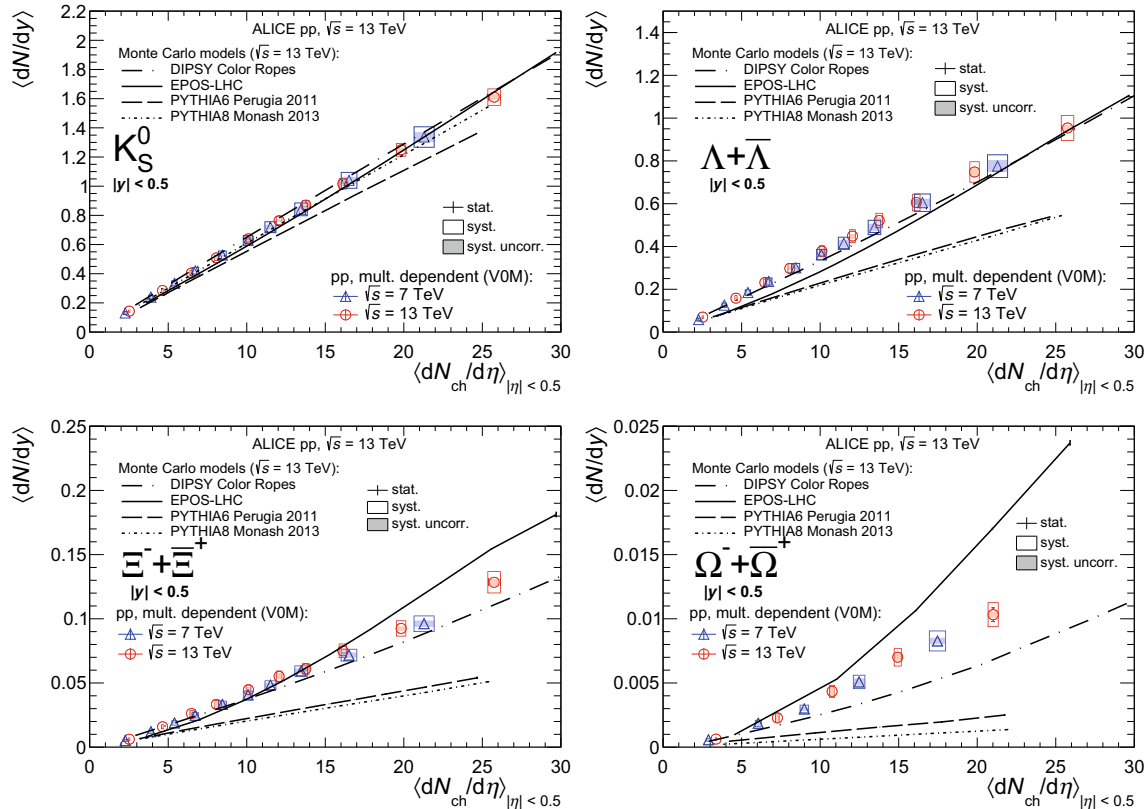


Fig. 14 Integrated yields of K_S^0 , Λ , Ξ , and Ω as a function of $dN_{ch}/d\eta$ in VOM multiplicity event classes at $\sqrt{s} = 7$ and 13 TeV. Statistical and systematic uncertainties are shown by error bars and empty

boxes, respectively. Shaded boxes represent uncertainties uncorrelated across multiplicity. The results are compared to predictions from several Monte Carlo models

main feature of this analysis is the usage of multiplicity estimators defined in different pseudorapidity regions, providing a different sensitivity to the fragmentation and MPI components of particle production and allowing for a detailed study of the selection biases due to fluctuations.

Hardening of the p_T spectra with the increase of the multiplicity is observed, as already reported for pp [40] and p-Pb collisions [22] at lower energies.

The p_T -integrated yields of strange hadrons increase as a function of multiplicity faster than the ones of unidentified charged-particles. This behaviour is even more pronounced for particles with higher strangeness content, confirming the earlier observations in pp collisions at $\sqrt{s} = 7$ TeV [21]. This leads to a multiplicity-dependent increase of the ratio of the strange baryons Ξ and Ω over K_S^0 , while the Λ over K_S^0 ratio turns out to be constant within uncertainties. In the context of a canonical thermal model, an increase of the relative strangeness abundance depending on the strange quark content can be understood as a consequence of an increase in the system volume leading to a progressive removal of canonical suppression.

Comparing the 13 TeV results to the 7 TeV ones, the data exhibit an interesting scaling property: for multiplicity esti-

mators selected in the forward region, the strange hadron yields turn out to be independent of the centre of mass energy, the $\sqrt{s}$ increase resulting just in harder p_T spectra.

The use of high-multiplicity triggered data collected during the full Run 2 period will allow us to test these scaling ansatzes extending the measurement to higher multiplicities, comparable to peripheral Pb-Pb collisions at the LHC ($dN_{ch}/d\eta \sim 40$ –50).

The color reconnection effects implemented in PYTHIA 8 and DIPSY as well as the collective hydrodynamic expansion implemented in EPOS-LHC could not account quantitatively for all the reported results. Some of the qualitative features of the data are described by these models, including the indication of a non-linear increase of the yields of strange hadrons at high p_T , similar to what was previously reported in [52] for J/ψ and D meson production. Although none of the tested models can reproduce the data from a quantitative point of view, EPOS-LHC and DIPSY do provide a better qualitative description of the strangeness enhancement.

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Data Availability Statement This manuscript has associated data in a data repository. [Authors' comment: The numerical values of the data points will be uploaded to HEPData.]

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ALICE Collaboration

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To whom it may concern,

On behalf of the Local Organizing Committee of the 19th International Conference on Hadron Spectroscopy and Structure (HADRON 2021), held virtually from July 26 to July 31, 2021 and hosted by Universidad Nacional Autónoma de México (UNAM, Mexico), I certify that Ms. Meenakshi Sharma (University of Jammu, India) attended the conference and gave the presentation entitled “Recent results on strange and multi-strange particle production with ALICE at the LHC” within the “Hadrons in hot and nuclear environment including hypernuclei” session of the conference.

Best regards,

Pablo Roig

Co-chair of the 19th International Conference on Hadron Spectroscopy and Structure

Permanent researcher

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Certificate

Meenakshi Sharma

has successfully participated and presented a poster titled

Energy dependence of strangeness production in small
collision systems with ALICE

XXIV DAE-BRNS Symposium on High Energy Physics

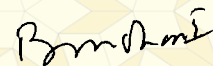
held **ONLINE** by the

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during

December 14-18, 2020



Prof. Bedangadas Mohanty
(Chair)



Prof. Sanjay Swain
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