

Production of strangeness in partonic interactions at the LHC

Production of strangeness in partonic interactions at the LHC

by Oliver Matonoha



LUND
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Thesis for the degree of Doctorate

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*Dedicated to
Humpty – Dumpty
bla bla blat*

Contents

List of publications	iv
Acknowledgements	v
Populärvetenskaplig sammanfattning på svenska	vi
Production of strangeness in partonic interactions at the LHC	1
I Fundamental theory	3
1 Introduction to quantum chromodynamics	5
1.1 Standard Model of elementary particles	5
1.2 Coordinate systems and kinematic observables	8
1.3 Quantum electrodynamics, electrons, and photons	8
1.4 Quantum chromodynamics, quarks, and gluons	10
1.5 From partons to hadrons	12
1.5.1 Initial and Final State Radiation	12
1.5.2 Factorisation theorem	13
1.5.3 Parton distribution functions	14
1.5.4 Parton fragmentation and the Lund string	15
1.6 Multiple partonic interactions	17
1.6.1 Colour reconnection	19
1.7 Underlying event	20
1.8 From hadrons to partons: deconfined QCD matter	21
1.8.1 Bag model of hadrons	21
1.8.2 Lattice QCD	22
1.8.3 QCD phase diagram, chiral symmetry restoration	23
2 QCD phenomena in high energy hadronic collisions	25
2.1 Collisions of heavy nuclei	25
2.1.1 Collision geometry, centrality, multiplicity	25
2.1.2 MC Glauber model	26
2.2 Quark-gluon plasma	28
2.2.1 Quarkonium dissociation and sequential suppression	29
2.2.2 Strangeness enhancement	31

2.2.3	Collective flow	32
2.2.4	Jet quenching	35
2.2.5	Cold nuclear matter effects	36
2.3	QGP phenomena in small systems	36
2.3.1	Role of multiplicity	41
2.4	Phenomenological models	42
2.4.1	Pythia	42
2.4.2	String interactions and Ropes	44
2.4.3	EPOS LHC	45
II	Experimental Setup and Methodology	49
3	ALICE's adventures at the LHC	51
4	Events, Tracks, and Vertices	53
III	Author's measurements	55
5	Reconstruction of neutral strange particles with ALICE	57
5.1	Analysed datasets	57
5.2	Identification of V^0 s using ALICE	58
5.3	Signal extraction	58
5.3.1	Validation using simulations	60
5.4	Normalisation	61
5.5	Corrections to the reconstructed production	62
5.5.1	Secondary contribution correction	62
5.5.2	Reconstruction efficiency	64
5.6	Transverse momentum spectra	65
5.6.1	Comparisons with previously published results	65
5.7	Systematic uncertainties	66
5.7.1	Variation of selection criteria	68
5.7.2	Feeddown correction	69
5.8	Transverse spherocity	70
5.8.1	Motivation for studying event topology	70
5.8.2	S_O and $S_O^{(p_T=1.0)}$ as experimental observables	72
5.8.3	Relationship between $S_O^{(p_T=1.0)}$ and S_O	74
5.8.4	Track and event selection	75
5.8.5	Multiplicity selection and its interplay with $S_O^{(p_T=1.0)}$	76
5.8.6	Comparison of V^0 production with MC generators	80
5.9	Systematic uncertainties	81
5.9.1	Experimental bias	82
5.9.2	Correlation of uncertainties with $S_O^{(p_T=1.0)}$	83

5.9.3	Summary	83
5.10	Transverse momentum spectra vs. $S_O^{(p_T=1.0)}$	85
5.10.1	Ratios of neutral kaons to charged kaons	87
5.11	Mean transverse momenta and integrated yields	87
5.12	Ratios to pions	89
5.13	Baryon-to-meson ratio	91
5.14	Ratio of integrated yields vs. $S_O^{(p_T=1.0)}$	92
6	Underlying Event Activity	95
7	Discussion of results, conclusions, and outlooks	97
7.1	Summary of the research goals	97
7.2	Highlights of the $S_O^{(p_T=1.0)}$ measurement	98
7.3	Highlights of the R_T measurements	99
7.4	Conclusions and Outlooks	101
IV	Appendices	103
A	List of Acronyms	105
B	Mathematical Derivations	107
C	Complementary Material	109
D	Scientific Publications	III
	Author contributions	III
	Paper I: Title paper 1	III
	Paper II: Title paper 2	III
	Paper I: Title paper 1	II3
	Paper II: Title paper 2	II7
	References	121

List of publications

This thesis is based on the following publications, referred to by their Roman numerals:

I Title paper 1

S. Doctor, B. Someone

The Journal of Physical Chemistry A, 2020, 124(19), pp. 3943-3946

II Title paper 2

S. Doctor, B. Someone, C Another

Physical Chemistry Chemical Physics, 2020, 22(24), pp. 13659-13665

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Acknowledgements

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Production of strangeness in partonic interactions at the LHC

Part I

Fundamental theory

Chapter I

Introduction to quantum chromodynamics

This chapter serves as an introduction to particle physics, QCD, and phenomenology of high energy QCD interactions, with the focus on multiple partonic interactions and string formations. Furthermore, it introduces the physics of QCD matter and the deconfinement of hadrons.

I.1 Standard Model of elementary particles

The Standard Model (SM) of particle physics is a set of theories describes *elementary* constituents of matter and their interactions via fundamental forces of the Universe. It has been formulated in the 1970s, combining frameworks of quantum field theory (QFT), gauge symmetries, and spontaneous symmetry breaking.

Matter particles in the SM are classified into two main categories: quarks and leptons. Quarks come in six flavours (*up, down, charm, strange, top, and bottom*) and form hadrons, i.e. *baryons* (qqq) and *mesons* ($q\bar{q}$). The lepton sector also comprises six flavours (*electron, muon, tau, and their corresponding neutrinos*). Quarks and leptons are fermions with an intrinsic spin 1/2. Furthermore, matter particles in the SM also come with associated antiparticles, which have opposite quantum numbers but the same mass.

The interactions between matter particles in the SM are mediated by an exchange of gauge bosons. There are three fundamental forces in the SM, described by four types of vector bosons (ordered by their typical strength):

1. *Strong force*, mediated by the gluon.
2. *Electromagnetic force*, mediated by the photon.
3. *Weak force*, mediated by the massive bosons W^\pm and Z^0 .

Moreover, the interactions are associated with local gauge symmetries, which determine their mathematical structure. The symmetry group for the SM is $SU(3) \times SU(2) \times U(1)$, corresponding to the strong and the electroweak sector. [1]

In addition, the SM also includes the scalar Higgs boson, which is responsible for giving mass to other elementary particles. This is achieved via the Higgs mechanism [2, 3], which involves the spontaneous breaking of the electroweak symmetry in the early universe. The Higgs boson was discovered experimentally at the Large Hadron Collider (LHC) in 2012 by the ATLAS [4] and CMS [5] collaborations, confirming a key prediction of the SM.

Nevertheless, the SM has several limitations, including its inability to account for dark matter, explain why the particle masses span over several orders of magnitude, and the CP violation problem related to the observed matter-antimatter asymmetry in the Universe. These are actively investigated in Beyond Standard Model (BSM) theories.

Quarks

$2.2^{+0.5}_{-0.4} \text{ MeV}$	$1.28 \pm 0.03 \text{ GeV}/c^2$	$172.76 \text{ GeV}/c^2$
u <i>up</i> $+\frac{2}{3}e$	c <i>charm</i> $+\frac{2}{3}e$	t <i>top</i> $+\frac{2}{3}e$
1964 - GIM mechanism	1974 - J/ψ meson	1995 - Tevatron
$4.7^{+0.5}_{-0.4} \text{ MeV}/c^2$	$4.7^{+0.5}_{-0.4} \text{ MeV}/c^2$	$4.7^{+0.5}_{-0.4} \text{ MeV}/c^2$
d <i>down</i> $-\frac{1}{3}e$	s <i>strange</i> $-\frac{1}{3}e$	b <i>bottom</i> $-\frac{1}{3}e$
1964 - GIM mechanism	1964 - kaon decay	1977 - Upsilon meson

Leptons

$0.511 \text{ MeV}/c^2$	$105.66 \text{ MeV}/c^2$	$1.77686 \text{ GeV}/c^2$
e <i>electron</i> $-1e$	μ <i>muon</i> $-1e$	τ <i>tau</i> $-1e$
1895 - X-rays	1936 - cosmic ray	1975 - SLAC
$< 2.2 \text{ eV}/c^2$	$< 0.17 \text{ MeV}/c^2$	$< 18.2 \text{ MeV}/c^2$
ν_e <i>electron neutrino</i> 0	ν_μ <i>muon neutrino</i> 0	ν_τ <i>tau neutrino</i> 0
1956 - nuclear reactor	1962 - Brookhaven	2000 - Fermilab

Gauge Bosons

0	0	$80.379 \text{ GeV}/c^2$
γ <i>photon</i> 0	g <i>gluon</i> 0	W W^\pm $\pm 1e$
1900 - Max Planck's quanta	1979 - PETRA	1983 - UA1, UA2
$91.1876 \text{ GeV}/c^2$	$125.10 \text{ GeV}/c^2$	
Z Z^0 0	H <i>Higgs</i> 0	2012 - ATLAS, CMS
1983 - UA1, UA2		

Figure 1.1: Particle of the Standard Model, listed together with their mass, electric charge, and the year and means of discovery (going clockwise from the top). VALUES NEED TO BE FIXED.

1.2 Coordinate systems and kinematic observables

Particles in HEP processes are described by their Lorentz-invariant four-vectors, $\mathbf{x} = (ct, x, y, z)$ and $\mathbf{p} = (E/c, p_x, p_y, p_z) = (E/c, \vec{p}_T, p_z)$, where $|\vec{p}_T| \equiv \sqrt{p_x^2 + p_y^2}$. In LHC experiments, the coordinate system is defined such that the x -axis points in the direction of the centre of the LHC, and the z -axis points in the direction of the beam, as shown in Fig. 1.2. In addition to the standard Cartesian coordinates, two observables, φ (azimuthal angle) and η (pseudorapidity), are used to describe the position and momentum of particles relative to the interaction point, which is located at $x = y = z = 0$. Pseudorapidity is defined as a function of the polar angle θ , where

$$\eta = -\ln(\tan(\theta/2)) . \quad (1.1)$$

For high-momentum particles ($E \simeq pc$), pseudorapidity is an approximation of the rapidity relative to the beam, given by

$$y = \frac{1}{2} \ln \frac{E + p_z c}{E - p_z c} . \quad (1.2)$$

Rapidity is a convenient quantity to use because it transforms additively under Lorentz boosts, unlike velocity. In these coordinates, the following relations hold:

$$p_x = |\vec{p}_T| \cos \varphi, \quad p_y = |\vec{p}_T| \sin \varphi, \quad p_z = |\vec{p}| \sinh \eta . \quad (1.3)$$

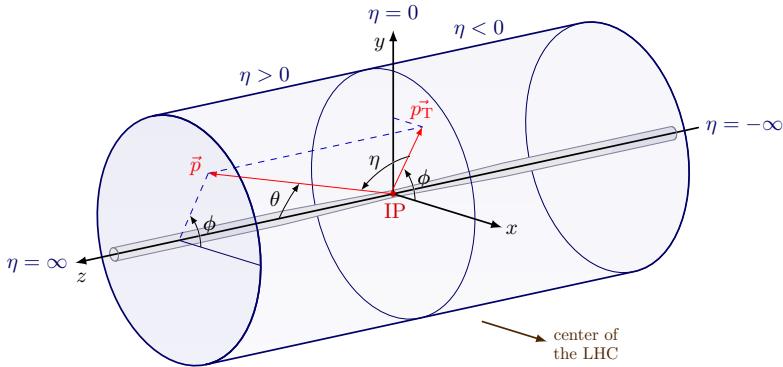


Figure 1.2: Coordinate system of an LHC experiment, with the interaction point in the centre.

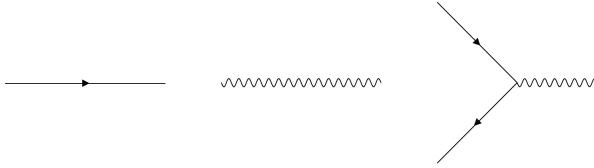
1.3 Quantum electrodynamics, electrons, and photons

In many aspects, QCD is a very similar theory to the simpler and better explored theory QED. In QFT, dynamics of particles can be provided in terms of its Lagrangian

density \mathcal{L} , from which equations of motions can be derived and which is also used to calculate interaction probabilities. The QED theory with a local U(1) symmetry has its \mathcal{L}_{QED} defined as:

$$\mathcal{L}_{\text{QED}} = \bar{\psi}(i\cancel{d} - m)\psi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} - e\bar{\psi}\cancel{A}\psi , \quad (1.4)$$

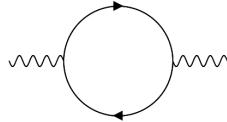
where ψ is the electron field with mass m and electric charge e , $F_{\mu\nu}$ the electromagnetic field-strength tensor, and Feynman slash notation is employed. The first part describes the dynamics of the electron fields, the second part describes the dynamics of the electromagnetic field, and the last part describes the interaction of electrons and photons with a coupling strength e . Using Feynman diagrams, they can be visualised as:



In QED interactions, each vertex depicted in the Feynman diagrams contributes to the probability of the process with a coefficient α (to the matrix elements as $\sqrt{\alpha}$, which is the coupling constant defined as $\alpha = e^2/4\pi$). This constant is generally small, which allows for interactions to be calculated using perturbation theory as an expansion series in α . The contributions to the series correspond to different Feynman diagrams representing the possible interaction processes, and they are ordered in powers of α based on the complexity of the diagrams.

Contributions from higher orders, such as the electron loop depicted below, lead to “screening” of the effective charge at large distances/small momenta, which makes the coupling constant dependent on the scale of the process μ . For example, at low energies corresponding to atomic scales, $\alpha \approx 1/137$, but at scales of the Z^0 boson mass, $\alpha \approx 1/127$ [6]. The *running* of this coupling is given by the β function, $\beta(\alpha) \equiv \frac{\partial \alpha}{\partial \log \mu^2}$, and it can be calculated by quantifying the effective coupling strengths at various orders of perturbation theory and using renormalisation group tools [7], although renormalisation is a more general concept. In QED, the screening leads to a positive sign in β calculated at lowest order, which means that when solving for α by integrating, α grows with energy scale¹.

¹The scale at which QED eventually breaks down due to this increase is well above the Plack mass.



Renormalisation is also used when calculating physical quantities where loop contributions lead to divergences, which are then absorbed into the parameters of the theory. The success of these procedures and the QED theory as a whole is validated by excellent prediction power for experimental measurements, such as the magnetic moment of an electron [8].

I.4 Quantum chromodynamics, quarks, and gluons

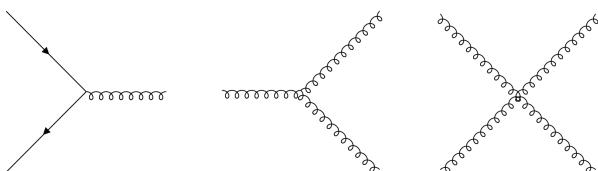
In QCD, particles have an additional quantum number called colour charge: red, green, and blue. Thus, there are three quarks of each flavour and eight gluons mediating interactions between. Gluons also carry colour charges, which allows them to interact with each other, making the theory non-Abelian. The QCD Lagrangian is symmetric under local SU(3) transformations and takes the shape of:

$$\mathcal{L}_{\text{QCD}} = \sum_f^{n_f} \bar{\psi}_i^{(f)} (iD_{ij}^\mu - m_f \delta_{ij}) \psi_j^{(f)} - \frac{1}{4} \sum_a^8 F_a^{\mu\nu} F_{\mu\nu}^a , \quad (\text{I.5})$$

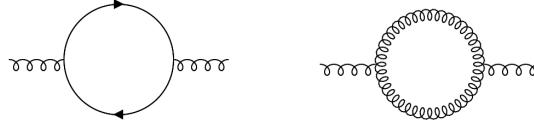
$$D_{ij}^\mu \equiv \partial^\mu \delta_{ij} + i g_s t_{ij}^a A_a^\mu , \quad (\text{I.6})$$

$$F_{\mu\nu}^a \equiv \partial_\mu A_\nu^a - \partial_\nu A_\mu^a - g_s f_{abc} A_\mu^b A_\nu^c , \quad (\text{I.7})$$

where ψ are the quark fields of n_f different flavours with mass m_f and colour i, j , D_{ij}^μ the covariant derivative with the coupling strength g_s and eight SU(3) generators given by matrices t_{ij}^a , and A_a^μ are the gluon fields [9]. Lastly, $F_{\mu\nu}^a$ is the field-strength tensor with f_{abc} being structure constants. The Lagrangian now also contains terms with interactions between gluons. In the representation of Feynman diagrams, for the interactions, there is:



Similarly to the QED case, the strong coupling constant can be defined as $\alpha_s = g^2/4\pi$. However, when considering its modifications due to virtual corrections, in addition to the quark loop, there is also a gluon loop.:



The gluon loop contributes to the β function in an opposite and larger way than the quark loop, and so overall, there is an *anti-screening* effect instead and a negative sign in the calculated one-loop β function. The running of the coupling can be calculated as:

$$\alpha_s(\mu) = \frac{1}{b_0 \log(\mu^2/\Lambda_{\text{QCD}}^2)} , \quad (1.8)$$

where b_0 is a constant computed from the loop calculations, $b_0 = \frac{11 - \frac{2}{3}n_f}{4\pi}$ [10]. The introduced Λ_{QCD} is a scale parameter of the theory corresponding to the energy where the coupling becomes infinite, and depends on the definition of α_s and the number of available quark flavours n_f [9]. It ranges between 200 and 300 MeV/ c^2 [11].

From the running, it is evident that the coupling strength decreases with increasing energy, which is known as *asymptotic freedom* [10, 12] and corresponds to the fact that strong interaction is short-ranged. On the other hand, at low values, α_s diverges, which is related to the fact that quarks are bound to hadrons – *quark confinement*. The evolution of α_s also limits the applicability of perturbation theory at low energy regimes; calculations from perturbative QCD (pQCD) are relevant at leading orders starting typically from $1 - 2$ GeV/ c^2 . The measured $\alpha_s(\mu)$ is shown in Fig. 1.3, and at scales of the Z^0 boson mass is approximately 0.1185 ± 0.0006 [13, 14].

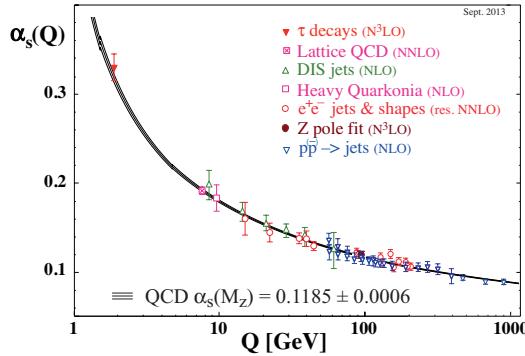


Figure 1.3: Strong coupling constant determined at different energy scales through various measurements and numerical calculations (data points) and compared with theoretical predictions from QCD. [13]

1.5 From partons to hadrons

1.5.1 Initial and Final State Radiation

In QFT, charged particles are surrounded by a cloud of virtual particles, which can be thought of as fluctuations in the particle's field. For example, the electron state can be described as a superposition of the bare electron plus additional massless bosons:

$$|e\rangle_{\text{phys}} = |e\rangle + |e\gamma\rangle + |e\gamma\gamma\rangle + \dots \quad (1.9)$$

and, at higher orders, pairs of virtual electrons. The fluctuations continuously form and recombine, with their lifetime depending on their energy and momentum. Specifically, the lifetime of a fluctuation with energy ω and transverse momentum k_T can be approximated as:

$$\tau \approx \frac{\omega}{k_T} \quad . \quad (1.10)$$

This implies that fluctuations with smaller- k_T live longer. [15]

As illustrated in Fig. 1.4, the coherent mixed state of the bare charge and the field fluctuations can be disturbed by the presence of an interaction. Intuitively, this interaction can change the energy and momentum of the fluctuations, their formation and recombination, and lead to the emission of radiation in two ways:

1. a fluctuation is kicked on-shell by the interaction and part of the field continues in its original direction, which leads to Initial State Radiation (ISR);

- as a result of the field of the scattered particle rearranging itself , which can be a source of Final State Radiation (FSR).

In both of the cases, a larger momentum transfer implies more radiation. *For hard, wide angle emissions, cross sections can be calculated perturbatively at fixed orders.*

Soft and collinear emissions, however, lead to infra-red divergences ($\propto \frac{1}{\omega}, \propto \frac{1}{k_T^2}$) and thus, need to be factorised away from the amplitudes or the cross sections and then described using resummation techniques. Without any emissions, the probabilities of finding electrons and photons of fractional momentum x with respect to the whole system are:

$$f_e(x) = \delta(1 - x), \quad f_\gamma(x) = 0, \quad (I.11)$$

When considering the emissions above some scales parametrised by the resolution parameter Q^2 , these probabilities, however, evolve according to the DGLAP equation [16] :

$$\frac{\partial}{\partial \ln Q^2} \begin{pmatrix} f_e(x, Q^2) \\ f_\gamma(x, Q^2) \end{pmatrix} = \frac{\alpha_{\text{em}}}{2\pi} \int_x^1 \frac{dz}{z} \begin{pmatrix} P_{ee}(z) & P_{e\gamma}(z) \\ P_{\gamma e}(z) & P_{\gamma\gamma}(z) \end{pmatrix} \begin{pmatrix} f_e\left(\frac{x}{z}, Q^2\right) \\ f_\gamma\left(\frac{x}{z}, Q^2\right) \end{pmatrix}, \quad (I.12)$$

where $P_{ij}(z)$ are the splitting probability functions of a particle i emitting a particle j .

In QCD, the behaviour is analogous, with $\alpha_{\text{em}} \rightarrow \alpha_s$, $e \rightarrow q$, and $\gamma \rightarrow g$. [15]

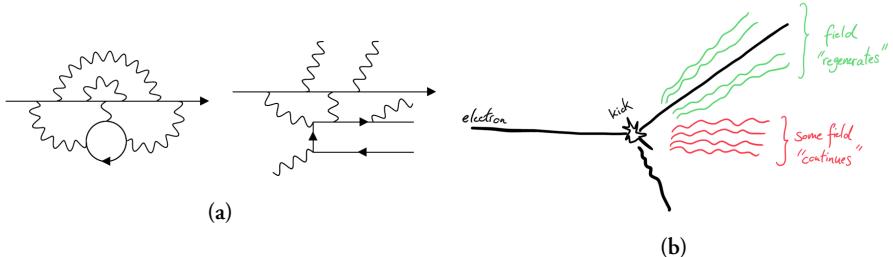


Figure 1.4: (a) Illustration of the field fluctuations before and after the state coherence gets disturbed by an external actor. (b) Illustration of emmisions of radiation in a scattering process.

I.5.2 Factorisation theorem

The evolution equation (I.12) implies that the probabilities of observing emissions with a fractional momentum x depend on the resolution Q^2 . In QCD,

1. when applied to the initial state, they are known as parton distribution functions (PDFs) and determine the probabilities of finding partons² in the composite hadronic state.
2. When applied to the final state, they are called fragmentation functions, and determine the probabilities of measuring fragments of the outgoing particles.

This leads to the factorisation theorem [17] for processes involving collisions of two hadrons, which separates the perturbatively calculable partonic cross section from the non-perturbative partonic evolution and hadronisation. The theorem can be expressed as follows:

$$\sigma = f_i^A(x_i, \mu_F) f_j^B(x_j, \mu_F) \otimes \hat{\sigma}_{ij \rightarrow n}(\mu_F, \mu_R) \otimes D_{n \rightarrow n'} . \quad (1.13)$$

Here, i and j are the initial partons, $\hat{\sigma}_{ij \rightarrow n}$ is the partonic cross section, $D_{n \rightarrow n'}$ is the process-specific fragmentation function for evolving the partons n into the particles' final state n' , and μ_F and μ_R are the factorisation and renormalisation scales, respectively. The factorisation scale, μ_F , determines the scale below which the emissions are absorbed into the PDFs. The theorem is depicted in Fig. 1.5.

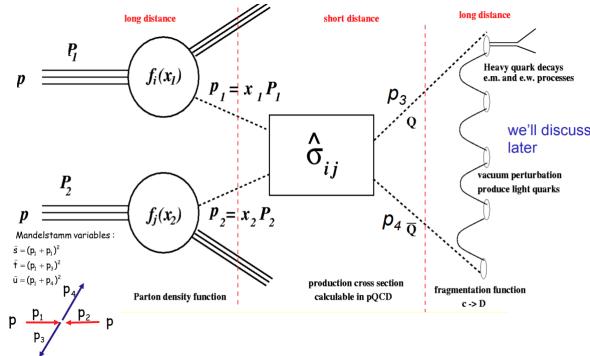


Figure 1.5: Illustration of the factorisation theorem. (NEEDS TO BE REMADE).

1.5.3 Parton distribution functions

The PDFs defining the probabilities of finding quarks and gluons in nucleons can be determined experimentally at hadron-electron colliders such as HERA [18]. They are determined from measurements of deep inelastic scatterings in a range of energies and momentum transfers. They are displayed in Fig. 1.6 as a function of the fractional momentum x (also called Björken x).

²Partons refer to the valence quarks, sea quarks, and gluons inside hadrons.

According to collider kinematics, $x \propto \frac{1}{\sqrt{s_{ey}}}$, therefore, the partonic composition of ultra-relativistic hadrons is dominated by gluons. Following unitarity principles and BK evolution equation [19], it is expected that gluons start recombining and the gluonic content saturates as $x \rightarrow 0$. This is actively researched [20], however, not directly measured yet. Additionally, it should be noted that in ultra-relativistic heavy nuclei, the partons are modified in the contracted nuclear environment and the PDFs are referred to as nPDFs.

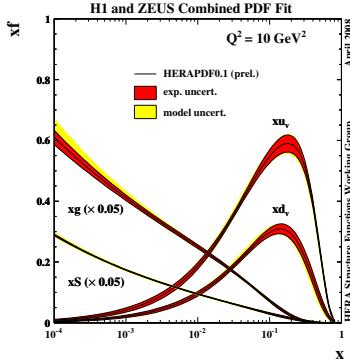


Figure 1.6: Parton distribution functions determined in ep scatterings at HERA as a function of the fractional momentum for the up, down, sea quarks, and gluons. [18]

1.5.4 Parton fragmentation and the Lund string

After the scattering process, the produced partons continue to fragment by emitting more partons in a process called the parton shower. Since the coupling strength in QCD increases with decreasing the energy scale of the splitting, this leads to the production of many soft, collimated emissions known as jets. The partonic evolution continues until the virtuality of the partons reaches the hadronization scale ($\approx \Lambda_{\text{QCD}}$). There are multiple frameworks within QCD to describe the evolution of partons into their final state, such as using the DGLAP equations or the so-called dipole formalism.

Once the partonic final state is reached, the partons hadronise into the observable mesons and baryons. The hadronisation process is not calculable in QCD and requires phenomenological models to describe it. One such model is the Lund string model [21], which describes hadronisation as the breaking of a color string between the quarks in the final state. In this model, the energy stored in the color string is converted into the mass of new hadrons.

According to confinement, hadronisation should involve at least two partons with

complementary colours. In QCD, the $q\bar{q}$ potential takes the shape of

$$V_{q\bar{q}} \approx -\frac{4}{3} \frac{\alpha_s \hbar c}{r} + \kappa r , \quad (1.14)$$

where κ is a parameter with value around $1\text{GeV}/\text{fm}$. In the non-perturbative regime (long distances), the potential is dominated by the linear part, which is reminiscent of a system bound by a string with tension κ . This is taken advantage of by the Lund string model – a q and \bar{q} pair separated by distance Δx is bound by a color field (string) with energy $\kappa\Delta x$.

If the q and \bar{q} continue separating as a result of the scattering, the energy stored in the color field increases. At some point, it can become energetically favourable to produce a new $q\bar{q}$ pair out of vacuum, which is a quantum mechanics tunnelling phenomenon characterised by the probability:

$$\frac{dP}{dm_T} \propto \exp\left(-\frac{\pi m_T^2}{\kappa}\right) , \quad (1.15)$$

where m_T is the transverse mass of the produced quarks. Otherwise, the $q\bar{q}$ system starts contracting and oscillates with a period $T = 2E_{\text{kin}}/\kappa$, where E_{kin} is its maximum kinetic energy. The produced q and \bar{q} then connect by new color fields to the original pair. This process repeats itself resulting in a cascade of many $q\bar{q}$ pairs connected by many color strings. In this description, baryons can also be created by double tunnelling of a $qq\bar{q}\bar{q}$ pair. The process is illustrated in Fig. 1.7.

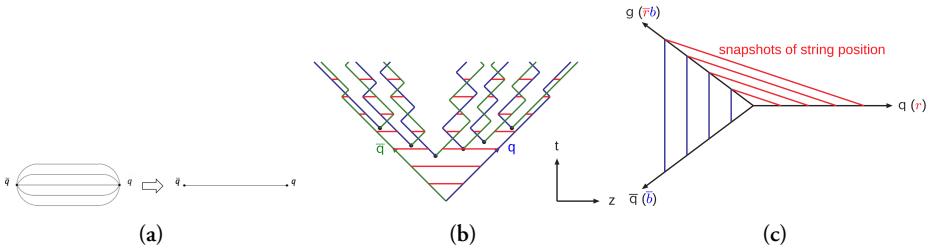


Figure 1.7: (a) Illustration of the color field between two quarks and its simplified representation with a string [21]. (b) Illustration of the string splitting by producing new $q\bar{q}$ in the $t-z$ plane [21]. (c) Visualisation of the treatment of gluons in the Lund string model [22].

Equation 1.15 also implies that production of strange quarks is suppressed by a factor of

$$\rho = \exp\left(-\frac{\pi(m_s^2 - m_{u,d}^2)}{\kappa}\right) . \quad (1.16)$$

This parameter is typically tuned to data, as substituting constituent ($m_s \approx 0.5 \text{ GeV}/c^2$, $m_{u,d} \approx 0.33 \text{ GeV}/c^2$) versus current masses ($m_s \approx 0.1 \text{ GeV}/c^2$, $m_{u,d} \approx 0$) leads to considerable differences underestimating and overestimating data, respectively.

For a $q\bar{q}g$ system, in this model, the gluon connects to the quark and antiquark and is effectively treated as a “kink” on the color field, adding energy and momentum to the $q\bar{q}$ string (stretching it in its direction), as visualised in Fig. 1.7.

It should be noted that in the paradigm of AA collisions, hadron production can be alternatively modelled by hadronisation at the QGP’s phase boundary by *coalescing* free quarks.

1.6 Multiple partonic interactions

Results from Sp \bar{p} S in the 1980s sparked motivations for considering interactions of multiple partons between the two composite protons. For example, the AFS experiment observed an abundance of 4-jet events, displayed in Fig. 1.8, that could not be explained by calculations considering a double gluon bremsstrahlung from a single partonic scattering[23]. Furthermore, UA5 measurements studying energy dependence of multiplicity distributions $P(N_{\text{ch}})$ saw the so-called KNO scaling[24], where $P(N_{\text{ch}})/\langle N_{\text{ch}} \rangle$ does not depend on energy, but revealed a broadening in high-multiplicity events with increasing \sqrt{s} [25, 26], which was not reproducible in the context of N_{ch} being produced from a single string [27]. This further suggested the presence of multiple production sources.

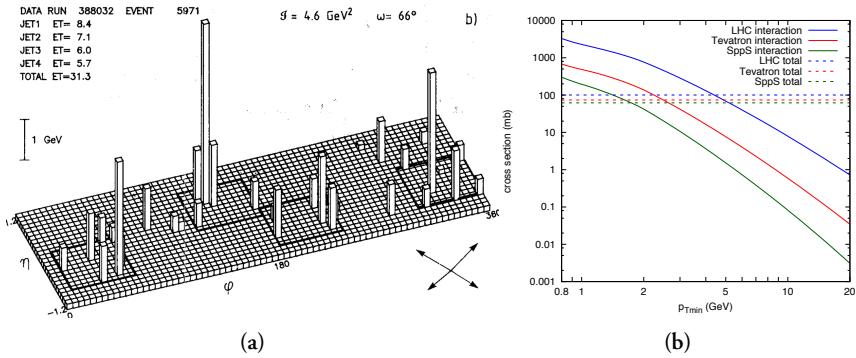


Figure 1.8: (a) Event display of an event with a 4-jet, where the pillars correspond to transverse energy deposits. [23] (b) Dependence of the integrated parton-parton cross section on the cutoff parameter $k_{\perp\min}$ for Sp \bar{p} S at $\sqrt{s} = 0.63 \text{ TeV}$, Tevatron at $\sqrt{s} = 1.96 \text{ TeV}$, and the LHC at $\sqrt{s} = 13 \text{ TeV}$, modelled with Pythia. [27]

These findings prompted further development of Regge theory and approaches that incorporated multiple pomerons, which were successful in describing the N_{ch} distributions. However, this approach is fully decoupled from descriptions of the perturbative primary scattering. Subsequently, much of the phenomenology related to multiple partonic interactions was developed within the framework of the Pythia MC event generator, which is discussed individually in Chapter ?? [27]. However, nowadays, the relevance of the concept of MPIs in hadronic collisions extends beyond this generator. A scattering with double partonic interactions is illustrated in Fig. 1.9.

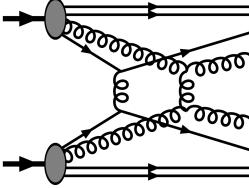


Figure 1.9: Diagram showing a double partonic interaction, a case of $n_{\text{MPI}} = 2$. [15]

In the Pythia approach, MPI are treated as additional perturbative scatterings. In QCD, the $2 \rightarrow 2$ cross section (dominated by the gluon exchange t-channel) diverges as $\propto \alpha_S^2(k_\perp^2)/k_\perp^4$, so a cutoff parameter $k_{\perp\min}$ must be introduced, and using (??) leads to:

$$\frac{d\sigma}{dk_\perp^2} = \sum_{ij} \int dx_1 dx_2 f_i(x_1, \mu_F^2) f_j(x_2, \mu_F^2) \frac{d\hat{\sigma}_{ij}^H}{dk_\perp^2}, \quad (1.17)$$

$$\sigma_{\text{int}}(k_{\perp\min}) = \int_{k_{\perp\min}^2}^{s/4} \frac{d\sigma}{dk_\perp^2} dk_\perp^2. \quad (1.18)$$

The choice of cutoff can be tuned to experimental data, and for the Sp \bar{p} S energy of $\sqrt{s} = 630$ GeV, a value of around 1.6 GeV/c was typical [27]. The dependence of this parton-parton scattering cross section is shown in Fig. 1.8.

The total pp cross-section, which is on the order of 100 mb at $\sqrt{s} = 13$ TeV, is given by

$$\sigma_{\text{pp}} = \sigma_{\text{elastic}} + \sigma_{\text{single dif.}} + \sigma_{\text{double dif.}} + \sigma_{\text{non-dif.}}, \quad (1.19)$$

where the inelastic cross sections $\sigma_{\text{inel}} \approx \sigma_{\text{double dif.}} + \sigma_{\text{non-dif.}}$ corresponds to approximately 60% of the total. The mean number of MPIs, $\langle n_{\text{MPI}} \rangle$, can be estimated using:

$$\langle n_{\text{MPI}} \rangle(k_{\perp\min}) = \frac{\sigma_{\text{int}}(k_{\perp\min})}{\sigma_{\text{inel}}} \quad (1.20)$$

However, the actual treatment is more complex and involves considerations of other parameters such as the dampening factor k_\perp^0 to account for the confinement nature of partons, modifications of multiparton PDFs, energy-momentum conservation effects, x -dependent source geometry, and the intertwinedness of partonic evolutions.

In summary, MPIs represent several subcollisions that take place in an average pp collision with p_T scales of a few GeV. They are colour-connected to the beam remnants, which in the Lund model are represented by strings. Since a string with $\kappa = 1\text{GeV}/\text{fm}$ yields, as a rule of thumb, approximately one hadron per unit rapidity, and the average pp collision at the LHC at $\sqrt{s} = 13\text{ TeV}$ has $\langle dN_{\text{ch}}/dy \rangle \approx 6$, the typical number of partonic interactions is around six [27].

Finally, the observation of QGP-like phenomena in pp collisions at the LHC has renewed interest in MPI phenomenology, as discussed in the following chapter. Such observations do not contradict the concept of MPIs; rather, they suggest the possibility of incorporating collective behavior among the MPIs, such as interactions between strings, local modifications of string tensions, or, alternatively, the formation of a multipartonic state with QGP-like properties.

1.6.1 Colour reconnection

The incorporation of MPIs improved the description of the N_{ch} distributions and their dependence on \sqrt{n} . However, there were also observations of $\langle p_T \rangle(N_{\text{ch}})$ increasing as a function of N_{ch} , which could not be explained. More MPIs lead to more strings, which in turn leads to the production of more particles, but the p_T is mostly unaffected. This would predict a weaker dependence of $\langle p_T \rangle$ on N_{ch} , contrary to the data [27]. The issue was resolved by implementing a possible color reconnection mechanism, which rearranges the color fields between partons.

TBA Insert diagrams of the processes!

One can envision the following process:

$$e^+ e^- \rightarrow W^+ W^- \rightarrow q_1 \bar{q}_2 q_3 \bar{q}_4.$$

In this scenario, a color reconnection mechanism could rearrange the colour-connected $q_1 \bar{q}_2$ and $q_3 \bar{q}_4$ into $q_1 \bar{q}_4$ and $q_3 \bar{q}_2$ if it were energetically favourable, depending on the phase-space configurations. Measurements at LEP [28] of this process have indeed shown that such final-state corrections must be taken into account to explain the data on W masses and widths. They also reported that the reconnection probabilities for such events are on the order of 50%, further indicating that colour reconnection is an important factor to consider.

Pythia implements CR by minimizing the total length of strings in the system, analogous to minimising potential energy [29]. This mechanism, illustrated in Fig. 1.10, explains the rising trend of $\langle p_T \rangle$ as a function of N_{ch} : shorter strings imply fewer hadrons to split the transverse boost across, and the more MPI, the bigger this effect. Moreover, CR also helped describe the absolute value of $\langle p_T \rangle$. With this approach, no further modifications of fragmentation parameters were necessary, in line with the concept of jet universality. However, it should be noted that there are various CR implementations and all rely on parameters obtained from tuning to data.

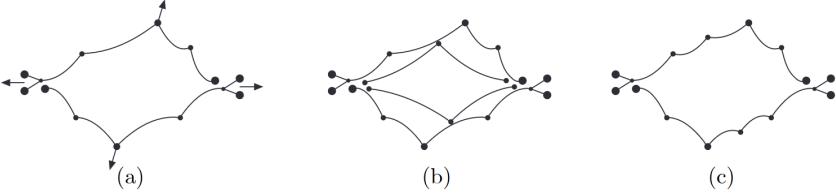


Figure 1.10: Depiction of the CR process: (a) in a hard parton subcollision, the outgoing gluons are connected to the beam remnants through colour. Additional gluon kinks may occur through initial state radiation, which are ordered by rapidity. (b) A second hard scattering should theoretically result in two new strings connected to the remnants. (c) In order to minimise the total string length, gluons are colour reconnected. [30]

It is also worth noting that the p_T boost acquired through color reconnection may depend on mass and whether a hadron is a baryon or meson, which somewhat mimics the hydrodynamic signatures of collective flow observed in AA collisions [31].

1.7 Underlying event

The underlying event (UE) in high-energy collisions refers to the additional hadronic activity that accompanies the primary hard scattering process, but is not directly related to it. This includes the fragmentation products of the beam remnants, ISR and FSR, as well as the effects of the previously discussed MPIs, and is visualised in Fig. 1.11. The UE (with its name coined at Tevatron [32]) is typically characterized by a distribution of softer particles around and far outside of the hard process and was first observed at SppS in the 1980s [33]. These measurements saw a constant plateau of transverse energy E_T outside of the jet core, with its height independent of the jet energy.

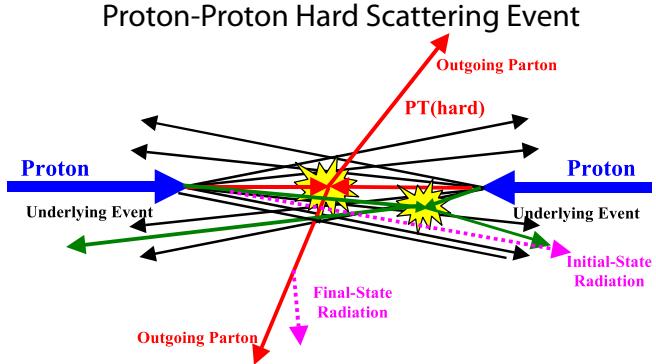


Figure 1.11: Cartoon illustrating a pp collision and its components: the hard scattering process, beam remnants, initial/final state radiation, and the MPIs. The last three contribute to the underlying event. [32, 34]

It is important to note that particle production in UE is different from the MB production, as it is biased by the presence of hard scattering. Additionally, the magnitude of the UE can fluctuate significantly from event to event.

1.8 From hadrons to partons: deconfined QCD matter

A simple argument can be made that confinement of quarks inside hadrons cannot be sustained when the density of partons is too large compared to that inside ordinary hadrons. Or alternatively, when the partonic kinetic energies are much larger than the confining part of the $q\bar{q}$ potential in (1.14). The following sections introduce some theoretical frameworks predicting the deconfinement of quarks and gluons. The calculations presented use natural units ($\hbar = c = k_B = 1$).

1.8.1 Bag model of hadrons

In this very naive approach [35, 36], hadrons are treated as spherical cavities (“bags”) with radius R of free massless quarks. These cavities exist in the non-perturbative QCD vacuum, which exerts a confining pressure B . The lowest-energy solution of the Dirac equation for the quarks, which in this case is $\not{p}\psi = 0$, is the $s_{1/2}$ -state given by:

$$\psi(r, t) = N \begin{pmatrix} j_0(\omega r)U \\ i\sigma \cdot \hat{r} j_1(\omega r)U \end{pmatrix} \exp(-i\omega t) , \quad (1.21)$$

where N is a normalisation constant, j_0 and j_1 are spherical Bessel functions, ω is the quark energy, and U the two-component spinor. The assumption that quarks are confined within the cavity volume represents the boundary conditions that the quark scalar density $\psi\bar{\psi}$ becomes zero at $r = R$, which is equivalent to $j_0(\omega R) = j_1(\omega R)$:

$$\hat{j}_0^2(\omega R) - \vec{\sigma} \cdot \hat{r} \vec{\sigma} \cdot \hat{r} j_1^2(\omega R) = 0 \quad (1.22)$$

$$j_0(\omega R) = j_1(\omega R) , \quad (1.23)$$

which happens when $\omega \approx \frac{2.04}{R}$. Thus, energy of the system can be given by

$$E(R) \approx n_q \cdot \frac{2.04}{R} + \frac{4\pi}{3} BR^3 . \quad (1.24)$$

Here, the first term represents the kinetic energy of n_q quarks in the cavity and second term is the cavity volume energy. Gluon solutions should also be considered but are neglected in this approach. The first term acts to expand the cavity, whereas the second term acts to contract it. Finding an optimum of this energy with respect to R leads to

$$B^{1/4} \approx \left(\frac{2.04n_q}{4\pi} \right)^{1/4} \frac{1}{R} . \quad (1.25)$$

Finally, assuming values for a proton, $R \approx 0.8$ fm, and three valence quarks $n_q = 3$, the confining pressure can be approximated as $B^{1/4} \approx 206$ MeV. [35]

To relate the confining pressure to a critical temperature at which deconfinement occurs, T_c , one can assume a gas of relativistic massless fermions and bosons with energy density ρ [37?]. Using Stefan-Boltzmann law, the equation of state is,

$$P = \frac{1}{3}(d_b + \frac{7}{8}d_f)\rho = (d_b + \frac{7}{8}d_f)\frac{\pi^2}{90}T^4 \quad (1.26)$$

where d_b and d_f are the degeneracy numbers for bosons and fermions, in this case gluons and quarks, respectively. Summing together possible colours, polarisations, and flavours for particles and antiparticles, one gets $d_b = 16$ and $d_f = 24$. Inserting the cavity pressure B value calculated in (1.25), T_c can be estimated as approximately 145 MeV.

1.8.2 Lattice QCD

Lattice QCD (LQCD) is a technique allowing calculation of processes involving the strong interaction in the non-perturbative regime, from first principles, without phenomenological assumptions. In this approach, the space-time continuum is discretised into a four-dimensional lattice, which allows QCD path integrals to be solved

numerically. Smallest squares on the lattice are called *plaquettes*, with the lattice links representing gluon fields and lattice sites representing quark fields. [38]

Lattice calculations are computationally extremely intensive³, thus, a sufficiently coarse lattice spacing must be chosen to reduce computational costs and make the approach feasible. Often, simulations are also performed with unphysical quark masses (e.g. $m_q \sim m_\pi$), for the same reason. The results are then extrapolated using highly complex methods. Another limitation in LQCD is the so-called *sign problem*, discussed in detail for instance in Ref. [40], which arises when evaluating highly oscillatory complex integrals in finite-density environment.

Despite its challenges, LQCD has been successful in various predictions, notably, the *ab initio* calculation of the mass of neutron (uud) using the mass of Ω (sss) as an input [41]. Furthermore, it allows for the calculation of thermodynamics of QCD matter and predicts its equation of state, as well as a phase transition at $T_c \approx 150$ MeV [42, 43], which is usually identified with the formation of quark-gluon plasma QGP, discussed in Chapter ???. The dependence of energy density and pressure on temperature can be seen in Fig. 1.12.

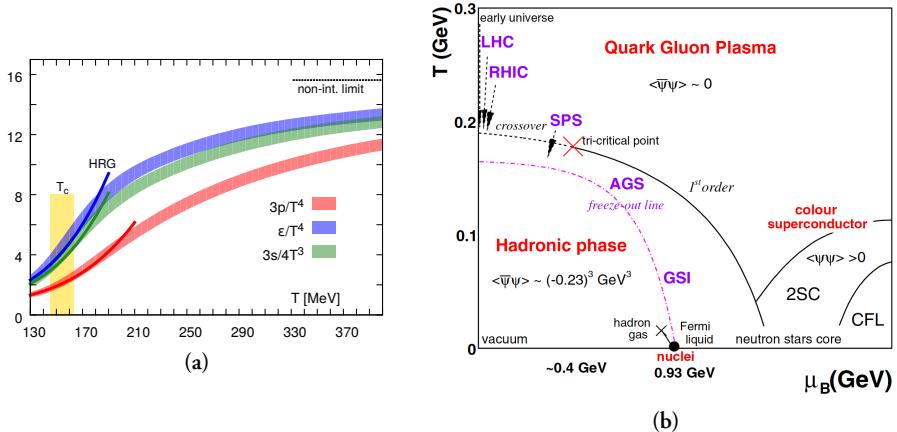


Figure 1.12: (a) Dependence of the pressure (red), energy density (blue), and entropy density (green) on temperature, determined with LQCD. [43] (b) Phase diagram of QCD matter. [44]

1.8.3 QCD phase diagram, chiral symmetry restoration

Phase transitions of QCD matter are investigated to explore the QCD phase diagram with respect to temperature T and baryon chemical potential μ_B , which corresponds

³In fact, in the past, LQCD was among important drivers of the advancement of GPU computing and it is also used as benchmark in high-performance computing [39].

to net baryon density. Figure 1.12 visualizes the different areas of the QCD phase diagram probed by various experiments [44]. Measurements of Pb-Pb collisions at the LHC access high T and almost zero μ_B , as the nucleons of ultrarelativistic Pb nuclei escape the interaction volume before the plasma develops, and the high energy subsequently leads to a sizable baryon production balanced by anti-baryons due to conservation laws. Experiments studying collisions of slower and heavier nuclei reach higher μ_B regions.

Furthermore, looking back at the QCD Lagrangian in (1.5) and neglecting quark current masses $m_{u,d} \rightarrow 0$, it can be seen that it is invariant when switching the up and the down quark, corresponding to a SU(2) isospin symmetry. This allows the fermion fields to be rewritten in terms of their left and right chiralities and the QCD Lagrangian then exhibits a larger *chiral symmetry*.

However, it is known that the vacuum expectation value of a $q\bar{q}$ state is much larger than the current masses $m_{u,d}$:

$$\langle 0 | q\bar{q} | 0 \rangle = \langle 0 | u\bar{u} + d\bar{d} | 0 \rangle \approx (250 \text{ MeV})^3 \quad , \quad (1.27)$$

where the value is taken from average masses of light flavour mesons. This means that the QCD vacuum spontaneously breaks the chiral symmetry. [45]

It can be expected that in the plasma of deconfined quarks and gluons, chiral symmetry is *restored* [46, 47]. This is actively studied in AA collisions, for example in searches of the so-called chiral magnetic effect [48] or degeneracy of normally chiral partners [49], such as the ρ ($J^P = 1^-$) and a_1 ($J^P = 1^+$) states.

Chapter 2

QCD phenomena in high energy hadronic collisions

The aim of this chapter is to give an introduction to the physics of heavy ions and the various phenomena related with the quark-gluon plasma QGP. Furthermore, a detailed summary of the findings of QGP phenomena in small systems, i.e. pp and pA collisions, is given. Lastly, some Monte Carlo event generators based on phenomenological modelling of hadronic collisions relevant to this dissertation are summarised.

2.1 Collisions of heavy nuclei

2.1.1 Collision geometry, centrality, multiplicity

Collisions of heavy nuclei, composed of many fluctuating nucleons, may occur under various initial state configurations. Some quantities used to describe them are the impact parameter b , defined as the distance between the two nuclei centers, number of participating (scattered) nucleons N_{part} , and the number of binary nucleonic collisions N_{coll} .

Determining these quantities is important because:

1. Soft processes, such as light flavor particle production, are expected to scale with the interaction volume, which $\propto N_{\text{part}}$.
2. Hard processes, such as jet and heavy flavor production, are expected to scale with the number of large momentum transfer interactions given by N_{coll} .

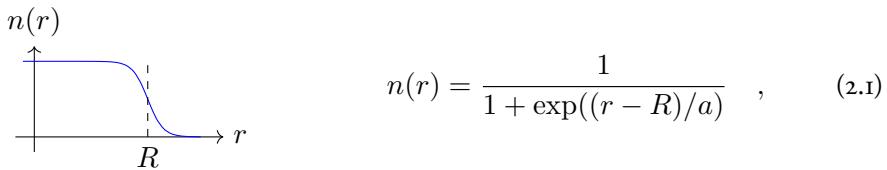
3. b , disregarding the fluctuations of nucleonic positions, defines the shape and anisotropy of the overlap region, which are important initial state conditions.

Since these quantities cannot be directly measured, they need to be modelled. The charged particle *multiplicity* is commonly used for this purpose, as $\langle N_{\text{ch}} \rangle$ increases monotonically with N_{part} , N_{coll} , and decreasing b . Multiplicity N_{ch} can be measured experimentally, e.g. with tracking detectors. The concept of *centrality* is also used, which is defined as quantiles of the total nuclear cross-section. For example, a centrality of 0 – 5% refers to low b values and the top 5% of N_{ch} values (central events), while 95 – 100% centrality refers to high b values and the bottom 5% of N_{ch} values (peripheral events). Centrality can also be inferred from other *event activity* classifiers, such as amplitudes of scintillators at forward rapidity, transverse energy in calorimeters, or energy from beam remnants in zero-degree-calorimeters.

In AA collisions, these relationships are well-defined, and thus the models perform well. The most popular model is the MC Glauber model [50]. Other models include MC-KLN [51] and IP Glasma [52].

2.1.2 MC Glauber model

The MC Glauber model [50] takes on a very simple albeit powerful approach. The two nuclei are simulated in three dimensions in a way that satisfies their respective nuclear density profiles, usually modelled by sampling the positions of nucleons from the Woods-Saxon distribution:



where R is the nuclear radius and a the nuclear skin thickness.

The nucleonic densities can be represented by uniform disks, or more accurately by Fermi-distributions or Gaussian profiles to account for fluctuations of their densities. Their parameters are left free and are tuned to the data.

A random impact parameter is then chosen or sampled. The collision is then treated as a sequence of independent binary nucleon-nucleon collisions, where

- i. nucleons remain travelling in straight lines,

2. the inelastic nucleon-nucleon cross section σ_{NN} does not depend on the number of interactions,
3. two nucleons are considered to interact if their transverse relative distance $d \leq \sqrt{\sigma_{\text{NN}}/\pi}$.

Figure 2.1 illustrates an example of a Glauber Monte Carlo event for a Au+Au collision. By simulating numerous collisions, the average N_{part} and N_{coll} are determined¹, and their relations to centrality and event activity observables are determined by fitting to experimental data.

Recent studies have extended the MC Glauber model to include sub-nucleonic structures. Such efforts show that the production of charged hadrons at mid-rapidity scales linearly with the number of participating partons. Comparisons with LHC data at $\sqrt{s_{\text{NN}}} = 5.02 \text{ TeV}$ suggest that the number of sub-nucleonic degrees of freedom ranges from 3 to 5 [53].

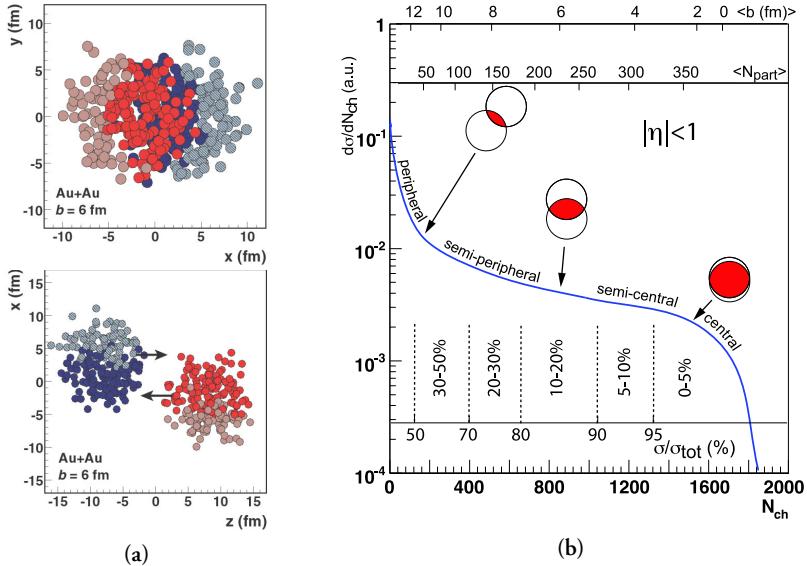


Figure 2.1: (a) Glauber Monte Carlo event of a Au+Au collision at $\sqrt{s_{\text{NN}}} = 200 \text{ GeV}$ shown in the transverse plane (top panel) and along the beam axis z (bottom panel). The darker circles represent participating nuclei and their area corresponds to σ_{NN} . (b) Illustrative diagram of the dependence of final-state event multiplicity on the initial-state quantities, number of participants N_{part} and impact parameter b . [50]

¹It also shows the scaling between the numbers of participants and binary collisions, which is approximately $N_{\text{coll}} \approx 0.35 N_{\text{part}}^{4/3}$.

2.2 Quark-gluon plasma

In agreement with lattice QCD predictions [54], the QGP has been measured in ultra-relativistic collisions of heavy nuclei at RHIC [55, 56], LHC [57], and even SPS [58]. Although it cannot be observed directly, a wealth of evidence from three decades of research combining various observables reveals the effects of the produced QGP medium. Whilst somewhat context-dependent, the following features make QGP the most extreme phenomenon observed in terms of its:

- *Temperature*: QGP temperatures reach values on the order of hundreds of MeV, which corresponds to approximately 2×10^{12} K.² [59]
- *Viscosity*: the shear viscosity to entropy density ratio η/s reaches the minimum quantum limit of $1/4\pi$ (for water at room temperature, $\eta/s \approx 8$), making it an almost perfect liquid. [44]
- *Vorticity*: in semi-peripheral collisions, the rotating plasma reaches a vorticity of approximately $9 \times 10^{21} \text{ s}^{-1}$. [60]
- *Magnetic field*: in non-central collisions, the magnetic fields in the system may peak at $\sim 10^{19}$ T. [61]

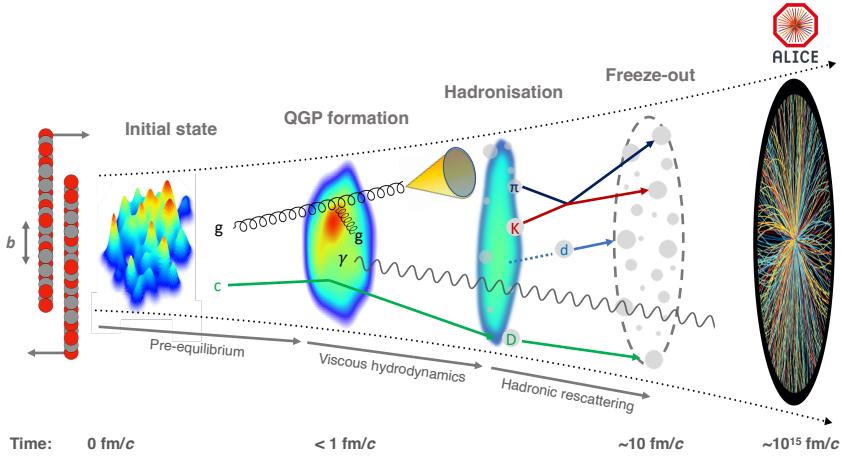


Figure 2.2: Evolution of heavy nuclei collisions at LHC energies, depicting the different stages. [62]

²Contrasting some of the lowest temperatures required for the super-conducting magnets of the LHC, $T \approx 1.9$ K.

Figure 2.2 illustrates the common paradigm for the different stages of a collision between heavy nuclei:

1. The Lorentz-contracted heavy nuclei approach each other at ultra-relativistic speeds.
2. *Pre-hydrodynamisation stage* ($\tau \equiv \sqrt{t^2 - z^2} \leq 1$ fm/c): “hard” particles are produced in scatterings with the highest momentum transfer Q^2 . Produced matter expands rapidly in the longitudinal direction and starts expanding in the radial direction.
3. *Hydrodynamisation* ($1 \leq \tau \leq 10$ fm/c): abundantly produced partons create a deconfined medium, which can be described by hydrodynamic equations.
4. *Chemical freeze-out* ($\tau \sim 10$ fm/c): the system cools down and hadronises. The produced hadrons then stop interacting inelastically and the system’s chemical content is stabilised.
5. *Kinetic freeze-out* ($\tau \lesssim 20$ fm/c): hadrons no longer interact elastically and their kinematics stabilize.
6. Long-lived particles can be measured in the detector volume.

The following subsections outline some of the essential phenomena related to the production of QGP.

2.2.1 Quarkonium dissociation and sequential suppression

Heavy quarkonia are vector meson states consisting of $c\bar{c}$ and $b\bar{b}$. They include J/ψ , $\psi(2S)$, $\Upsilon(1S)$, $\Upsilon(2S)$, $\Upsilon(3S)$, which can be relatively easily measured in LHC experiments via their di-lepton decay channels. They are created solely in the first phases of the collision and then experience the entire evolution of the QGP medium:

$$t_{\text{creation}}^{Q\bar{Q}} < t_{\text{creation}}^{\text{QGP}} < t_{\text{lifetime}}^{\text{QGP}} \ll t_{\text{lifetime}}^{Q\bar{Q}} . \quad (2.2)$$

Additionally, due to their large binding energies, their radii may remain smaller than the plasma screening radius $r_D(T)$, and thus, survive the dissociation. For instance, considering their in-vacuum radii determined from the $q\bar{q}$ potential, $r_{\Upsilon(1S)} \sim 0.14$ fm, $r_{\Upsilon(2S)} \sim 0.28$ fm, $r_{\Upsilon(3S)} \sim 0.39$ fm, which contrast the $r_\pi \sim 0.7$ fm [63]. This implies that different temperatures result in the dissociation of different states, and measuring the production of different states can help infer the QGP temperature, as illustrated in Fig. 2.3 [64, 65].

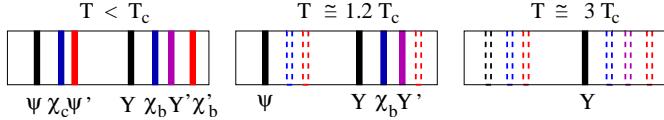


Figure 2.3: Spectral lines corresponding to various charmonium and bottomonium states for different medium temperatures , relative to the QGP critical temperature. [65]

The production of heavy quarkonia in AA collisions is compared to that in pp collisions (scaled by the average number of binary collisions) through the nuclear modification factor, R_{AA} . This quantity is widely used in various other measurements and is defined as:

$$R_{AA} = \frac{dN_{AA}/dp_T}{\langle N_{coll} \rangle dN_{pp}/dp_T} . \quad (2.3)$$

R_{AA} can take on the following values:

1. $R_{AA} = 1$: The result one would expect if the AA collision is a mere superposition of nucleon-nucleon collisions. There is no net effect on the production, corresponding to the absence of the QGP medium and other nuclear effects, or their mutual cancellation.
2. $R_{AA} < 1$: The production is overall suppressed, for example, due to dissociation.
3. $R_{AA} > 1$: The plasma and nuclear effects systematically enhance the measured production.

At LHC energies, the abundance of charm quarks in the QGP is high enough that charmonia can be formed after dissociation, which somewhat complicates the interpretation of charmonia suppression. However, the $\Upsilon(3S)$ bottomonium has R_{AA} consistent with zero at $\sqrt{s_{NN}} = 5.02$ TeV, as shown in Fig. 2.4 [59]. This complete suppression is a clear signature of the QGP and can be used together with models to estimate the QGP temperature at these energies as $T \approx 630$ MeV [66].

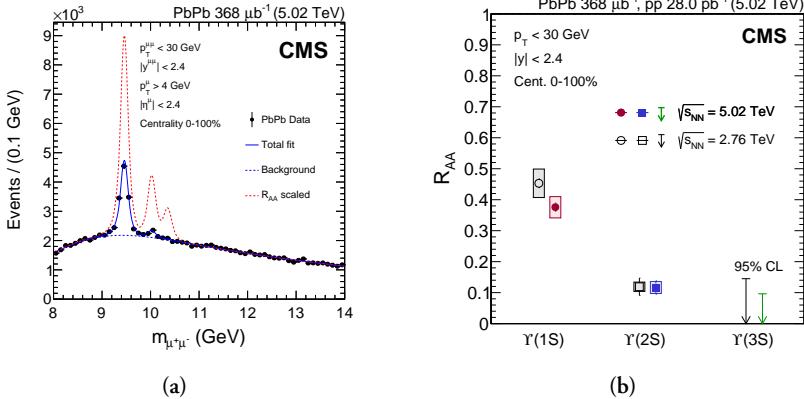


Figure 2.4: (a) Muon invariant mass distributions in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, also showing a scaled result from pp collisions at same energies (red dashed line). (b) Nuclear modification factors of the three bottomonium states. [59]

2.2.2 Strangeness enhancement

In the production of hadrons in vacuum, strangeness is suppressed relatively to light quarks not only due to the higher mass of the strange quark ($m_s \approx 0.1 \text{ GeV}/c^2$), but also due to the much higher constituent mass ($m_K \approx 0.5 \text{ GeV}/c^2$). However, in the QGP, due to the high gluon densities and $T \sim m_s$, strangeness production may equilibrate with u and d quarks through gluon fusion:

$$gg \rightarrow s\bar{s} \quad .$$

This phenomenon was proposed as one of the first signatures of QGP observation in colliders [67, 68]. Indeed, an enhancement in the production of strange hadrons is observed in AA collisions, which is dependent on the event activity and increases with increasing strangeness content of the hadron [69]. Figure 2.5 displays these results.

Furthermore, the yields of hadrons measured in AA collisions can be accurately described by statistical models [70, 71] which, generally, assume that the dense system is in thermal and chemical equilibrium at the point of its freeze-out. In these models, strangeness is assumed to be conserved on average, which corresponds to a grand-canonical ensemble with a strange chemical potential μ_S .

In small systems, the conservation of strangeness must be taken into account for each interaction, locally. This necessitates the use of a canonical ensemble and introducing a parameter, V_0 , to describe the volume of this locality requirement [72]. With this approach, strangeness enhancement can be reproduced by increasing V_0 and transi-

tioning from the canonical ensemble in small systems to the grand-canonical ensemble in AA collisions, as depicted in Fig. 2.5.

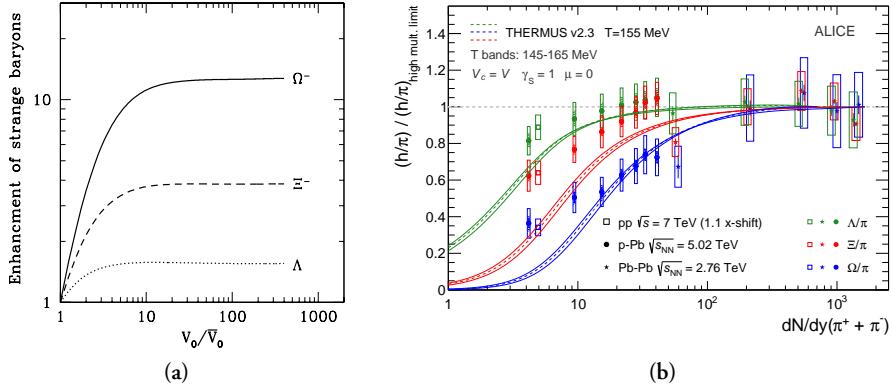


Figure 2.5: (a) Dependence of strange baryon densities on the parameter V_0 characterising the volume where strangeness is locally conserved in models describing strangeness suppression in small systems as canonical suppression. The volume is normalised to a typical AA value of $\bar{V}_0 = 7.4 \text{ fm}^3$. [72]. (b) Ratios of yields of strange baryons to pions in pp, p-Pb, and Pb-Pb collisions as a function of the pion multiplicity normalised to the high-multiplicity limit in 0 – 60% most central Pb-Pb collisions. The results are compared with a statistical model combining the canonical and grand-canonical approach. [69, 71]

2.2.3 Collective flow

The strongly interacting plasma exhibits a collective expansion which can be described by hydrodynamic equations, since the mean free paths of the constituents are much smaller than the system size ($\lambda \ll L$). The non-uniform energy density in the initial state results in varying pressure gradients, which drive this expansion. Since the centre of the plasma has greater pressure than its outside regions, common expansion velocity field develops, which results in the so-called *radial flow*. Similarly, the medium also translates the directionally-dependent anisotropies in the initial state, which stem from the almond-shape geometry of the collision overlap region as well as nucleonic fluctuations, to the final-state. This is the so-called *anisotropic flow*.

Together with hadronic re-scattering, the flows are reflected in the kinematics of the final-state hadrons. When comparing p_T spectra in central AA collisions to those in peripheral or in pp collisions, a broadening as well as a momentum boost can be observed (see Fig. 2.6), caused by the radial expansion as well as the less important thermal motion [73, 74, 75]. The expansion effect depends on the mass of the hadrons, as the amount of additional p_T acquired is proportional to their mass and the collective expansion velocity field, $p \approx m\beta c$. A notable exception to this trend is the ϕ

quarkonium; although comparable with the proton ($m_\phi \approx 1.02 \text{ GeV}/c^2 \sim m_p$), its scattering cross-section is much smaller [76].

The p_T spectra influenced by radial flow can be described by the Blast-Wave parametrisation [77]. In this approach, the radial expansion is accounted for as a common velocity field profile $\beta(r)$ affecting thermal spectra,



where β_s , R , and n are the expansion velocity on the surface of the plasma, its radius, and an extra parameter usually ranging $0.7 - 1.0$ in central collisions [73], respectively. The effects of radial flow can also be reproduced in AA collisions with hydrodynamic models using an equation of state from LQCD and hadronic re-scattering [76], and in pA collisions with the EPOS3 model, which also incorporates hydrodynamic evolution in QGP droplets [78].

Ratios of baryons to mesons, such as p/π or $(\Lambda^0 + \bar{\Lambda}^0)/(2K_s^0)$, as a function of event activity are often used to demonstrate the effect of radial flow, as shown in Fig. 2.6. In these ratios, the modification of p_T spectra results in the following effects on the high-event-activity ratios:

1. The peak in the ratio is shifted to higher p_T by up to $1.5 \text{ GeV}/c$,
2. there is an enhancement of baryons in the intermediate p_T $1.5 < p_T < 6 \text{ GeV}/c$ region,
3. and a corresponding depletion of baryons at low p_T .

Figure 2.7 shows a typical shape of the initial state with its azimuthal anisotropy and the resulting pressure gradients. Anisotropic flow can be quantified by decomposing the azimuthal particle distribution into its Fourier series [79]:

$$\frac{dN}{d\varphi} \propto 1 + 2 \sum_{n=1}^{\infty} v_n e^{in(\varphi - \Psi_n)} , \quad v_n = \langle \cos[n(\phi - \Psi_n)] \rangle , \quad (2.5)$$

where Ψ_n is the symmetry plane of the n -th harmonic and v_n is the Fourier coefficient corresponding to that harmonic, also known as the flow coefficient. In this context, a finite initial state ellipticity ϵ_2 leads to a finite *elliptic flow* v_2 , triangularity ϵ_3 to a *triangular flow* v_3 , and so on [80, 81].

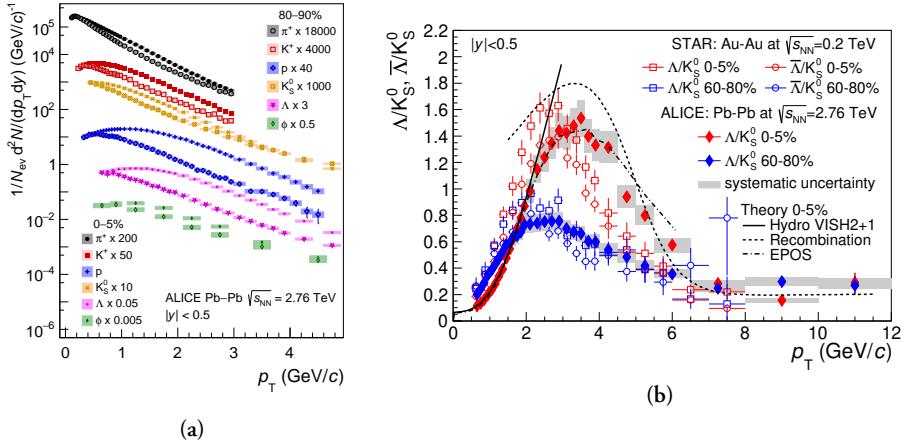


Figure 2.6: (a) Transverse momentum spectra of light-flavour hadrons in central 0–5% and peripheral 80–90% Pb-Pb collisions scaled by arbitrary factors to enhance the visibility. [62, 73, 74, 75] (b) Λ to K_S^0 ratios of transverse momentum spectra in Pb-Pb collisions at the LHC and Au-Au collisions at RHIC for central 0 – 5% events (red) and peripheral 60 – 80% events (blue). [74]

The flow coefficients can be experimentally extracted using various methods, including two-particle azimuthal correlations (as shown in Fig. 2.7), and are typically studied as a function of event multiplicity. It is important to note that these azimuthal correlations between particles due to anisotropic flow are long-range, i.e. present consistently across the entire pseudorapidity range $\Delta\eta$ (the so-called “ridge”) [82], which makes them distinguishable from similarly appearing “non-flow” short-range correlations coming from jet fragmentation and resonance decays.

Moreover, measurements of v_2 in AA collisions for different particle species reveal a mass dependence in the low- p_T region, and a baryon/meson dependence in the intermediate p_T region, with baryons having approximately 1.5 times higher values [83]. This suggests that the flow of hadrons is built up from its deconfined constituents.

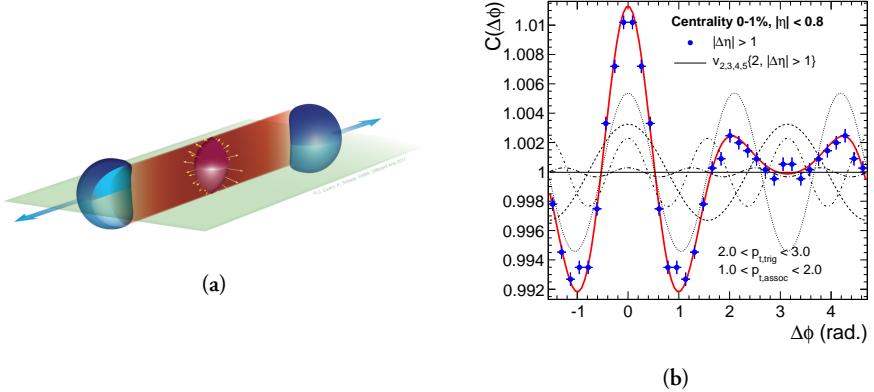


Figure 2.7: (a) Illustration of a collision of ultra-relativistic heavy nuclei and the overlapping region with pressure gradients (yellow). [84] (b) Correlation function of the relative azimuthal angle between a trigger particle and an associated particle, separated by a pseudorapidity gap, measured in central Pb-Pb collisions. The contributions from the elliptical, triangular, quadrupolar, and pentapolar harmonics are shown as different dashed lines. [81]

2.2.4 Jet quenching

In AA collisions, partons produced in hard scattering processes interact with the colour charges in the quark-gluon plasma, resulting in the loss of energy through collisions and gluon bremsstrahlung. This phenomenon is known as jet quenching and modifies or even "quenches" the parton shower [85, 86]. In the factorisation theorem in Eq. 1.13, this corresponds to the medium-modification of the fragmentation functions. Studies of parton energy loss and jet quenching often use the transport coefficient \hat{q} , which describes the average p_T loss of a parton per a mean free path in the plasma and corresponds to the medium opacity

Jet quenching is one of the most important probes into the structure and dynamical properties of the QGP, since the hard partons experience its entire evolution, similarly to the case of heavy quarkonia discussed in Sec. ???. It can also be compared with a large body of sound theoretical calculations [87] and MC simulations [88] based on QCD. Experimentally, jet quenching can manifest as suppression of the jet yield or even its complete disappearance, due to the energy loss of partons in the medium and their re-scattering. [62]

Figure 2.8 displays such an AA event with a large jet imbalance and juxtaposes it with a pp event, where the leading and the recoil jet need to be balanced in azimuth due to conservation laws. Further effects and fields of study of jet quenching include the modification of jet substructures and shapes.

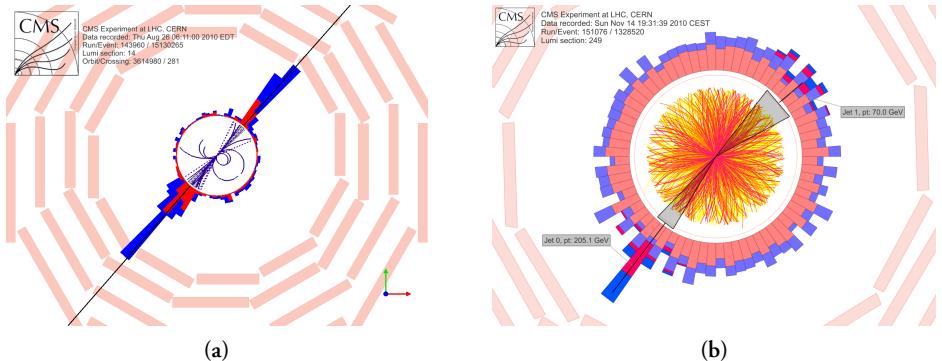


Figure 2.8: Event displays from CMS in azimuthal plane showing a collision of (a) pp and (b) Pb-Pb with a dijet. The red and blue columns correspond to energy deposits in the detector calorimeters. [57, 89]

2.2.5 Cold nuclear matter effects

It should be noted that apart from the QGP, other effects come into play due to the fact that the collision involves two nuclei instead of two protons. These effects are important caveats to bear in mind and include:

1. Nuclear (anti-)shadowing: Reflects the modification in production due to differences in nPDFs and PDFs. [90]
2. Cronin effect: Describes the initial parton energy loss due to scatterings in the nuclear medium and broadens measured p_T spectra. [91]
3. Nuclear absorption: Describes the dissociation of particles due to their interactions with the passing-by nuclear remnants [92]. It is generally negligible at LHC energies.
4. Co-mover absorption: This is the effect of inelastic interactions with the hadron gas. [93]

These effects can be isolated and quantified in pA or very peripheral AA collisions.

2.3 QGP phenomena in small systems

Measurements within the last decade have shown that certain QGP phenomena can also be observed in high-multiplicity events of pp collisions at LHC energies, which

challenges the traditional assumption that QGP is only produced in AA collisions. This has sparked debates about the existence of QGP in pp collisions and, to a lesser degree, about the absence of QGP in AA collisions, despite the extensive experimental evidence.

Furthermore, the observed behavior of these phenomena indicates that the role of event multiplicity N_{ch} may be more significant than the collision system size. This has led to ongoing efforts to establish a consistent and seamless link between the paradigms of pp and AA collisions.

Strangeness and charm enhancement

ALICE measurements on Λ/π , Ξ/π , and Ω/π ratios demonstrate that the production rates of particles containing strange quarks increase faster with multiplicity than those containing only u and d quarks [94]. This also depends on the strangeness content – the effect is the strongest for Ω and vanishes for protons. Furthermore, the evolution to larger systems seems to be continuous with respect to N_{ch} . The measurements can be seen in Fig. 2.9 [62].

To contrast the strangeness measurements with charm, the $J/\psi/\pi$ ratio also shows a clear increase in yield with increasing N_{ch} in pp collisions, as is shown in Fig. 2.9 [95, 96]. However, this comes with an important caveat: high-multiplicity events are biased to have enhanced hard processes, as discussed further in Chapter ???. Moreover, the evolution of this phenomenon is also not continuous with N_{ch} when going from pp collisions at $\sqrt{s} = 13 \text{ TeV}$ to $\sqrt{s_{\text{NN}}} = 5.02 \text{ TeV}$, which can also be explained by the fact that charm quarks are produced solely in hard scattering processes, the rates of which depend on the collision system and center-of-mass energy.

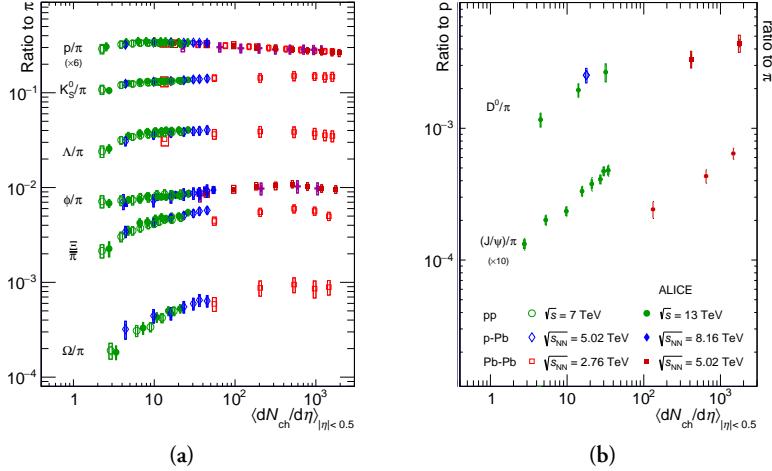


Figure 2.9: Ratios of integrated yields of (a) various light-flavour hadrons [94] and (b) charm mesons [95, 96] to pions as a function of multiplicity in pp, p-Pb, and Pb-Pb collisions. [62]

Anisotropic flow

Azimuthal correlations and anisotropic flow measurements in small collision systems exhibit features similar to those observed in AA collisions, hinting at the presence of collective expansion [97]. However, in small systems, these measurements are particularly challenging due to their large sensitivity to non-flow effects, such as jet fragmentation or resonance decays, which can mimic the features of collective flow.

While models using hydrodynamic-like descriptions seem to be able to describe v_2 results (despite the fact that their assumption $\lambda \ll L$ is not valid) [98], especially at high multiplicities, the interpretation of the results in small systems is still under investigation. The values of elliptic flow v_2 seem to be comparable to those in low-multiplicity Pb-Pb collisions, although the evolution of v_2 across different system sizes does not appear to be smooth. The measurements from CMS displaying a clear ridge in high-multiplicity events [97] and the v_2 results from ALICE [99] can be seen in Fig. 2.10.

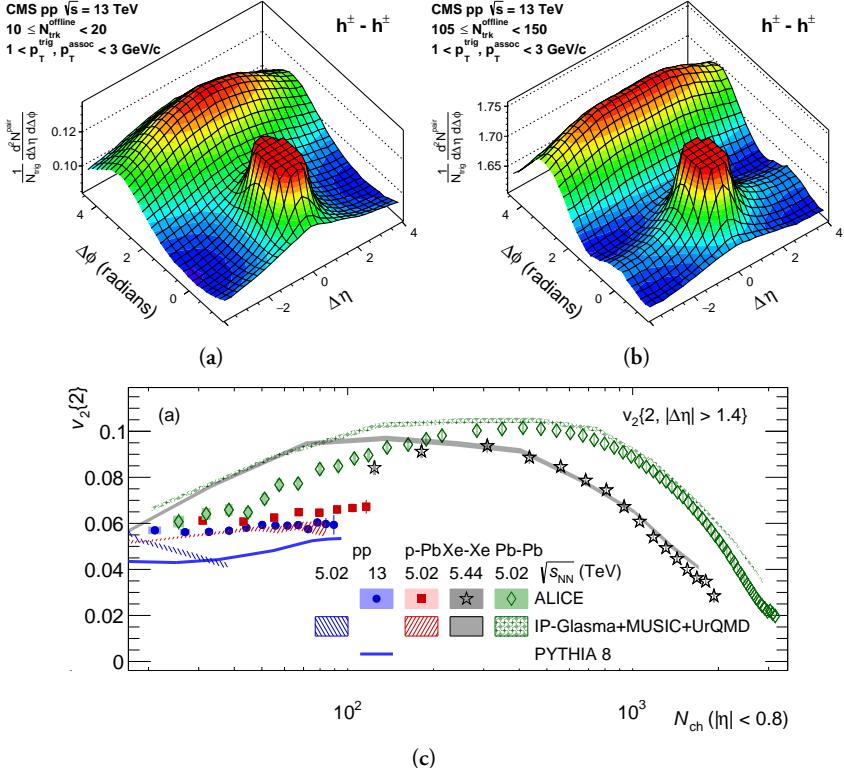


Figure 2.10: (a, b) Two-dimensional two-particle correlation functions of charged hadrons in low (left) and high (right) multiplicity events of pp collisions at $\sqrt{s} = 13 \text{ TeV}$. [97] (c) Elliptic flow measured using two-particle cumulants with a pseudorapidity separation in pp, p-Pb, Xe-Xe, and Pb-Pb collisions as a function of multiplicity. [99]

Radial flow

Measurements of the ratio of Λ to K_S^0 p_T spectra ratio were also studied in pp collisions with differing N_{ch} , see Fig. 2.11 [100]. The boost of a collectively expanding system, as expected in the context of radial flow, should have a greater impact on heavier hadrons, leading to an enhancement of the baryon-to-meson ratio at intermediate p_T . This enhancement is observed in the Λ/K_S^0 ratio, its magnitude increases with increasing N_{ch} and the peak position shifts towards higher values collisions, consistent with the hydrodynamic picture. The increase at intermediate momenta leads to a corresponding depletion at low p_T . High- p_T (as well as integrated) Λ/K_S^0 ratios exhibit essentially no (or minor) multiplicity dependence. This observation also applies to proton-to-pion ratios.

Recent studies have also investigated the charmed baryon-to-meson ratio Λ_c/D^0 , with

similar findings, although measurements with smaller uncertainties are still required. Fig. 2.11 presents the corresponding results.

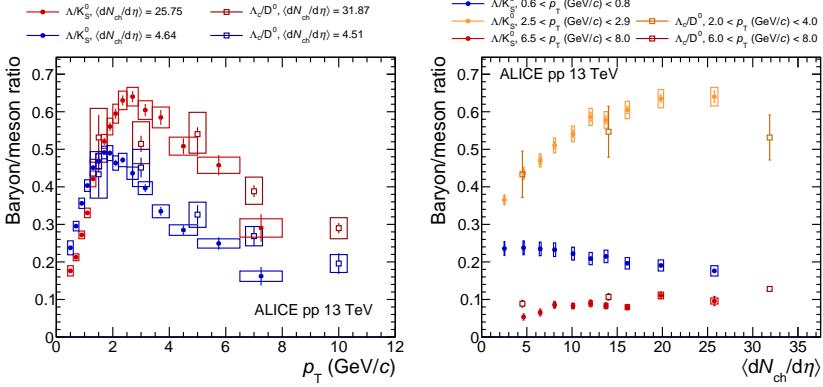


Figure 2.11: Baryon-to-meson ratios shown as the p_T differentials (left) and integrated yields in various p_T ranges as a function of multiplicity (right) for the Λ/K_S^0 and Λ_c/D^0 in pp collisions at $\sqrt{s} = 13$ TeV. [62, 96, 100]

Sequential suppression of Υ states

While defining R_{AA} to compare high-multiplicity and low-multiplicity events is unclear, and measuring yields as a function of N_{ch} is complicated by its biases related to the hardness of primary scatterings, it is worthwhile to investigate the ratio of excited-to-ground states of quarkonia as a function of N_{ch} .

Interestingly, these results [101] exhibit a decrease with increasing N_{ch} , resembling the pattern of sequential suppression due to QGP deconfinement. Even more remarkable, this dependence disappears in low-sphericity, jet-dominated, events (event shape observables such as sphericity are discussed in more detail in Chapter 5.7.2). These findings, reported in Fig. 2.12, suggest that the dependence on N_{ch} is solely influenced by the UE, rather than jets. As event multiplicity grows larger, excited Υ states become relatively less likely to be measured compared to the ground state.

These results indicate the need for a better understanding of Υ hadronization and the role UE may play in it. They also raise the question of whether the ground state is enhanced rather than the excited states being suppressed. Additionally, the effects of the mass differences must also be considered. However, the fact that low-sphericity, jet-dominated events have the same ratios as high-sphericity, UE-dominated events at low N_{ch} argues against these ideas.

An important caveat to note is that hadronic decays (which are dominant) of the

heavy Υ states may result in tens of produced particles [102]. Therefore, even minor discrimination against the excited states could hypothetically be correlated with a substantial but trivial increase in the accompanying N_{ch} . To the author's knowledge, there are currently no available phenomenological descriptions of the observed behavior, which further limits potentially groundbreaking interpretations.

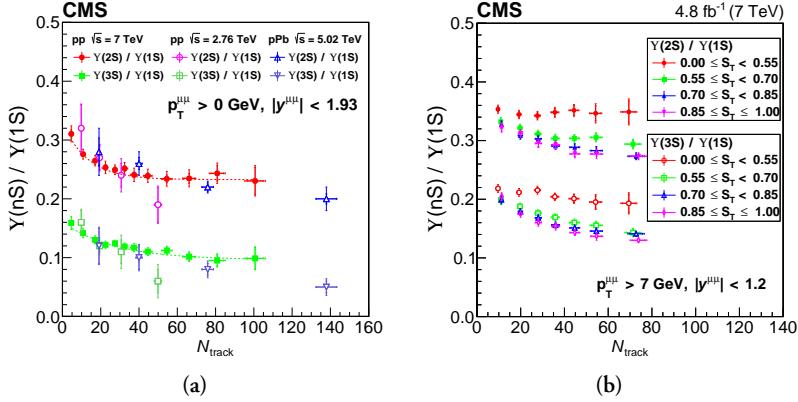


Figure 2.12: The $\Upsilon(2S)/\Upsilon(1S)$ and $\Upsilon(3S)/\Upsilon(1S)$ ratios of measured yields in pp collisions as a function of (a) multiplicity, compared with p-Pb results, and (b) multiplicity and transverse sphericity, with the di-muon transverse momentum $p_T^{\mu\mu} > 7 \text{ GeV}/c$. [101]

Other QGP signatures

If a QGP is formed in small systems with sufficient volumes, the effect of jet quenching should be observed. One can also expect to observe it, due to the fact that in small systems, high- p_T hadrons have been measured to have finite flow v_2 [103], which could indicate that hard partons interact with an expanding medium. Whilst theoretical approaches do not provide unambiguous answers on whether this phenomenon can be observed [104] or not [105], experimental results on jet quenching in both pp and p-Pb collisions are consistent with no observable effect, within uncertainties [106]. These results are mostly based on measuring jet yields as a function of event activity, although such measurements are challenging due to fluctuations and interplays between jet characteristics and event activity.

2.3.1 Role of multiplicity

The observations made above highlight the significance of studying the role of multiplicity N_{ch} . In contrast to AA collisions, high-multiplicity events in pp collisions do not arise from a mere increase in the amount of colliding matter, as the values of

N_{part} and N_{coll} are fixed:

$$N_{\text{part}} = 2, \quad N_{\text{coll}} = 1. \quad (2.6)$$

Additionally, due to the relatively constant initial system volume, high- N_{ch} pp events may exhibit energy densities that exceed the threshold for QGP formation, given that the highest N_{ch} values are similar to those observed in peripheral AA collisions, where QGP formation is observed.

Clearly, the picture is more complex and despite its simplicity as an event activity classifier, N_{ch} poses challenges when it comes to relating data to theory since it cannot be directly linked to the initial state, and multiplicities in different events may originate from entirely different processes.

To address these issues and gain a better understanding of the evolution between low and high multiplicities and the potential for QGP formation, this dissertation focuses on transverse spherocity $S_O^{(p_T=1.0)}$ and underlying event activity R_T measurements. The goal of these studies is to provide a deeper insight into the relevant degrees of freedom involved.

2.4 Phenomenological models

The next parts of this thesis give an overview to phenomenological models and event generators pertinent to the measurements in Chapters 5.7.2 and ???. Other generators, such as Herwig 7 [107] or Sherpa [108], are not discussed here.

2.4.1 Pythia

Pythia is a Monte Carlo event generator used to simulate full events of high-energy particle collisions, based mostly on approximately perturbative QCD, with some important non-perturbative aspects. With more than four decades of development, only a brief overview is given in this thesis, whilst detailed description can be found in Ref. [29, 109]. It has a modular structure to simulate different aspects of the collision process and includes the simulation of the initial kinematics, hard scattering, multiple parton interactions, parton showering, and hadronization, which were all discussed in Chapter ???. The various components of the event simulation, such as the matrix element for the primary scattering, or even the parton evolution, can also be replaced with external alternatives.

Its current and in ALICE most widely used version is Pythia 8, specifically its Monash tune [110], incorporating colour reconnections. Pythia has also included the implementation of Angantyr [111], a new model for the simulation of collisions of nuclei.

A pp collision event, as simulated by Pythia, can be seen in the illustration in Fig. 2.13 and crudely structured as:

1. Relevant parton kinematics are determined based on the PDFs of the beam particles and nature of the event is decided (e.g. Z^0 production). Produced resonances decay.
2. All subsequent partonic activity is simulated. This includes the initial- and final-state parton radiation, MPIs, and handling of the beam particle remnants. Eventually, after the full parton shower evolutions, a full partonic structure and string configuration is given.
3. Hadronisation occurs by the Lund string model fragmentation. This part is completely non-perturbative and fully phenomenological. Unstable particles decay and a full final-state particle collection is obtained. [112]

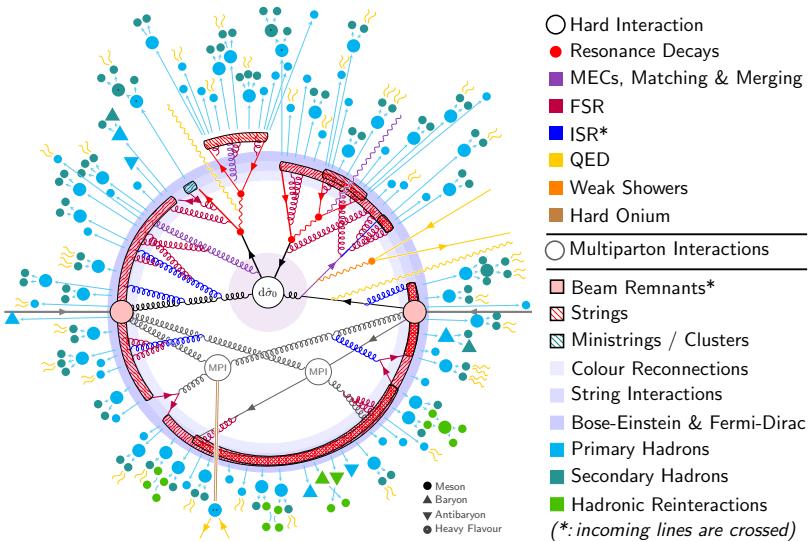


Figure 2.13: Diagram depicting a full simulation of pp collision event with its various components. For full description, see Ref. [29].

2.4.2 String interactions and Ropes

While some effects typically associated with QGP formation, such as radial flow, can be somewhat mimicked in Pythia by colour reconnection [31, 113], as touched upon in Sec. ??, it is not sufficient to describe the strangeness enhancement patterns observed in small systems.

Ideas that overlapping QCD strings in high-density environments interact and form higher-tension ropes date back to modelling AA collisions in 1984 [114] and have been explored further [115], specifically in the framework of the DIPSY model [116]. This “rope hadronisation” approach is also incorporated similarly in Pythia 8 and can be included in its simulations, which in the context of this dissertation will be referred to as the Pythia 8 Ropes tune.

In Pythia, strings are considered overlapping on purely geometrical considerations and utilise a parameter α to quantify the size of strings relatively to the proton radius. Combining two strings follows an algebra based on the SU(3) group, described below, following a more detailed discussion in Ref. [116].

A $q\bar{q}$ string can be viewed as a SU(3) triplet **3**. Stacking another string suggests adding another triplet and forming a multiplet with quantum numbers p , corresponding to the number of coherent triplets **3** (e.g. all red), and q , corresponding to the number of coherent antitriplets **3̄** (e.g. all anti-blue). Using a $\{p, q\}$ notation, the algebra for multiplets is as follows:

$$\{1, 0\} \otimes \{1, 0\} = \{2, 0\} \oplus \{0, 1\} , \quad (2.7)$$

$$\{1, 0\} \otimes \{0, 1\} = \{1, 1\} \oplus \{0, 0\} . \quad (2.8)$$

The first equation, physically, corresponds to merging of two colour strings with colour flows going in the same direction (same $q\bar{q}$ orientation), merging into a rope. When the colours are the same (e.g. both red), the result is a sextet rope **6**, $\{2, 0\}$. In other cases (e.g. red and blue), the result is an anti-triplet rope **3̄**, $\{0, 1\}$ (corresponding to anti-green—green string). This algebra is illustrated in Fig. 2.14.

The second equation describes stacking a triplet **3** with an anti-triplet **3̄** (opposite colour flows and $q\bar{q}$ orientation). This results either in a gluon octet **8**, $\{1, 1\}$, or a singlet **1**, $\{0, 0\}$, with destructive interference and no colour flow.

The tension of the produced rope $\tilde{\kappa}$ is proportional to the quadratic Casimir operator C_2 . When normalising to the tension of a single string κ , e.g. a $\{1, 0\}$ triplet, the

relative increase is given by

$$\frac{\tilde{\kappa}}{\kappa} = \frac{C_2(\{p, q\})}{C_2(\{1, 0\})} = \frac{p^2 + q^2 + pq + 3p + 3q}{4} , \quad (2.9)$$

so in the example of adding two red triplets, the resulting $\{2, 0\}$ rope tension is $\tilde{\kappa} = 5/2$.

During hadronisation, the rope is assumed to break not entirely at once, but rather one string at a time. For the purpose of considering the probability of creating new quarks in (1.15), the effective tension from the $\{p, q\} \rightarrow \{p - 1, q\}$ transition is given by (2.9) and corresponds to

$$\tilde{\kappa}_{\text{eff}} = \frac{2p + q + 2}{4} \kappa . \quad (2.10)$$

This means that in the example of the $\{2, 0\}$ rope breaking, the first quark creation comes with relative effective tension $\tilde{\kappa}_{\text{eff}} = 3\kappa/2$, and the second one with the normal value of $\tilde{\kappa}_{\text{eff}} = \kappa$.

The strangeness production suppression factor in (1.16) then becomes modified as

$$\tilde{\rho} = \rho^{\frac{\kappa}{\tilde{\kappa}_{\text{eff}}}} , \quad (2.11)$$

which makes it evident that overlapping many strings ($\tilde{\kappa}_{\text{eff}} \rightarrow \infty$) results in $\tilde{\rho} \rightarrow 1$ and strange quarks are produced at the same rate as up and down.

Furthermore, ideas for further interactions between strings, such as “string shoving”, wherein overlapping strings may repel each other due to transverse pressure from their excess energy, have also been developed. Such mechanisms produce effects similar to a hydrodynamically expanding medium, e.g. long-range anisotropic flow. [117]

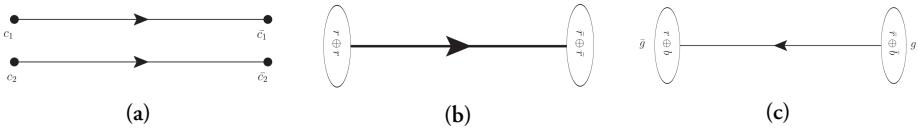


Figure 2.14: (a) Illustration of stacking two colour flow strings, triplets with colours c_1 and c_2 . (b) Rope sextet resulting from adding two coherent triplets, $c_1 = c_2, \{2, 0\}$. (c) Resulting rope anti-triplet coming from adding two incoherent triplets, $c_1 \neq c_2, \{0, 1\}$. Illustrations are by C. Bierlich.

2.4.3 EPOS LHC

EPOS is an event generator built on the Gribov-Regge theory [118], wherein several partons undergo multiple scatterings, each consisting of the hard scattering component as well as initial and final state linear parton emission. Together, they form

a so-called parton ladder and correspond to a “cut” pomeron exchange [119]. The parton ladder represents a (mostly) longitudinally flowing colour field, a “flux tube”, which may hadronise via pair production.

To model the full collision process, EPOS combines a two-component core-corona approach. When the density of flux tubes in a given volume exceeds a parameter ρ_0 , the core is formed. Conversely, flux tubes escaping the volume (usually with higher p_T) make up the corona. The core is assumed to evolve hydrodynamically, corresponding to a QGP droplet, and then hadronises collectively, where smaller core segments form hadrons following a statistical ensemble. Since the relative amount of core- and corona-related particle production can vary continuously, EPOS models can be used to describe pp, pA, and AA collisions using a single paradigm. [120]

EPOS models have been successful particularly at modelling soft-QCD physics and, apart from collider physics, are also widely used in studies of cosmic rays [121]. Throughout this dissertation, an adaptation of EPOS called EPOS LHC [120] is mostly used and shown, which only parametrises the flow dynamics of the core instead of implementing a full hydrodynamic simulation. Diagrams illustrating the partonic structure and core-corona mixing can be seen in Fig. 2.15.

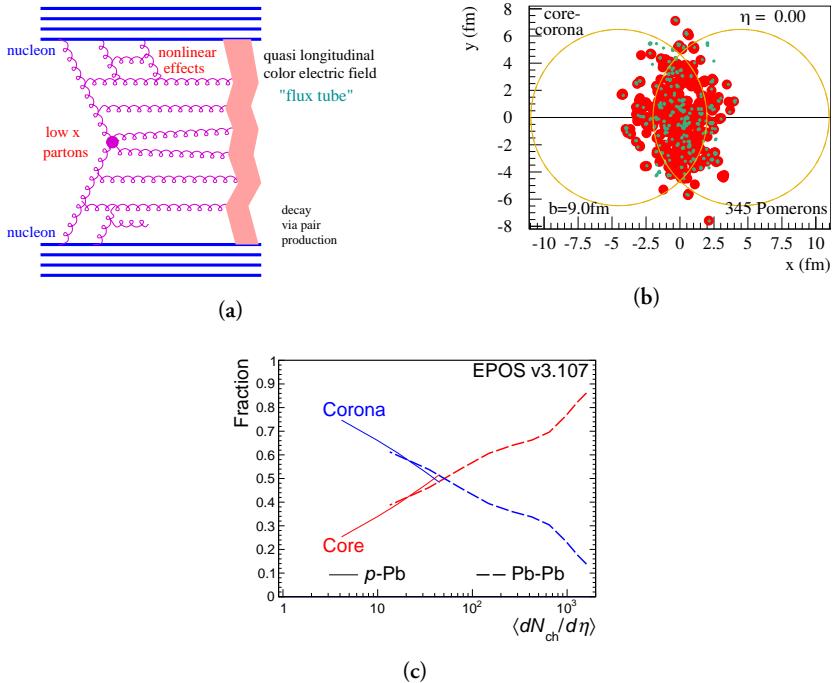


Figure 2.15: (a) Illustration of the partonic structure in the form of a parton ladder. [120] (b) Visualisation of the core (red) and corona (green) components in a peripheral 20 – 40% collision of Pb-Pb with 345 initial multiple scatterings, modelled by EPOS3 [122] (c) Fractions of particle production associated with the core (red) and corona (blue) regions in p-Pb and Pb-Pb collisions, modelled by EPOS3. [123]

Part II

Experimental Setup and Methodology

Chapter 3

ALICE's adventures at the LHC

Chapter 4

Events, Tracks, and Vertices

Part III

Author's measurements

Chapter 5

Reconstruction of neutral strange particles with ALICE

Hadrons K_S^0 and Λ ($\bar{\Lambda}$) are unstable neutral primary particles that usually decay within the volume of the detector through the weak interaction. Their mean lifetimes are $\sim 2.7 \text{ cm}/c$ and $\sim 7.9 \text{ cm}/c$, respectively.[147] Their dominant decay channels, which are also used for their measurement, are:

$$K_S^0 \rightarrow \pi^+ \pi^- \quad (5.1)$$

$$\Lambda \rightarrow p \pi^- \quad (5.2)$$

$$\bar{\Lambda} \rightarrow \bar{p} \pi^+ . \quad (5.3)$$

Because of how these hadrons' decay topologies appear in the detector (an undetectable neutral particle decaying into a V-shaped pair of detectable tracks), they are commonly nicknamed V^0S^1 .

5.1 Analysed datasets

TBA Description of data, collection years, some QA Monte Carlo The Monte Carlo data are simulated using a physics event generator (in this measurement, Pythia 8[29]) and a model describing the propagation of particles through the detector environment (GEANT[148, 149]).

¹Not to be confused with V0A and V0C—the forward calorimeters in ALICE, or V0M—the related multiplicity estimator using the calorimeters' signal.

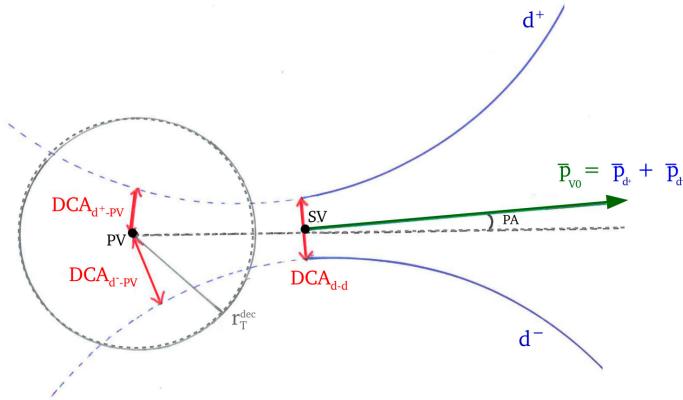


Figure 5.1: Typical topology of V^0 decay. PV stands for primary vertex, SV for secondary vertex. [? , p. 102]

5.2 Identification of V^0 s using ALICE

A centrally developed ALICE algorithm, the ALICE V^0 finder, is used to collect suitable V^0 candidates from pairs of oppositely charged tracks with the relevant topology. This typical topology is illustrated in Fig. ?? Additional selection criteria (“cuts”) are further applied to suppress the background among those candidates. These include:

- cuts on kinematics of the mother and the daughters,
- constraints on the topology of the decay,
- constraints on the reconstruction quality of the daughter tracks,
- cuts on the specific ionisation energy loss of the daughters,
- rejection of contributions from pile-up using “fast detector” information,
- rejection of other competing V^0 candidates based on their invariant mass.

The full list of used cuts is listed in Tab. 5.1.

5.3 Signal extraction

The V^0 signal is separated from the background in distributions of M_{inv} in several p_T intervals using the so-called sideband method. Assuming the signal peaks around $\Delta m_{V0} = M_{\text{inv}} - M_{V0} = 0$ and approximating the background in this region as linear, the subsequent procedure is followed:

Table 5.1: Cuts used in the identification of the K_S^0 , Λ , and $\bar{\Lambda}$ particles.

Cut Variable	Cut Value for K_S^0 (Λ , $\bar{\Lambda}$)
Topology	
V^0 pseudorapidity	$-0.8 < \eta < 0.8$
Transverse momentum	$1.0 < p_T < 25.0 \text{ GeV}/c$
V^0 DCA	$\text{DCA}^{d-d} < 1.0$
Pointing angle	$\cos \text{PA} > 0.97(0.995)$
Decay radius	$0.5 \text{ cm} < R_{xy}$
Daughter Tracks Selection	
DCA of daughters to PV	$\text{DCA}_{xy}^{\text{d-PV}} > 0.06 \text{ cm}$
TPC PID of daughters	$< 5\sigma$
Track pseudorapidity	$-0.8 < \eta < 0.8$
TPC crossed rows	$N_{\text{cr}} > 70$
TPC crossed rows to findable ratio	$N_{\text{cr}}/N_f > 0.8$
Candidate Selection	
Proper lifetime (transverse)	$(R_{xy} \times m_{(\Lambda, \bar{\Lambda})}/p_T < 30 \text{ cm})$
Competing mass	$> 4\sigma$

- the sideband regions are defined. The M_{inv} spectra are fitted in the $-0.03 < M_{\text{inv}} < -0.03 \text{ GeV}/c^2$ interval using a χ^2 -fit with the distribution

$$f = [0] + [1] \cdot M_{\text{inv}} + [2] \cdot \mathcal{N}(\mu, \sigma_1^2) + [3] \cdot \mathcal{N}(\mu, \sigma_2^2), \quad (5.4)$$

where \mathcal{N} is a Gaussian distribution. This is done in all p_T bins and illustrated in Fig. 5.4.

- In each p_T bin, parameter σ is obtained as the RMS of $[2] \cdot \mathcal{N}(\mu, \sigma_1^2) + [3] \cdot \mathcal{N}(\mu, \sigma_2^2)$. To calculate the RMS, the distribution is sampled 10^5 times.
- Variables μ_{V^0} and σ_{V^0} as functions of p_T are interpolated using χ^2 fit and the parametrisations:

$$\mu_{K_S^0}(p_T) = \begin{cases} [0] + [1] \cdot p_T + [2] \cdot p_T^2 & \text{if } p_T < 1.6 \text{ GeV}/c, \\ [3] & \text{if } p_T \geq 1.6 \text{ GeV}/c, \end{cases} \quad (5.5)$$

$$\mu_{\Lambda, \bar{\Lambda}}(p_T) = \begin{cases} [0] + [1] \cdot p_T + [2] \cdot p_T^2 & \text{if } p_T < 1.9 \text{ GeV}/c, \\ [3] + [4] \cdot p_T & \text{if } p_T \geq 1.9 \text{ GeV}/c, \end{cases} \quad (5.6)$$

$$\sigma_{V^0}(p_T) = [0] + [1] \cdot p_T + \frac{[2]}{p_T}. \quad (5.7)$$

The fitted parametrisations can be seen in Fig. 5.3.

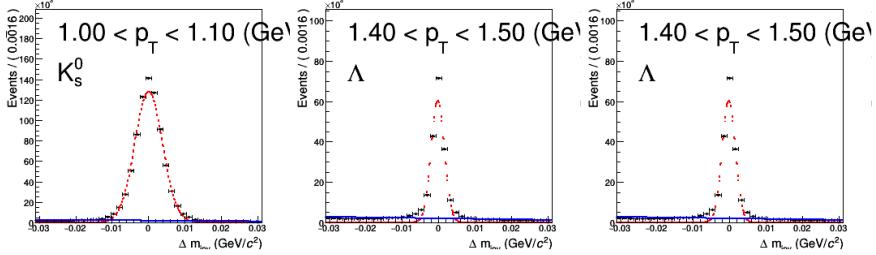


Figure 5.2: Determination of the signal peak mean and width using a fit of the Gaussian distribution for K_S^0 , Λ , and $\bar{\Lambda}$ particles.

4. In each p_T bin, we define the signal region N as $(\mu_{V0} - 6\sigma_{V0}; \mu_{V0} + 6\sigma_{V0})$ and the sidebands A and B as $(\mu_{V0} - 12\sigma_{V0}; \mu_{V0} - 6\sigma_{V0})$ and $(\mu_{V0} + 6\sigma_{V0}; \mu_{V0} + 12\sigma_{V0})$. In these regions, we sum together the entries and acquire N, A, B . The choice of $6\sigma_{V0}$ is rather liberal to avoid biases from incorrect determination of the μ_{V0} or the imperfect description of the signal peak width σ_{V0}
5. Since the background is assumed to be linear, the sum of the two sideband integrals is an accurate estimation of the background in the signal region. Particle yields Y and the corresponding statistical uncertainties σ_Y are calculated as

$$Y = N - A - B \quad (5.8)$$

$$\sigma_Y = \sqrt{N + A + B} , \quad (5.9)$$

due to the fact that the statistical uncertainties in the signal and sideband regions are fully uncorrelated. Illustrations of this step can be seen in Fig. ??.

5.3.1 Validation using simulations

The accuracy of the sideband method is tested with “MC closure”—in MC simulated data, the p_T -spectra acquired blindly from the V^0 candidates are compared with p_T -spectra of identified V^0 . The ratios can be seen in Fig. ?? and show a $\sim 5\%$ effect at high- p_T . This is caused by the fact that in ALICE MC simulations, the V^0 mass peaks have somewhat longer tails than in data and thus the signal can enter the background regions. This has to be taken into account when defining reconstruction efficiency using MC data.

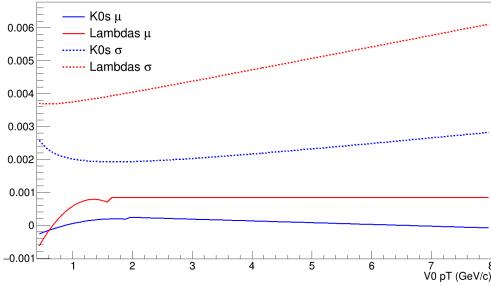


Figure 5.3: Parametrisation of the signal peak mean and width as a function of p_T .

Alternative approach

Originally, methods involving a likelihood fit and an unbinned likelihood fit of two Gaussian distributions as well as other background descriptions were tested. However, although more sophisticated, these methods proved considerably less precise. This is due to the fact that the signal peaks cannot be accurately described by the two Gaussian distributions, particularly in highly populated p_T bins. That said, they are sufficient to determine the σ_{V^0} for above-stated purposes.

Mass resolution of secondary Λ and $\bar{\Lambda}$ particles

Approximately 20% of the Λ yields measured are produced as secondary particles coming from decays of the Ξ baryon in most cases – also called feeddown. Investigations of the simulated data revealed that the invariant mass of these secondaries suffers from a worse resolution (ca. 3 times higher σ). Subsequently, this gives our signal extraction a ca. 75% efficiency for secondaries, and ca. 95% efficiency for inclusive Λ yields at intermediate p_T . This has to be taken into consideration when calculating corrections for the feeddown yields. This effect can be seen in Fig. 5.5.

5.4 Normalisation

The reconstructed K_S^0 , Λ , and $\bar{\Lambda}$ yields $Y(\eta, p_T)$ are normalised according to

$$\frac{d^2N^{\text{raw}}}{dydp_T} = \frac{1}{N_{\text{ev}}} \frac{1}{J} \frac{1}{\Delta\eta} \frac{1}{\Delta p_T} Y(\eta, p_T) \quad , \quad (5.10)$$

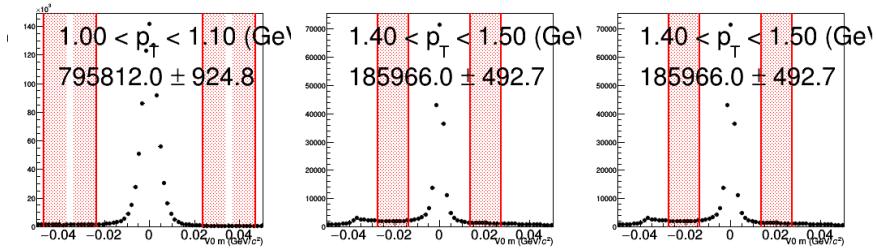


Figure 5.4: Visualisation of the sideband regions, from which the background is estimated, for K_S^0 , Λ , and $\bar{\Lambda}$ particles.

where N_{ev} is the number of selected events, J the Jacobian of the $\eta \rightarrow y$ transformation, and $\Delta\eta$ and Δp_T the widths of the pseudorapidity and transverse momentum intervals, respectively.

TBA Jacobian

TBA Event loss correction

5.5 Corrections to the reconstructed production

To acquire results with scientific relevance, the raw yields of V^0 s observed with ALICE need to be corrected for geometrical acceptance, detector effects, and, in the case of Λ ($\bar{\Lambda}$), also for secondary contribution.

5.5.1 Secondary contribution correction

Only ca. 80% of the measured inclusive Λ and $\bar{\Lambda}$ yields are produced directly in the pp collision or near-instantaneously in non-weak decays of resonances, as primary particles. The remainder is produced secondarily, as products of weak decays of heavier baryons. The dominant, and the only relevant, reactions are:

$$\Xi^- \rightarrow \Lambda \pi^- , \quad (5.11)$$

$$\Xi^0 \rightarrow \Lambda \pi^0 , \quad (5.12)$$

$$\Xi^+ \rightarrow \bar{\Lambda} \pi^+ , \quad (5.13)$$

$$\bar{\Xi}^0 \rightarrow \bar{\Lambda} \pi^0 . \quad (5.14)$$

For the K_S^0 , the secondary production (such as from ϕ mesons) is negligible.

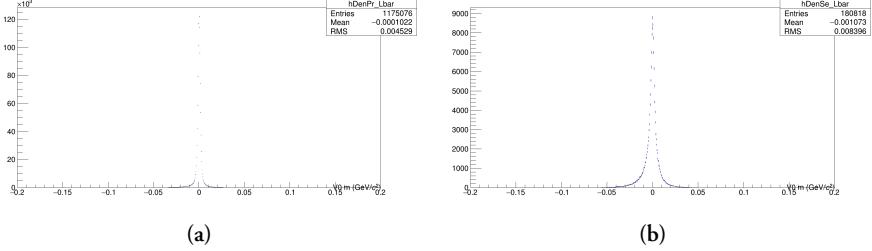


Figure 5.5: TBA.

The primary Λ yields can be estimated using the following equation,

$$\Lambda^{\text{raw}}(p_T^i) = \Lambda_{\text{measured}}^{\text{raw}} - \Lambda_{\text{secondary}}^{\text{raw}} \quad (5.15)$$

$$= \Lambda_{\text{measured}}^{\text{raw}} - \sum_j F_{ij}^{\Lambda} \int_{p_T^j} dN \frac{dN}{dp_T}(\Xi^-) \quad , \quad (5.16)$$

where F_{ij} is the so-called feeddown matrix giving the probabilities of a produced Ξ^- or Ξ^0 particle in a p_T interval j decaying into reconstructed Λ in a p_T interval i , and $\frac{dN}{dp_T}(\Xi^-)$ the measured Ξ^- spectra. This approach assumes that the Ξ^0 decay contribution is identical to Ξ^- and is used because Ξ^0 baryons are challenging to measure. For the $\bar{\Lambda}$, the equation is analogous but uses Ξ^+ .

The feeddown matrix is calculated in ALICE MC simulations of MB events,

$$F_{ij}^{\Lambda} = 2 \cdot \frac{N_{\text{rec.}}(\Lambda)|_{p_T^{\Lambda}=i}^{p_T^{\Xi}=j}}{N_{\text{gen.}}(\Xi)|_{p_T^{\Xi}=j}} \quad , \quad (5.17)$$

where Ξ represent both Ξ^- and Ξ^0 . There is an assumption that the probabilities, and thus, the matrix, do not depend on multiplicity of the event. It is taken into account in systematic uncertainties.

An alternative approach is constructing F_{ij}^{Λ} from charged Ξ solely, and then multiplying $\Lambda_{\text{secondary}}^{\text{raw}}$ by two and was used to determine the systematic uncertainty.

As discussed previously, due to the worse mass resolution of secondary Λ , a M_{inv} cut of $5\sigma_{V0}$ (determined in the sideband definition procedure). Since a large amount of the secondaries enter the background regions, a negative weight -1 has to be applied to achieve the best MC closure validation. Other configurations ($6\sigma_{V0}$ and -1 weight, $4\sigma_{V0}$ and 0 weight) were also tested.

The feeddown matrices F_{ij}^{Λ} , $F_{ij}^{\bar{\Lambda}}$ are displayed in Fig. 5.6.

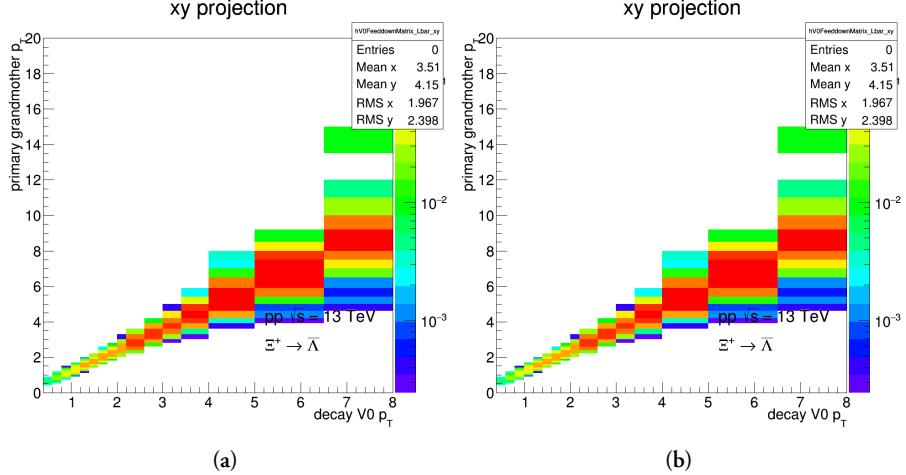


Figure 5.6: Feeddown matrices (a) F_{ij}^Λ and (b) $F_{ij}^{\bar{\Lambda}}$ from Ξ baryons.

Ξ spectra

Fitting. TBA

5.5.2 Reconstruction efficiency

The total reconstruction efficiency, including the acceptance, for V^0 s in our events with ALICE can be determined using the Monte Carlo simulated data. It is calculated as

$$\epsilon(p_T) = \text{acceptance} \times \epsilon_{\text{rec}} \quad (5.18)$$

$$= \frac{\# \text{ associated reconstructed } V^0\text{s}}{\# \text{ generated } V^0\text{s within } |\eta| < 0.8}, \quad (5.19)$$

in events that passed the selection criteria. The association is done by comparing the mother's and daughters' PDG ID as well as the MC generator label. Particles in the numerator have to satisfy all selection cuts. The reconstruction efficiency for K_S^0 , Λ , and $\bar{\Lambda}$ is plotted in Fig. 5.8.

As mentioned before, in ALICE simulations, the M_{inv} resolution worsens with increasing p_T ; in high- p_T bins, the simulated V^0 s are sometimes reconstructed with higher M_{inv} than what is considered realistic. This would lead to a lower efficiency as those V^0 s can fall out of the signal region, and an overestimation of the total measured spectra. For this reason, a $4\sigma_{V^0}$ cut is required for the V^0 s M_{inv} in the numerator.

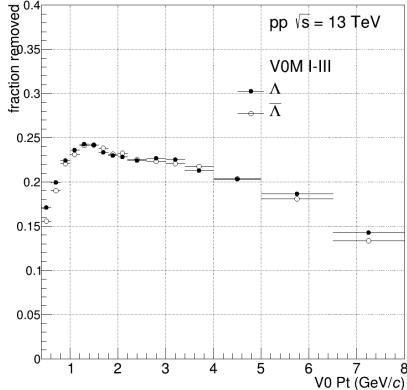


Figure 5.7: TBA.

Alternatively, one could use a cut of $6\sigma_{V^0}$ and applying a negative weight -1 in cases where it is not satisfied.

The reconstruction efficiency is defined in MB events, assuming the reconstruction in pp collisions does not largely depend on multiplicity, geometrical event classification, or event sub-structure. This assumption is taken into account in systematic uncertainties.

5.6 Transverse momentum spectra

Using the corrections on the normalised yields, one acquires the measured transverse momentum spectra, which are comparable with production cross sections and thus theoretical predictions.

$$\frac{d^2N}{dydp_T} = \epsilon(p_T) \times \frac{d^2N_{\text{primary}}^{\text{raw}}}{dydp_T} \quad (5.20)$$

5.6.1 Comparisons with previously published results

The acquired results were tested against previously published measurements of K_S^0 , Λ , and $\bar{\Lambda}$ transverse momentum spectra at the ALICE experiment in MB as well as high-multiplicity (V0M I and V0M III) events in pp collisions at $\sqrt{s} = 13$ TeV.

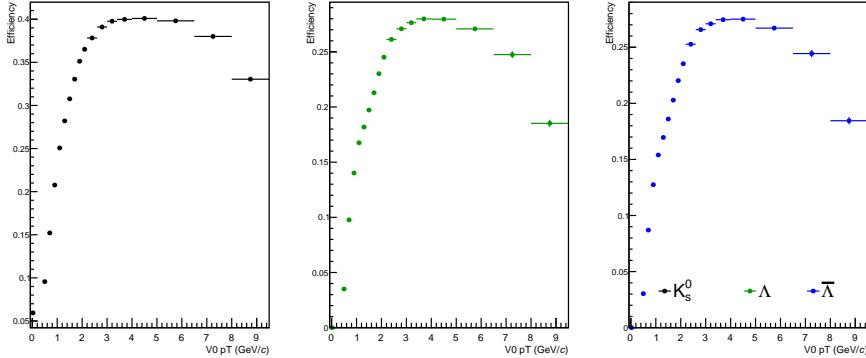


Figure 5.8: TBA.

$$K_S^0$$

The published K_S^0 results were measured in kINT7 events. [] Thus, in order to compare on an equal footing, a trigger efficiency scaling factor $\epsilon_{\text{trig}} = 0.7448$, taken over from [], was applied to this analysis.

The comparison of this analysis to the published results can be seen in Fig. 5.9a. In high-multiplicity events, the spectra are in a good agreement across the entire p_T range (most points lie within $\sim 5\%$ difference). In MB events, there is a difference ($\sim 10\%$) at the lowest p_T values. This is understood as a loss of signal in events with no reconstructed charged tracks and is usually corrected for. Since the correction plays a role only in MB – events which are of little interest to this thesis’ work – it is not taken into account.

$$\Lambda + \bar{\Lambda}$$

The published $\Lambda + \bar{\Lambda}$ results were measured in same events as this analysis, (INEL>0), therefore, ϵ_{trig} was not applied. [] They are compared to this analysis in Fig. 5.9b and show a satisfactory agreement (most points lie within $\sim 5\%$ difference).

5.7 Systematic uncertainties

Experimentally measured values always come with uncertainties – statistical and systematic. Whereas statistical uncertainties are caused by the limited number of measurements and can be decreased by increasing the statistical sample analyzed, system-

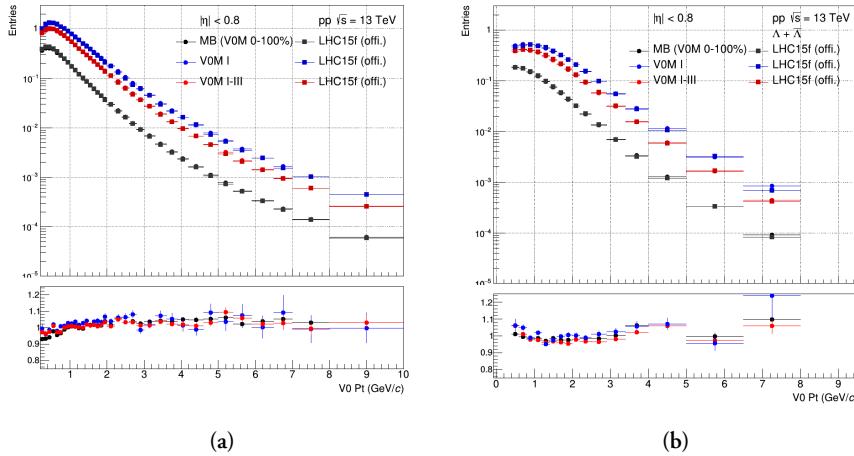


Figure 5.9: Cross-checks of this analysis' p_T spectra of (a) K_S^0 and (b) $\Lambda + \bar{\Lambda}$ in MB, VOM I, and VOM I-III events in pp collisions at against $\sqrt{s} = 13$ TeV results previously published by ALICE.

atic uncertainties represent the imprecision or the bias of the experimental methodology itself. Calculation of statistical uncertainties is given directly from frequentist statistics. Definition of systematic uncertainties, however, is not always straightforward – one cannot simply re-do the measurement with several completely different experimental setups and data analysis techniques. Therefore, a lot of effort needs to go into identifying all possible sources of systematic uncertainties.

In this measurement, the following sources of systematic uncertainty were identified as relevant:

- **Variation of selection criteria**

In determining the reconstruction efficiency, it is assumed that in ALICE MC simulations, all observables used for the identification of V^0 s and for assuring the quality of daughter tracks represent reality. Their inaccurate description, however, results in a bias. This bias is estimated by testing the sensitivity of the final results to varying the selection criteria on these observables.

- **Signal extraction method**

The biases of the sideband background estimation procedure are tested against increasing and reducing the signal and background regions, by varying the number of σ_{V^0} . Variations of 5 and 7 σ_{V^0} were used.

- **Multiplicity dependence of $\epsilon(p_T)$**

Studies of the reconstruction efficiency in pp collisions reveal a small, albeit

significant dependence on the collision final state. A constant uncertainty of $\sim 2\%$ is applied on the spectra to account for this.

- **Feeddown correction**

Three sources of uncertainty on the contribution of secondary particles were identified – variation of the Ξ yields, multiplicity dependence of the feeddown matrix, and an alternative method.

- **Material budget**

This uncertainty reflects that implementing ALICE’s material composition in simulations comes with limitations. Previous studies in ALICE [] which varied parameters of the description of the apparatus showed that this effect corresponds to a constant 4% uncertainty on the measured spectra.

When testing the default method A against an alternative method B , one can implement the deviation of the ratio of their measured values $\Delta = B/A$ from unity as an uncertainty. To ensure that this difference is statistically significant and not just an effect of a limited data sample, the deviation is considered only if it exceeds its own uncertainty, defined as

$$\sigma_\Delta = \frac{\sqrt{|\sigma_B^2 - \sigma_A^2|}}{A} , \quad (5.21)$$

where σ_A and σ_B are the uncertainties of the results from methods A and B , respectively.

5.7.1 Variation of selection criteria

To investigate the differences between description of variables in measured data and ALICE simulations, and determine sensible cut variations λ_i , raw yield loss F was studied. It was measured in MB events and defined as

$$F(\lambda) = 1 - \frac{Y(\lambda)}{Y(\lambda_0)} , \quad (5.22)$$

where $Y(\lambda)$ is the raw yield as a function of the cut value λ and λ_{LOOSEST} the loosest variation (corresponding to the highest yield).

For most observables, the systematic effect can be estimated from alternative methods using λ_{LOOSEST} and $\lambda_{\text{TIGHTEST}}$. To ensure the stability and possible non-linearity, less strict λ_{LOOSE} and λ_{TIGHT} are also tested. If applicable it is reasonable to choose λ_i such that $F(\lambda_i)$ does not exceed approximately 10%.

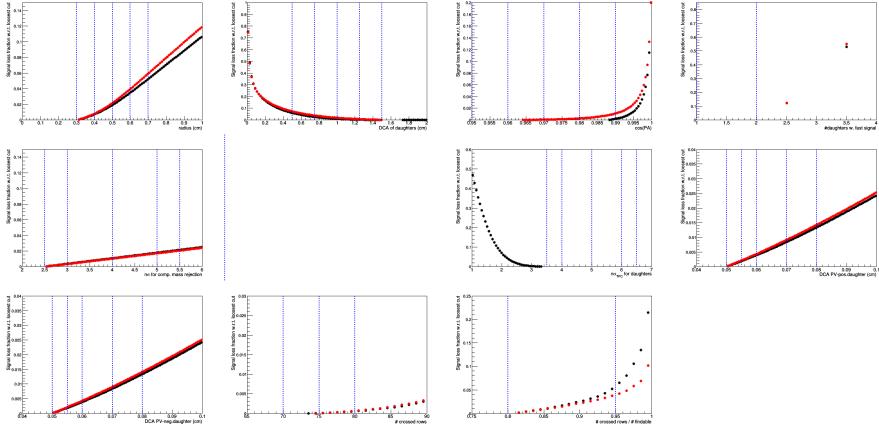


Figure 5.10: TBA

The $F(\lambda)$ for the different selection criteria, and with the chosen λ_i are shown in Fig. 5.10, Fig. 5.11, and Fig. 5.12 for K_S^0 , Λ , and $\bar{\Lambda}$, respectively. The pile-up rejection cut, which requires “fast detector” information for at least one daughter is of a binary nature. So, its variation was tested by requiring a different amount of “fast detector” hits between the two daughters. The selected values of λ_i are summarised in Tab. 5.2.

Table 5.2: Cut variation parameters for the K_S^0 (Λ and $\bar{\Lambda}$).

Quality	loosest	loose	default	tight	tightest
radius	0.3	0.4	0.5	0.6	0.7
DCA between daughters	1.5	1.25	1.0	0.75	0.5
cos PA	0.95 (0.993)	0.96 (0.994)	0.97 (0.995)	0.98 (0.996)	0.99 (0.997)
pile-up removal cut	-	-	1	2	-
comp. mass number of σ	2.5	3.0	4.0	5.0	5.5
lifetime	-	(35.0)	(30.0)	(25.0)	-
TPC PID number of σ	6.5	6.0	5.0	4.0	3.5
DCA to PV of pos. track	0.05	0.055	0.06	0.07	0.08
DCA to PV of neg. track	0.05	0.055	0.06	0.07	0.08
TPC crossed rows	-	-	70	75	80
TPC find. ratio	-	-	0.8	0.95	-

5.7.2 Feeddown correction

As mention before, first, the Ξ spectra, from which the feeddown is calculated, are varied within their reported uncertainties. In both variations, the yields are then extracted using a fit. Second, similarly to $\epsilon(p_T)$, the assumption of no multiplicity

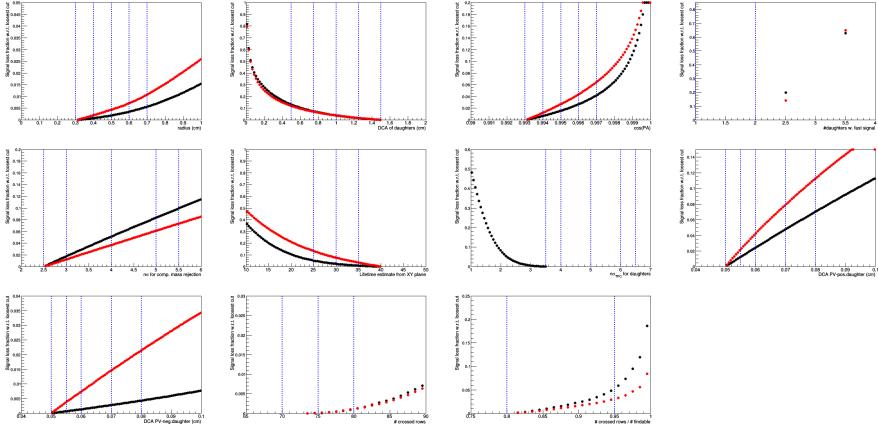


Figure 5.11: TBA

dependence of the feeddown matrix is accompanied by a constant uncertainty of 2% on the secondary yields (corresponding to ca. 0.6% uncertainty on the primary yields).

Lastly, an alternative method of estimating the feeddown just from charged Ξ baryons, and multiplying by a factor of two, was also tested and contributes a systematic uncertainty. It is considered significant and applied when $|\Delta - 1| > \sigma_\Delta$. The difference between the two methods can be seen in Fig. 5.13. It should be noted that whilst the secondary yields suffer from a rather large systematic uncertainty, the effect on the primary spectra is significantly smaller, as the uncertainties enter as $\frac{1-B}{1-A}$ and the secondary yields do not exceed $\sim 30\%$.

In this chapter, measurements of K_S^0 , Λ , and $\bar{\Lambda}$ are reported as a function of transverse spherocity $S_O^{(p_T=1.0)}$, a measure of the event's topology in the transverse xy -plane.

5.8 Transverse spherocity

5.8.1 Motivation for studying event topology

As explained in Sec. 2.3, there is overwhelming evidence that some phenomena associated with QGP, such as collective flow and strangeness enhancement, also arise in pp and p–A collisions at LHC energies and high event multiplicities. This challenges the conventional assumption that the hadron densities and densities of colour fields between partons are too low to interact with each other. Consequently, high-multiplicity pp (and p–A) collisions cannot be treated as a superposition of mostly independent parton-parton (or parton-hadron) scatterings and a more in-depth ap-

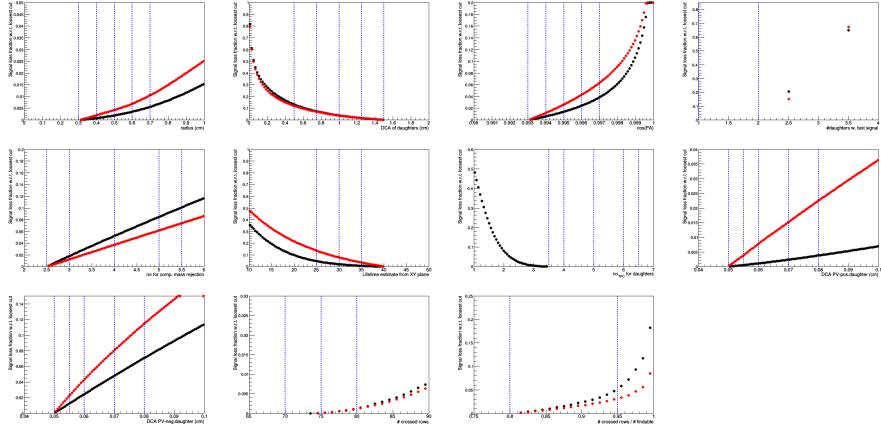


Figure 5.12: TBA

proach is required to fully understand these phenomena.

Event shape observables have been used historically in lepton experiments to study fundamental QCD properties such as the gluon spin [150], and also at Tevatron and the LHC in events with very high p_T ($\gtrsim 100$ GeV/c) jets to further test pQCD predictions [151]. There are various observables, including sphericity S_T , spherocity S_O , thrust T , F-parameter, and Ellis-Karliner angle, most of which are collinear- and infrared-safe and therefore moderately easily calculable[152]. An illustration of two events with different topologies and calculated values of selected observables can be seen in Fig. 5.14.

With the discoveries of QGP phenomena in high-multiplicity collisions of small systems, event shape observables become attractive for different reasons. This is because pQCD (“hard”) processes are responsible for a significant fraction of particle production and are likely to impact the character of QGP phenomena in non-trivial ways. The role of non-perturbative (“soft”) processes is particularly interesting to study as their mechanisms are not fully understood and their interpretation relies on phenomenological models that require clear experimental measurements with high discriminatory power.

Event shape observables allow us to quantify events according to the dominant contributing processes. For instance, collisions with single large p_T transfer scatterings are likely to lead into events with two back-to-back, highly collimated showers, which create a pencil-like shape in the transverse plane. Conversely, collisions with multiple lower p_T transfer partonic interactions will exhibit a high degree of azimuthal isotropy. Therefore, event shape measurements help us gain a deeper understanding of high-multiplicity events and a better control over the magnitudes of the hard

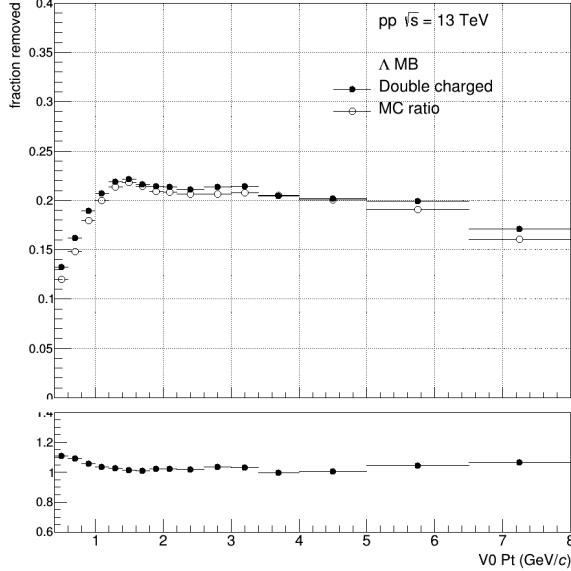


Figure 5.13: TBA

and soft contributions. Ultimately, these measurements may help determine whether QGP formation is necessary in small systems or uncover new physical behaviours.

5.8.2 S_O and $S_O^{(p_T=1.0)}$ as experimental observables

Traditionally, spherocity S_O is defined as:

$$S_O = \frac{\pi^2}{4} \min_{\hat{n}} \left(\frac{\sum_i |p_{T,i} \times \hat{n}|}{\sum_i |p_{T,i}|} \right)^2 , \quad (5.23)$$

where $p_{T,i}$ represents the vector of transverse momentum of a particle i and \hat{n} is the event-dependent unit vector that minimises the sum. The sum runs over all charged particles in the event within the detector acceptance.

Previous ALICE measurements [153] studied characteristics of charged particles in pp collisions and discovered their strong dependence of $\langle p_T \rangle$ on spherocity S_O , which validates the previously discussed motivation. This relationship is shown in Fig. 5.15. Additionally, phenomenological studies of S_O in Pythia 8 further demonstrate its classifying power by finding strong dependence of $\langle n_{\text{MPI}} \rangle$ as well as the mean number of reconstructed jets $\langle n_j \rangle$ on S_O [154, 155]. These results can be seen in Fig. 5.16.

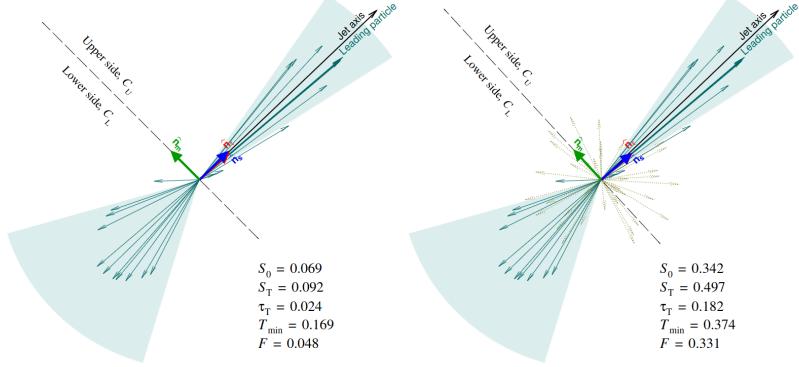


Figure 5.14: Visualisation of two events in the azimuthal plane: (**left**) more “jetty” and (**right**) more isotropic. Values for the transverse spherocity S_0 are given as well as for some other event shape observables: sphericity, thrust, thrust-minor, and F-parameter. The vector minimising the S_0 calculation is denoted as a blue arrow. [151]

This work uses a modified definition of this observable, *unweighted* transverse spherocity $S_O^{(p_T=1.0)}$, defined as follows:

$$S_O^{(p_T=1.0)} = \frac{\pi^2}{4} \min_{\hat{n}} \left(\frac{\sum_i |\hat{p}_{T,i} \times \hat{n}|}{N_{\text{trks}}} \right)^2 , \quad (5.24)$$

where $\hat{p}_{T,i}$ represents the *unit* vector of transverse momentum of a particle i and N_{trks} the number of charged particles entering the sum.

In this thesis, unless stated otherwise, the terms transverse spherocity and spherocity are both used to refer to this unweighted transverse spherocity $S_O^{(p_T=1.0)}$.

Applying the spherocity $S_O^{(p_T=1.0)}$, events in two geometrical limits can be studied:

- $S_O^{(p_T=1.0)} \rightarrow 0$: the “jetty” limit, where a pencil-like topology is selected. These events are dominated by hard pQCD processes. In this limit, with perfectly collimated back-to-back particles, \hat{n} aligns with them. Thus, the sum of vector products in Eq. 5.24 contains only zero values as $\sin 0 = \sin \pi = 0$.
- $S_O^{(p_T=1.0)} \rightarrow 1$: the “isotropic” limit, where a circular topology is selected. Such events are dominated by multiple softer non-perturbative processes². In this limit of $N \rightarrow \infty$ uniformly distributed unit vectors within $(0, 2\pi)$, the

²However, it is important to mention that anisotropic collective flow such as v_2 , a non-perturbative phenomenon, reduces the event isotropy.

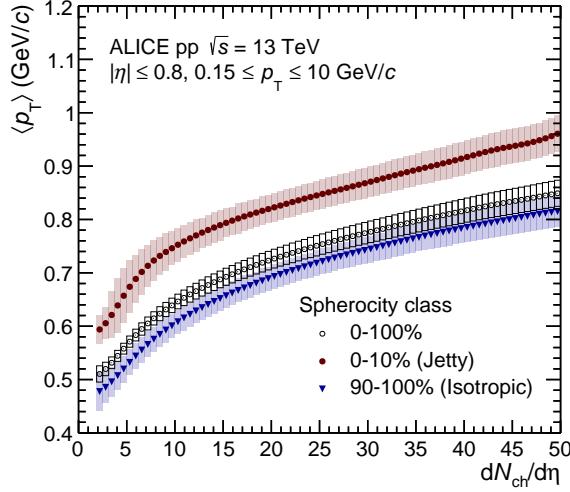


Figure 5.15: Dependence of the transverse momentum spectrum of charged particles as a function of event multiplicity at mid-rapidity. Cases for jetty, isotropic, and spherocity-unbiased events are displayed. Systematic uncertainties are indicated by the boxes and shaded areas. [153]

choice of \hat{n} becomes arbitrary and calculation of the sum in Eq. 5.24 leads to:

$$\frac{1}{N} \sum_{n=1}^N \left| \sin \frac{2\pi n}{N} \right| \approx \frac{1}{N} \int_0^N \left| \sin \frac{2\pi x}{N} \right| dx \quad (5.25)$$

$$= \frac{2}{N} \int_0^{N/2} \sin \frac{2\pi x}{N} dx = \frac{1}{\pi} \int_0^\pi \sin u du \quad (5.26)$$

$$= \frac{1}{\pi} [-\cos x]_0^\pi = \frac{2}{\pi} \quad (5.27)$$

and therefore $S_O^{(p_T=1.0)} = 1$.

Figure 5.17 illustrates how spherocity slowly approaches the circular limit value with increasing N compared to other event shape observables. This property makes spherocity favoured by experimentalists in these measurements, as it provides the highest discrimination power of isotropic events [151].

5.8.3 Relationship between $S_O^{(p_T=1.0)}$ and S_O

In ALICE, only charged particles are considered when calculating spherocities. This introduces biases when measuring charged and neutral species of hadrons. For instance, even topologically identical events with dominant high- p_T leading π^+ and π^0 can yield significantly different values of the traditional p_T -weighted spherocity S_O ,

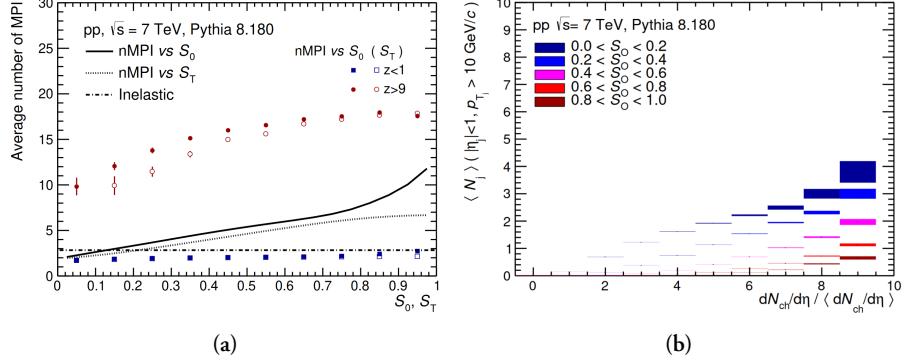


Figure 5.16: (a) Correlation between the mean number of MPIs and transverse spherocity (and sphericity) in pp collisions at $\sqrt{s} = 7$ TeV predicted by Pythia 8, shown as solid and dotted lines. The horizontal line shows the average for minimum bias inelastic events. Events with high (red) and low (blue) self-normalised event multiplicity z are also studied and depicted as datapoints. [154] (b) Average number of reconstructed jets as a function of self-normalised event multiplicity in pp collisions at $\sqrt{s} = 7$ TeV with varying values of spherocity, as predicted by Pythia 8. [155]

despite being comparable in all relevant aspects. In contrast, unweighted spherocity $S_O^{(p_T=1.0)}$ offers a more similar quantification of the two events, as shown in Fig. 5.18. However, it should be noted that this modified definition is only applicable to events with many tracks (i.e., $N_{trks} > 10$).

In addition, while not a large concern in high-multiplicity collisions [153], $S_O^{(p_T=1.0)}$ also offers improved resolution compared to S_O , as the failure to reconstruct a high- p_T track has a smaller impact. Overall, S_O and $S_O^{(p_T=1.0)}$ exhibit similar values and interpretations, with a strong correlation between the two, as illustrated in Fig. 5.19.

5.8.4 Track and event selection

The measurements are carried out on MB events with $INEL > 0$, requiring at least one hit in the V0A or V0C scintillators and one charged particle reconstructed within $|\eta| < 1$. The primary vertex is reconstructed using hits in the SPD and is required to be within 10 cm of the interaction point. The fast read-out time of the SPD allows rejection of out-of-bunch pile-up. In-bunch pile-up is further removed by excluding events with multiple reconstructed vertices. The presented results are based on high-multiplicity events, selected by the classifiers V0M (forward rapidity) and $N_{tracklets}^{|\eta| < 0.8}$ (mid-rapidity), and require a minimum of 10 reconstructed tracks within $|\eta| < 0.8$ and with $p_T > 0.15$ GeV/c.

To ensure a high level of azimuthal acceptance uniformity, which is important for

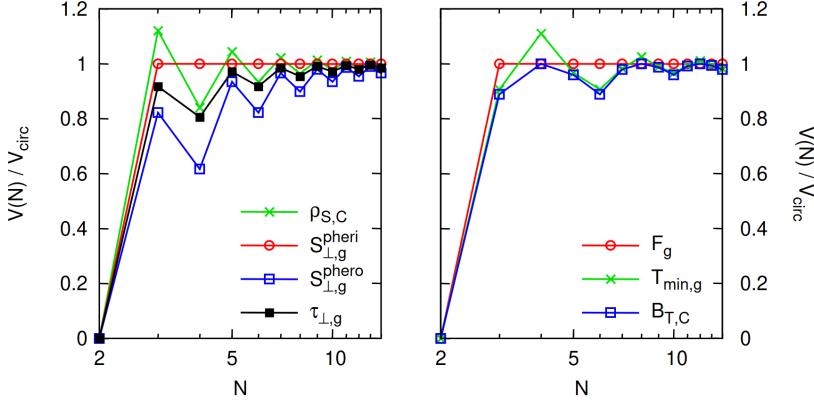


Figure 5.17: Graphs showing how rapidly the different event shape observables approach the circular limit V_{circ} corresponding to $N \rightarrow \infty$, based on the number of perfectly isotropically distributed particles N . Transverse spherocity, here denoted as $S_{\perp,g}^{\text{sphero}}$, exhibits the slowest rise, and never exceeds unity.[152]

event shape measurements, the following, rather loose, selection criteria are employed:

1. The SPD is not used due to its inactive sectors, at the expense of a lower momentum resolution.
2. A track is required to have at least 50 clusters in the TPC and be matched to hits in the ITS to improve tracking precision and further reject pile-up.
3. DCA cuts are applied in both the longitudinal ($|\text{DCA}_z| < 3.2 \text{ cm}$) and transverse ($|\text{DCA}_{xy}| < 2.4 \text{ cm}$) planes to ensure that the reconstructed TPC track points to the primary vertex.

It should be noted that charged decay products of V^0 s with low p_T ($\lesssim 1 \text{ GeV}/c$) and small decay radius may enter and influence $S_O^{(p_T=1.0)}$ determination.

5.8.5 Multiplicity selection and its interplay with $S_O^{(p_T=1.0)}$

Spherocity exhibits a twofold correlation with multiplicity that is not particularly informative. First, the definition of $S_O^{(p_T=1.0)}$ results in higher values for events with more uniformly distributed particles, as shown in Fig. 5.17. Second, in models such as Pythia, high multiplicity is often associated with more MPI, which tend to lead to higher isotropy due to the increased number of emission sources. To gain a more nuanced understanding of the relationship between spherocity and multiplicity, the

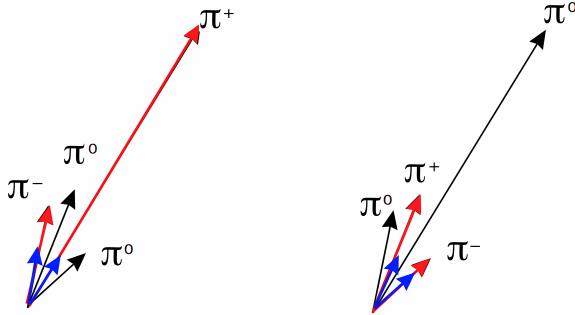


Figure 5.18: Illustration of two topologically identical events with swapped charged and neutral particles. Momentum vectors entering the p_T -weighted sphericity (red) and p_T -unweighted (blue) calculations are displayed, showing a larger difference between the two scenarios for the former than the latter. (Needs to be remade or credited)

effect of $S_O^{(p_T=1.0)}$ on measured particles was analysed in high-multiplicity events determined in two distinct rapidity regions, as described above. Specifically, the top 1% (10%) quantiles are used, denoted as V0M I and $N_{\text{SPD}_{\text{Trkls}}}$ I ($V_{\text{0M I-III}}$ and $N_{\text{SPD}_{\text{Trkls}} \text{ I-III}}$).

Figure 5.20 shows the effect of $S_O^{(p_T=1.0)}$ on the pion yields and $\langle p_T \rangle$. Pions are measured in the high-multiplicity events and in top and bottom 10% and 1% quantiles of $S_O^{(p_T=1.0)}$. The result reveals that when measuring multiplicity in forward rapidity (V0M I), the effect of $S_O^{(p_T=1.0)}$ causes a change of approximately 100% in the yields when going from jetty to isotropic limits, whereas the difference in $\langle p_T \rangle$ is only about 10%. Conversely, when determining the multiplicity in mid-rapidity ($N_{\text{SPD}_{\text{Trkls}}}$ I), the same region where the pion spectra are reconstructed, the change in the yields is only approximately 10% while the change in $\langle p_T \rangle$ is approx. 25%.

For this reason, the following combinations of multiplicity and $S_O^{(p_T=1.0)}$ selections are presented:

1. $N_{\text{SPD}_{\text{Trkls}}}$ I and $S_O^{(p_T=1.0)}$ top and bottom 10% quantiles: This selection emphasises the impact of extreme event topologies on the QCD processes whilst minimising the effect of multiplicity dependence.
2. $N_{\text{SPD}_{\text{Trkls}}}$ I-III and $S_O^{(p_T=1.0)}$ top and bottom 1% quantiles: This selection shows the effect of even more extreme event topologies but with overall less and somewhat varying multiplicity.
3. V0M I and $S_O^{(p_T=1.0)}$ top and bottom 10% quantiles: This selection highlights the effect of extreme event topologies with highly varying mid-rapidity

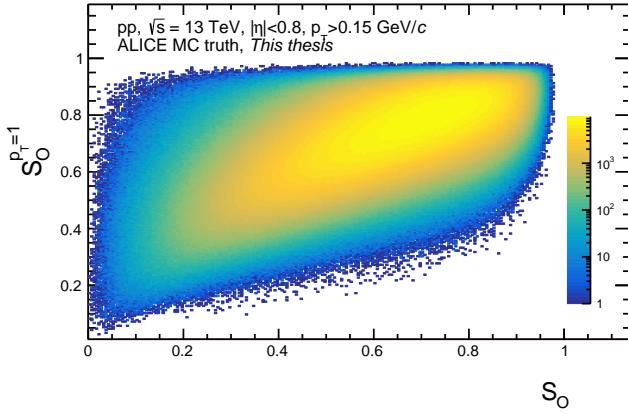


Figure 5.19: Two-dimensional histogram showing the correlation between the traditional transverse spherocity S_O and the unweighted transverse spherocity $S_O^{(p_T=1.0)}$ employed in this measurement, obtained from ALICE MC simulations.

multiplicity. It also allows for a comparison with $N_{\text{SPD}_{\text{Trklts}}}$ I-III and $S_O^{(p_T=1.0)}$ 1% selection, as the mid-rapidity multiplicity and the $\langle p_T \rangle$ variations are more similar.

The mid-rapidity multiplicities in the different high-multiplicity classes are reported in Tab. 5.3. The measured $S_O^{(p_T=1.0)}$ distributions in $N_{\text{SPD}_{\text{Trklts}}}$ I, $N_{\text{SPD}_{\text{Trklts}}}$ I-III, and V0M I-III are shown in Fig. 5.21. They are treated with Bayesian unfolding to account for reconstruction effects. They are also compared with theoretical predictions from Pythia 8 (Monash [110] and Ropes [116] tunes), EPOS LHC [120], and Herwig 7 [107]. Table 5.4 provides the $S_O^{(p_T=1.0)}$ cut values associated with the quantile selections in data.

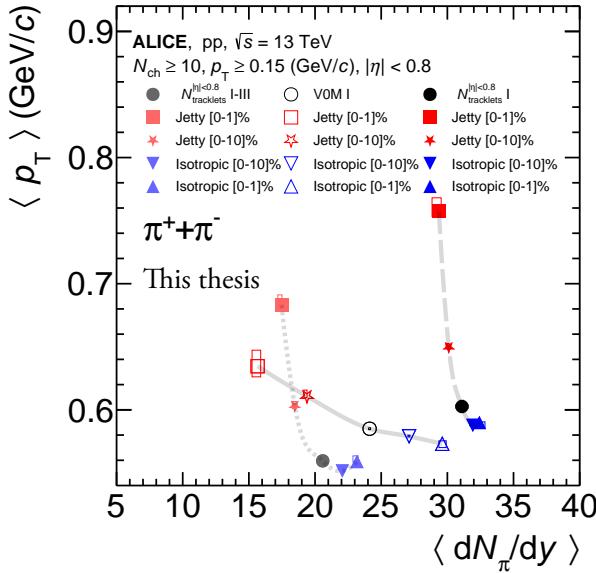


Figure 5.20: Relationship between the mean transverse momentum of charged pions and their yields in events with different $S_0^{(p_T=1.0)}$ values and event multiplicity determined at mid-rapidity (here denoted $N_{\text{tracklets}}^{|\eta|<0.8}$) and forward rapidity (VOM). Events in the same multiplicity classes are connected with lines for clarity's sake.

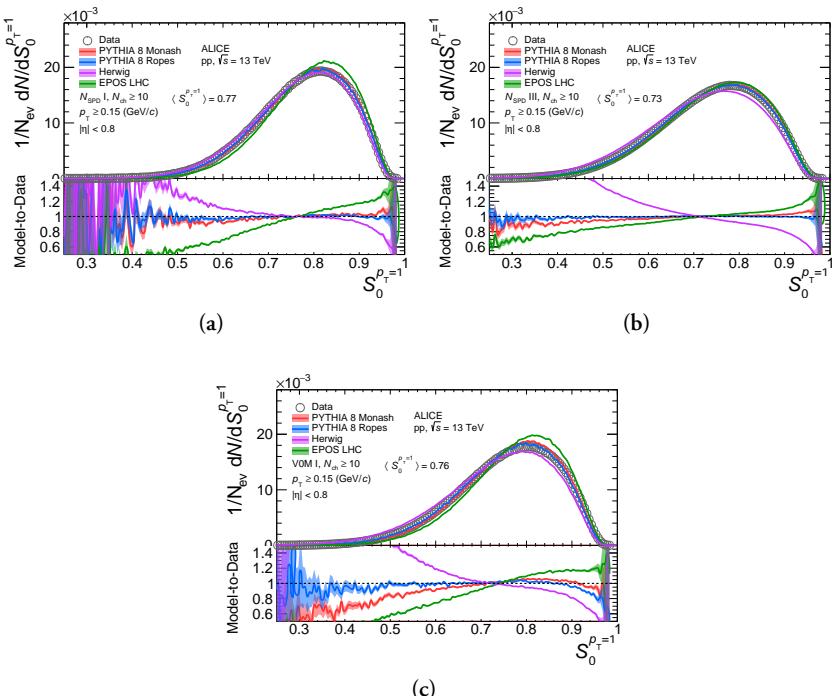


Figure 5.21: The measured and fully corrected $S_0^{(p_T=1.0)}$ distributions for both (a) $N_{\text{SPD Tracklets}}^{|\eta|<0.8}$: 0-1%, (b) 0-10% and (c) VOM I: 0-1%. The curves represent different model prediction, where the shaded area represents the statistical uncertainty of the models.

Table 5.3: Mean values of charged particle multiplicity at mid-rapidity in different high-multiplicity event classes.

Event class	$N_{\text{SPD}_{\text{Trkts}}} \text{ I}$	$N_{\text{SPD}_{\text{Trkts}}} \text{ I-III}$	V0M I
$\langle dN_{\text{ch}}/d\eta \rangle$	33.01 ± 0.55	21.57 ± 0.32	26.02 ± 0.35

Table 5.4: Cut values of the different quantiles of the uncorrected $S_O^{(p_T=1.0)}$ distribution used for the event selections in this analysis.

Event class	$N_{\text{SPD}_{\text{Trkts}}} \text{ I}$	$N_{\text{SPD}_{\text{Trkts}}} \text{ I-III}$	V0M I
Jetty			
$S_O^{(p_T=1.0)} \text{ 0-1\%}$	< 0.487	< 0.408	< 0.433
$S_O^{(p_T=1.0)} \text{ 0-10\%}$	< 0.624	< 0.561	< 0.589
Isotropic			
$S_O^{(p_T=1.0)} \text{ 90-100\%}$	> 0.892	> 0.871	> 0.882
$S_O^{(p_T=1.0)} \text{ 99-100\%}$	> 0.942	> 0.930	> 0.936

5.8.6 Comparison of V^0 production with MC generators

Further on in this chapter, the results of K_S^0 , Λ , and $\bar{\Lambda}$ as a function of $S_O^{(p_T=1.0)}$ are presented and compared with predictions from phenomenological models Pythia 8 [29], EPOS LHC [120], and Herwig 7 [107] obtained from MC simulations. To mitigate the effect of reconstruction on the experimental results and make the comparison with these predictions as comparable as possible, the following strategies were employed based on findings using the ALICE MC simulations:

- The results were compared using the same quantiles of the $S_O^{(p_T=1.0)}$ distributions in both the MC and the data, instead of relying on the experimental $S_O^{(p_T=1.0)}$ ranges determined by specific cut values. This approach reduced the effects of $S_O^{(p_T=1.0)}$ resolution.
- In the MC simulations, the $S_O^{(p_T=1.0)}$ calculations included neutral particles K_S^0 , Λ , and $\bar{\Lambda}$, despite their neutral charge. This helped minimize differences between the true and reconstructed/corrected MC results, possible due to the potential contribution of charged daughters to the $S_O^{(p_T=1.0)}$ calculation.

Any discrepancies that still persisted between the true and reconstructed/corrected transverse momentum spectra were accounted for as systematic uncertainties.

5.9 Systematic uncertainties

The systematic uncertainties associated with the p_T spectra of K_S^0 , Λ , and $\bar{\Lambda}$ were evaluated separately for $N_{\text{SPD}_{\text{Trkts}}}$ I and V0M I events with no $S_O^{(p_T=1.0)}$ selection, as well as for the top and bottom 10% isotropic and jetty quantiles, using the methodology described in Section ???. The relative systematic uncertainties obtained from these configurations were also applied to the $N_{\text{SPD}_{\text{Trkts}}}$ I-III and V0M I-III event classes with different jetty/isotropic quantiles.

Figures 5.22, 5.23, and 5.24 illustrate the maximal deviations of variations resulting from alternative cut values, extraction parameters, or feeddown methods in V0M I, $S_O^{(p_T=1.0)}$ -unbiased events for K_S^0 , Λ , and $\bar{\Lambda}$, respectively.

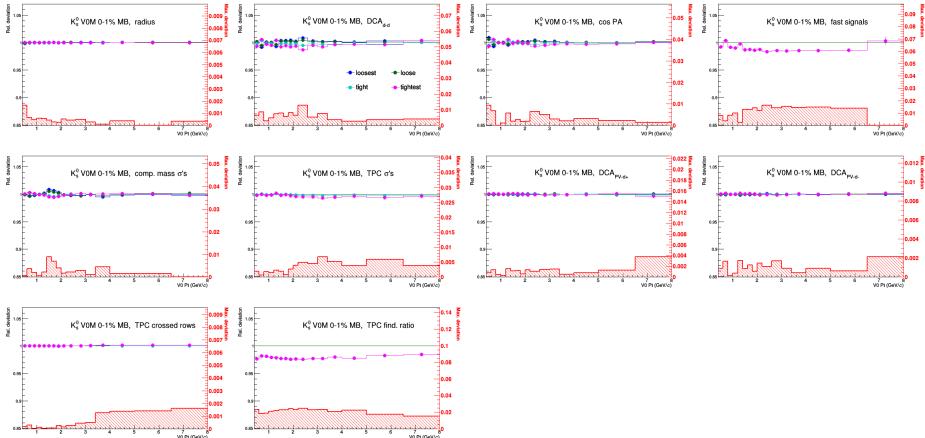


Figure 5.22: NEEDS TO BE REDRAWN IN CONSISTENT STYLE WITH RT CHAPTER:Deviations of the corrected spectra w.r.t. the different cut variations used for the K_S^0 . The maximum deviation is added in quadrature to the total, if it's larger than σ_{RB} (depicted as errorbars) from unity.

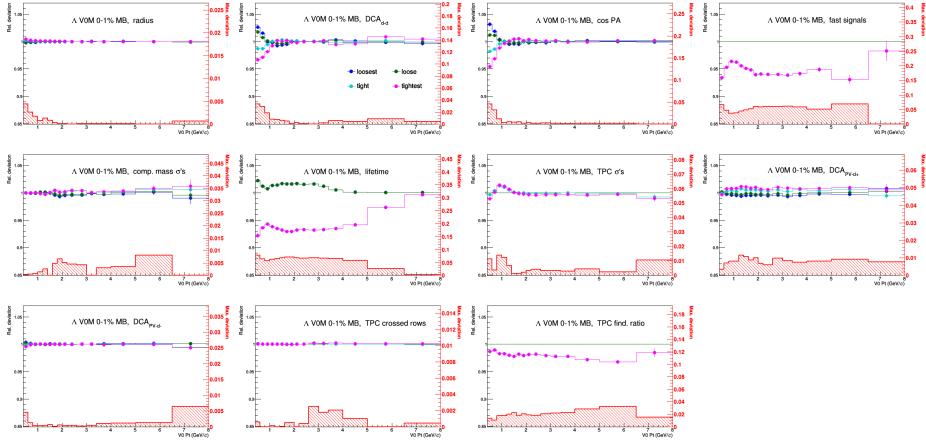


Figure 5.23: NEEDS TO BE REDRAWN IN CONSISTENT STYLE WITH RT CHAPTER:Deviations of the corrected spectra w.r.t. the different cut variations used for the K_S^0 . The maximum deviation is added in quadrature to the total, if it's larger than σ_{RB} (depicted as errorbars) from unity.

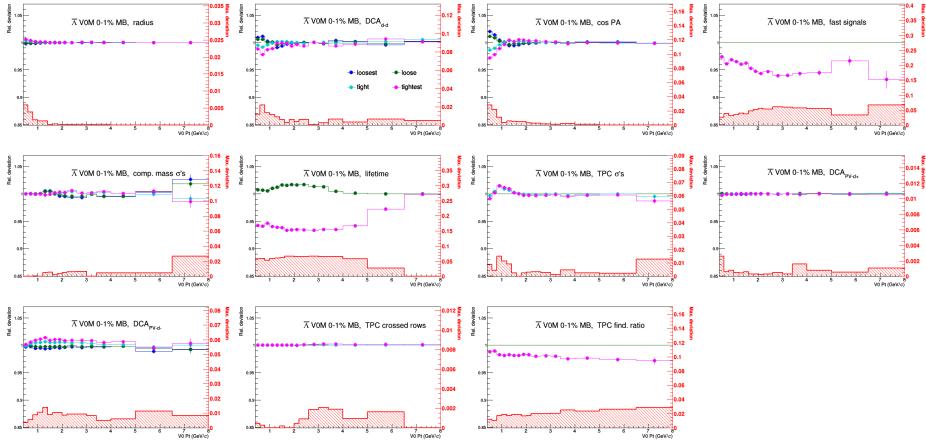


Figure 5.24: NEEDS TO BE REDRAWN IN CONSISTENT STYLE WITH RT CHAPTER:Deviations of the corrected spectra w.r.t. the different cut variations used for the K_S^0 . The maximum deviation is added in quadrature to the total, if it's larger than σ_{RB} (depicted as errorbars) from unity.

5.9.1 Experimental bias

To estimate the experimental bias of the $S_O^{(p_T=1.0)}$ selection, Monte Carlo closure tests were used, studying the ratios of true p_T spectra in jetty/isotropic quantiles of the true $S_O^{(p_T=1.0)}$ distribution to the measured and corrected p_T spectra in quantiles of the measured $S_O^{(p_T=1.0)}$ distribution.

Due to the loose DCA cuts used in the $S_O^{(p_T=1.0)}$ determination, it is expected that the V^0 daughters may enter its calculation. Thus, to make the predictions from simulations more comparable to the data, the K_S^0 , Λ , and $\bar{\Lambda}$ particles were included in the true $S_O^{(p_T=1.0)}$ calculation. Although insufficient below $p_T < 1 \text{ GeV}/c$, this works well and corresponds to 1% and 4% discrepancies for isotropic and jetty events, respectively [156].

Alternatively, as discussed further in this dissertation in the R_T measurements in Chapter ?? but not employed here, this effect could be accounted for experimentally, by making the two track sets (spherocity and V^0 daughters) explicitly disjunct by enforcing a DCA cut.

5.9.2 Correlation of uncertainties with $S_O^{(p_T=1.0)}$

Correlations of several systematic uncertainties with respect to the $S_O^{(p_T=1.0)}$ selection are expected. Since the ratios of jetty/isotropic results to $S_O^{(p_T=1.0)}$ -unbiased ones provide important insights, it is necessary to account for these correlations in order to not overestimate the uncertainties, following the methodology described in Fig. ??.

Moreover, the systematic uncertainty associated with the material budget is treated as fully correlated, while assuming the reconstruction efficiency is independent of multiplicity leads to an uncorrelated uncertainty. This assumption would lead to a factor of $\sqrt{2}$ in the ratios of jetty/isotropic to $S_O^{(p_T=1.0)}$ -unbiased for this uncertainty. However, in ALICE, this assumption is generally considered too conservative [146], and thus this factor is dropped. The same approach is used for the uncertainty associated with the multiplicity independence of the feeddown matrix. Detailed results can be found in Appendix ??.

5.9.3 Summary

The total systematic uncertainties are reported in Tab. 5.5 and visualised in Fig. 5.25. The dominant contributions are, in no specific order, selection cuts, experimental bias, signal extraction, and the enforcement of signals from fast detectors to prevent track pile-up.

Table 5.5: The most relevant systematic uncertainties for the long-lived particles K_S^0 and $\Lambda(\bar{\Lambda})$ as a function of $S_O^{(p_T=1.0)}$. “HM” in this table represents the $S_O^{(p_T=1.0)}$ -unbiased spectra. Uncertainties are p_T -dependent, and the ranges listed represents the minimum-maximum values presented in the final spectra (see text for details).

Topology:	Jetty	Iso	HM	Jetty/HM	Iso/HM
K_S^0					
Selection cuts	3%	3–4%	3–4%	Negl.	1%
Track pile-up	1%	1–3%	1%	0–2%	0–2%
Signal extraction	1–3%	1–3%	1–3%	Negl.	Negl.
Efficiency	2%	2%	2%	2%	2%
Material budget	4%	4%	4%	—	—
Experimental bias	4%	1%	—	4%	1%
Total uncertainty	7%	6–7%	5–6%	5%	2–3%
$\Lambda(\bar{\Lambda})$					
Selection cuts	1–5%	2–6%	4–5%	0–1%	0–3%
Track pile-up	4–5%	5%	3–5%	0–1.5%	0–1%
Signal extraction	2–6%	2–6%	2–6%	0–2%	0–1%
Feed-down correction	1.0–1.5%	1.0–1.5%	1.0–1.5%	Negl.	Negl.
Efficiency	2%	2%	2%	2%	2%
Material budget	4%	4%	4%	—	—
Experimental bias	4%	1%	—	4%	1%
Total uncertainty	8–10%	8–9%	7–9%	5%	3–4%

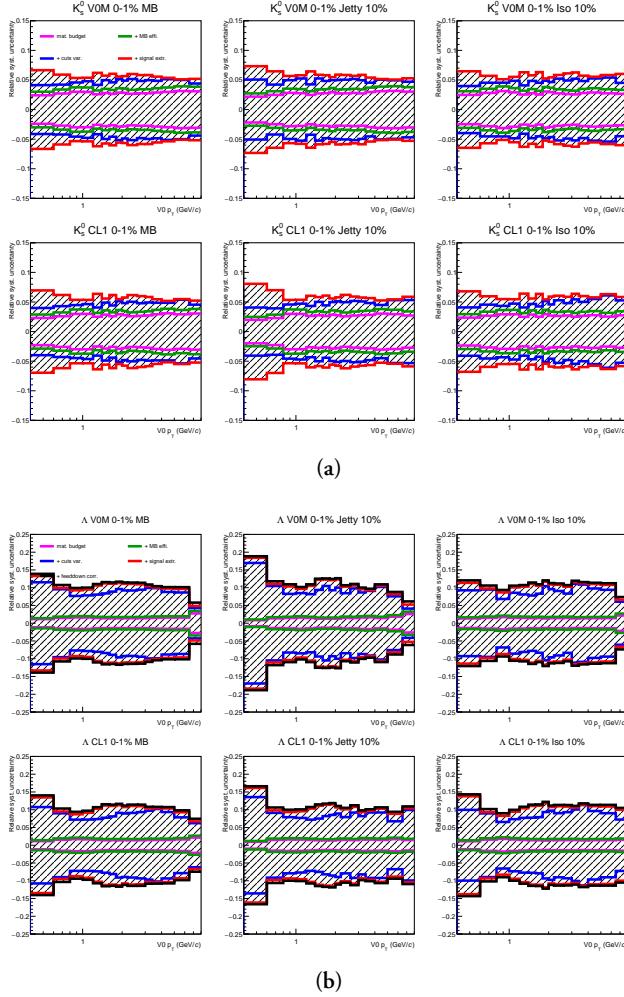


Figure 5.25: NEEDS TO BE REDRAWN IN CONSISTENT STYLE WITH RT CHAPTER: Total relative systematic uncertainty and individual contributions for the K_S^0 , Λ , and $\bar{\Lambda}$.

5.10 Transverse momentum spectra vs. $S_O^{(p_T=1.0)}$

The corrected spectra in V0M and $N_{\text{SPD}_{\text{Trkts}}}$ high-multiplicity events and the dependence on spherocity for the K_S^0 and $\Lambda + \bar{\Lambda}$ can be seen in Fig. 5.26 and Fig. 5.27, respectively. The trends observed in the spectra are consistent between K_S^0 and Λ and indicate a significant hardening (softening) in the low (high) spherocity selection, relative to the inclusive high-multiplicity event class. These trends are qualitatively generally well captured by all included model predictions, favouring Pythia 8 Ropes. Particularly, EPOS LHC overestimates the yields and Pythia 8 Monash fails in de-

scribing the Λ p_T spectra.

The separation between jetty and isotropic events is more pronounced in the N_{SPDTrkts} high-multiplicity events, and shows significant difference in the spectra slopes rather than just an offset, which is somewhat the case for the V0M events.

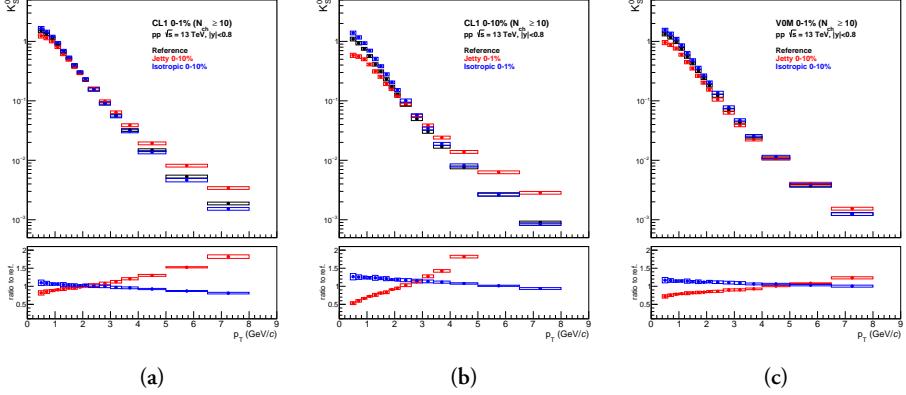


Figure 5.26: Transverse momentum spectra of K_S^0 in high-multiplicity jetty (red) and isotropic (blue) events of pp collisions at $\sqrt{s} = 13$ TeV measured in the (a) $N_{\text{SPDTrkts}} \geq 10$, (b) $N_{\text{SPDTrkts}} \geq 10$ (l-I-III), and (c) V0M event classes. The bottom panels display ratios to the $S_0^{(p_T=1.0)}$ -unbiased spectra. Statistical and systematic uncertainties are indicated with error bars and boxes, respectively. (Needs to be redrawn to also show MC)

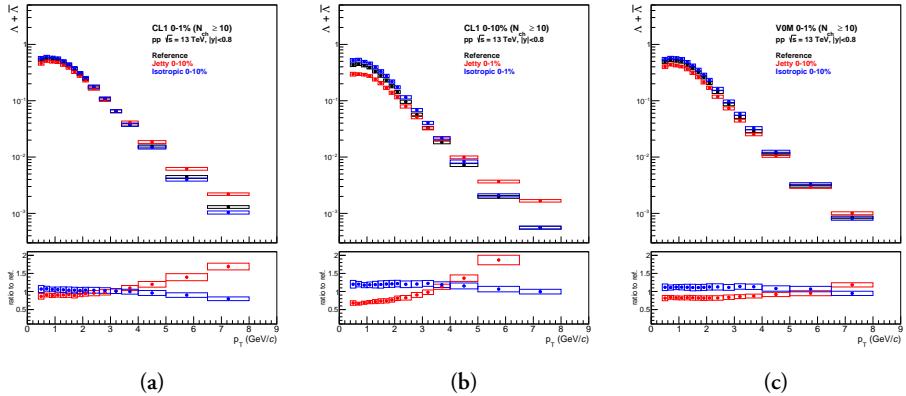


Figure 5.27: Transverse momentum spectra of $\Lambda + \bar{\Lambda}$ in high-multiplicity jetty (red) and isotropic (blue) events of pp collisions at $\sqrt{s} = 13$ TeV measured in the (a) $N_{\text{SPDTrkts}} \geq 10$, (b) $N_{\text{SPDTrkts}} \geq 10$ (l-I-III), and (c) V0M event classes. The bottom panels display ratios to the $S_0^{(p_T=1.0)}$ -unbiased spectra. Statistical and systematic uncertainties are indicated with error bars and boxes, respectively. (Needs to be redrawn to also show MC and be cut off from 1 GeV)

5.10.1 Ratios of neutral kaons to charged kaons

To verify the robustness of $S_O^{(p_T=1.0)}$ as an event observable, the p_T spectra of neutral and charged kaons are compared in the $N_{\text{SPD}_{\text{Trkts}}}$ I and V0M I classes. The ratios exhibit no dependence on $S_O^{(p_T=1.0)}$ and are consistent with unity, according to expectations. The slight depletion of K_S^0 in the $N_{\text{SPD}_{\text{Trkts}}}$ I class is interpreted as the multiplicity selection bias due to requiring a large number of charged tracks at mid-rapidity, and is also reproducible by simulations. The ratios are shown in Fig. 5.28.

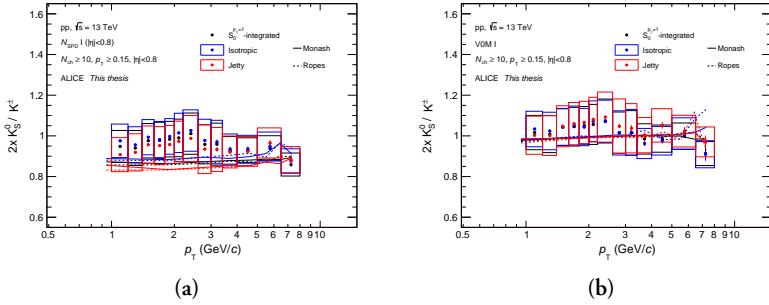


Figure 5.28: Ratios of transverse momentum spectra of K_S^0 to K^\pm in (a) $N_{\text{SPD}_{\text{Trkts}}}$ I and (b) V0M I high-multiplicity events, also showing the effect of spherocity selection. The results are compared with MC predictions of Pythia 8 tunes.

5.11 Mean transverse momenta and integrated yields

The $\langle p_T \rangle$ and particle yields (dN_{ch}/dy) in different $S_O^{(p_T=1.0)}$ bins for the $N_{\text{SPD}_{\text{Trkts}}}$ I class are reported in Fig. 5.29. The measured values of $\langle p_T \rangle$ quantify the observations in the spectra: there is a significant p_T -hardening in jet-like events, consistently seen in both K_S^0 and Λ . Furthermore, the $\langle p_T \rangle$ of the inclusive high-multiplicity events is consistent with that of the isotropic subsample. This result suggests that the average high-multiplicity events and $S_O^{(p_T=1.0)}$ -selected isotropic events are dominated by similar underlying physics processes.

Furthermore, the $S_O^{(p_T=1.0)}$ -integrated event class is not simply the arithmetic average of the jetty and isotropic subsamples. This implies that jetty events are rare outliers of a much more homogeneous group of high-multiplicity events, which will be further focused on in the next sections.

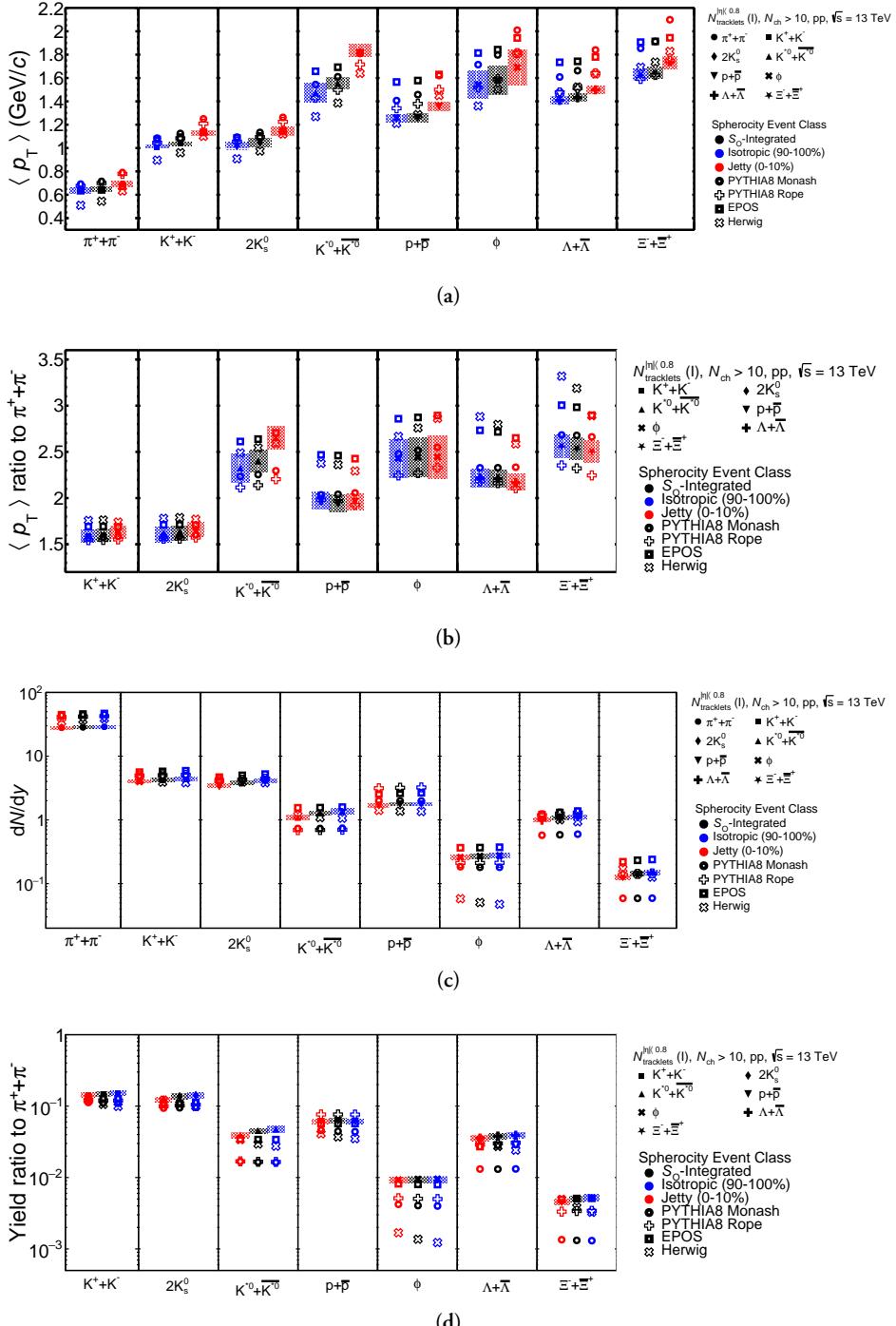


Figure 5.29: (a) Mean transverse momentum and (c) integrated yields of $\pi^+ + \pi^-$, K_S^0 , and $\Lambda + \bar{\Lambda}$ in jetty, isotropic, and $S_O^{(pT=1.0)}$ -unbiased high-multiplicity events determined at mid-rapidity. (b,d) Their ratios of K_S^0 and $\Lambda + \bar{\Lambda}$ to pions. CROP TO SHOW ONLY π , K_S^0 , Λ .

The results indicate that not only charged pions but also neutral K_S^0 and Λ show little variation in multiplicity between the different $S_O^{(p_T=1.0)}$ extremes in the $N_{\text{SPD}_{\text{Trkts}}}$ I class, although more so for the latter two. This could be because of their smaller correlation with $S_O^{(p_T=1.0)}$ as they are neutral, or it may suggest strangeness enhancement. The models generally provide good agreement with the measured $\langle p_T \rangle$ of K_S^0 , while the results for Λ are best described by the Pythia Ropes model. The yields for K_S^0 are also largely consistent with the models presented, whereas the Pythia Monash model does not match the Λ yields.

5.12 Ratios to pions

Rather than focusing on the models' inability to match the absolute values, this measurement can provide an opportunity to study the underlying dynamics that affect the heavier and stranger particles. To achieve this, the K_S^0 and $\Lambda + \bar{\Lambda}$ p_T spectra are divided by the $\pi^+ + \pi^-$ spectra. The results are presented in Fig. 5.30 and Fig. 5.31.

The data reveals overall suppression of K_S^0 in jetty events and enhancement in isotropic events in the measured p_T range. The effect is observed, although in different ways, in all three event classes: $N_{\text{SPD}_{\text{Trkts}}}$ I (where spherocity is assumed to change mostly the $\langle p_T \rangle$), $N_{\text{SPD}_{\text{Trkts}}}$ I-III (where spherocity is assumed to change both the $\langle p_T \rangle$ and multiplicity), and V0M I (where spherocity is assumed to change mostly the multiplicity). The Λ ratios behave similarly but an enhancement in low- p_T is observed in the V0M I case, resembling typical radial flow signatures.

The presented models cannot describe the ratios to pions, although the trends relative to the $S_O^{(p_T=1.0)}$ -unbiased case are reproduced well.

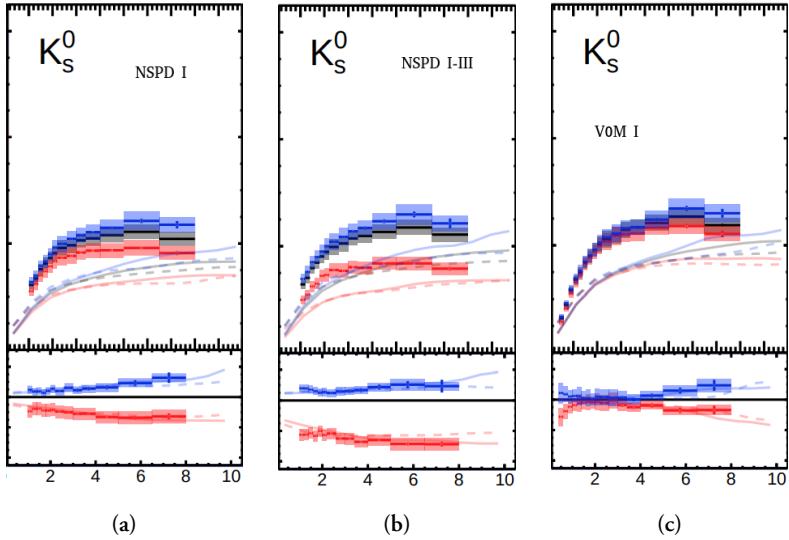


Figure 5.30: Ratios of transverse momentum spectra of K_S^0 to $\pi^+ + \pi^-$ in jetty, isotropic, and $S_0^{(p_T=1.0)}$ -unbiased in events with high-multiplicity classes (a) $N_{\text{SPD-Tkls}}^I$, (b) $N_{\text{SPD-Tkls}}^{I-III}$, and (c) V0M I. Statistical and systematic uncertainties are depicted as error bars and boxes, respectively.

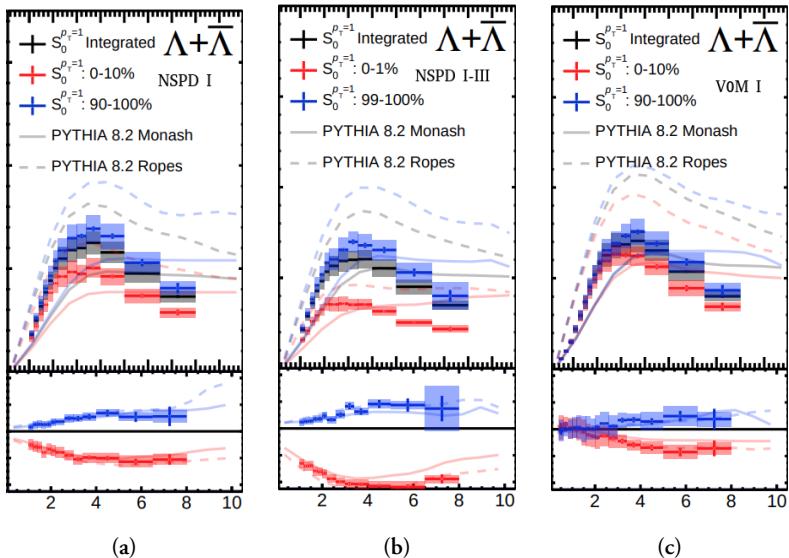


Figure 5.31: Ratios of transverse momentum spectra of $\Lambda + \bar{\Lambda}$ to $\pi^+ + \pi^-$ in jetty, isotropic, and $S_0^{(p_T=1.0)}$ -unbiased in events with high-multiplicity classes (a) $N_{\text{SPD-Tkls}}^I$, (b) $N_{\text{SPD-Tkls}}^{I-III}$, and (c) V0M I. Statistical and systematic uncertainties are depicted as error bars and boxes, respectively.

5.13 Baryon-to-meson ratio

The baryon-to-meson ratio $(\Lambda^0 + \bar{\Lambda}^0)/(2K_S^0)$ was investigated in this study, as it is a common observable used to measure the effects of radial flow, as discussed in Section X. To focus on the different functions of $S_O^{(p_T=1.0)}$ selection in the different high-multiplicity classes, the results for the $N_{\text{SPD}_{\text{Trkts}}}$ I and V0M I events are shown in Fig. 5.32. They are also juxtaposed with the p/π ratios [157]. The ratios reveal a significant increase when transitioning from jetty to isotropic events, indicating that the production of heavier Λ is systematically more suppressed in jetty events than K_S^0 . Similar observations have been made in other ALICE measurements studying jets (discussed further in ??) [158].

Although the increase in the ratio is consistent with the typical signatures of radial flow, the $N_{\text{SPD}_{\text{Trkts}}}$ I results do not reveal the depletion at low p_T and neither results show the shift of the peak to higher p_T , both of which are also its characteristic features. Additionally, the data are compared with the two tunes of Pythia 8, favouring the Ropes configuration.

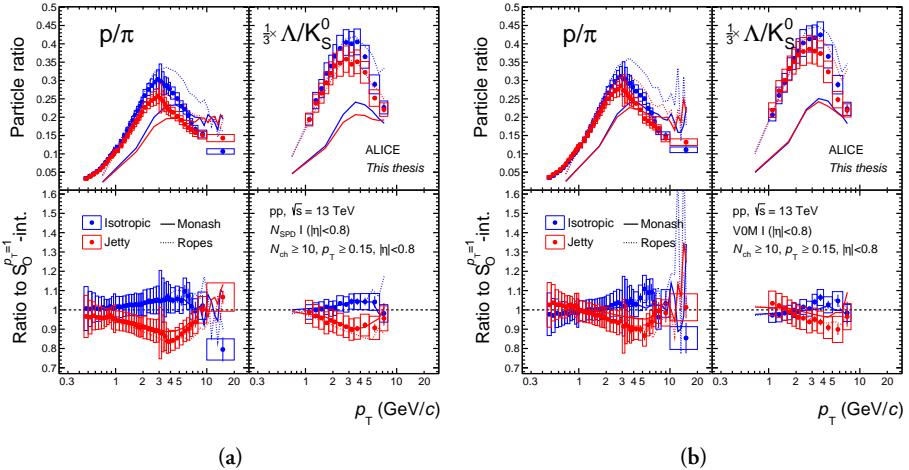


Figure 5.32: Baryon-to-meson ratios of the jetty (0 – 10%) and isotropic (90 – 100%) transverse momentum spectra of $p \rightarrow \pi$ and $\Lambda + \bar{\Lambda} \rightarrow K_S^0$ in the (a) $N_{\text{SPD}_{\text{Trkts}}}$ I and (b) V0M I high-multiplicity events of pp collisions at $\sqrt{s} = 13$ TeV. The bottom panels display the ratios to the $S_O^{(p_T=1.0)}$ -unbiased case. MC predictions are denoted by lines, statistical uncertainties as error bars, and systematic uncertainties as error boxes.

5.14 Ratio of integrated yields vs. $S_O^{(p_T=1.0)}$

The main contribution of this study is the investigation of strangeness production as a function of $S_O^{(p_T=1.0)}$. Yields of the Λ baryon were determined in $N_{\text{SPD}_{\text{Trkls}}}$ I and V0M I classes in the following intervals of $S_O^{(p_T=1.0)}$: 0–1%, 1–5%, 5–10%, 10–20%, 20–80%, 80–90%, 90–95%, 95–99%, and 99–100%. The integration was done in the measured p_T range rather than extrapolated. The $S_O^{(p_T=1.0)}$ -dependent Λ yields are then divided by $S_O^{(p_T=1.0)}$ -dependent pion yields as a reference. The effects of not extrapolating were studied and found to be unimportant, and while the extrapolation confirms the physics message of the default method, it also increases systematic uncertainties.

The results, together also with $p(|S| = 0)$ and $\Xi(|S| = 2)$ are reported in Fig. 5.33. For the $N_{\text{SPD}_{\text{Trkls}}}$ I events, there is a clear strangeness- and/or mass-dependent enhancement with increasing $S_O^{(p_T=1.0)}$. It is important to emphasize that in this event class, the N_{ch} multiplicity is basically fixed. Therefore, these effects are results of the different underlying dynamics of the collisions, specifically the varying hardness of involved scatterings, rather than merely an effect of increased N_{ch} .

Conversely, the V0M I events show no such dependence, and jetty and isotropic events appear to produce the same relative amount of strangeness. This goes against intuition that varying N_{ch} between jetty ($N_{\text{ch}} \rightarrow 15$) and isotropic events ($N_{\text{ch}} \rightarrow 30$, according to Fig. 5.20) should introduce an effect, in accordance with traditional studies of strangeness enhancement in pp collisions as a function of N_{ch} . This observation is not fully understood, but it suggests that the decrease in strangeness in jetty events due to the decrease in N_{ch} and increase in $\langle p_T \rangle$ is not trivial and counterbalanced by some other factors of the collision.

Finally, the results were compared with MC predictions, and only Pythia 8 Ropes accurately captures the observed trends and, to some extent, the magnitudes.

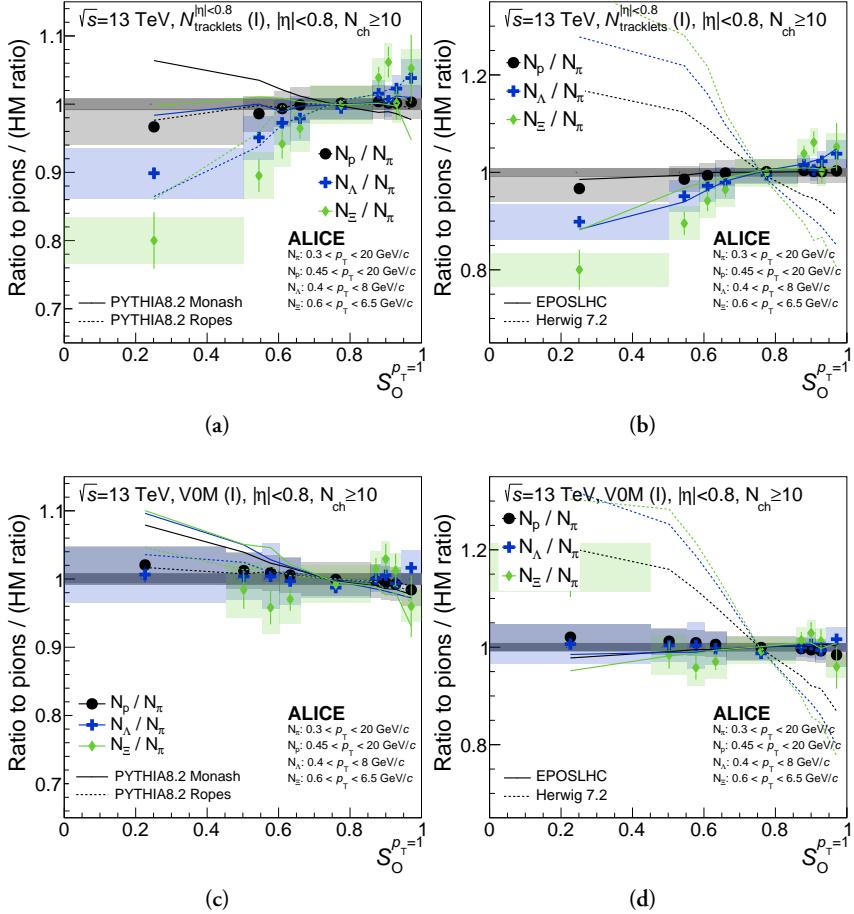


Figure 5.33: Double-ratios of integrated yields with respect to pions as a function of $S_O^{(p_T=1.0)}$ in the $N_{\text{SPD}_{\text{Trkts}}}^{\eta|<0.8}$ I and V0M I events. Statistical and total systematic uncertainties are shown by bars and boxes, respectively. The curves represent different model predictions of the same measurement: (a),(c) Pythia Monash and Ropes, (b),(d) EPOS LHC and Herwig.

Chapter 6

Underlying Event Activity

Chapter 7

Discussion of results, conclusions, and outlooks

7.1 Summary of the research goals

The first part of this dissertation provides an introduction to quantum chromodynamics (QCD), a theory of the fundamental strong force of the Universe, and explains why QCD interactions involving low-momentum (soft) transfers cannot be calculated from first principles, unlike hard processes, which can use perturbation theory. Next, the importance of studying QCD matter at extreme conditions is discussed, specifically the plasma of deconfined quarks and gluons QGP.

Furthermore, the dissertation presents an overview of experimental studies that investigate the properties and signatures of the QGP in AA collisions and their dependence on collision centrality, which can be related to energy densities as well as final-state multiplicities in the system. Special attention is given to the phenomena of strangeness enhancement, where strange particles are produced more abundantly in more active collisions, and collective flow, where the hydrodynamic behavior of the plasma affects the dynamics of final-state hadrons.

The dissertation also describes the challenges to traditional paradigms that pp collisions were thought to be incapable of producing extremes of QCD matter, and lists the observations resembling traditional QGP phenomena in these small system collisions. Moreover, it explains the intricacies of isolating the physics behind these phenomena, given that event multiplicity arises from non-perturbative (softer) as well as perturbative (harder) processes, is susceptible to large fluctuations, and cannot be directly linked to the energy density.

The goal of this dissertation is to provide a clearer and more differential study of QGP phenomena, namely strangeness enhancement and radial flow, and to elucidate the roles played by both soft and hard processes. The aim is to identify and utilise observables that isolate events with extreme activity resulting from non-perturbative processes, where novel physics may be at play, such as the formation of a QGP-like state or complex interactions of overlapping QCD fields. Two state-of-the-art phenomenological models are used to represent these paradigms: EPOS LHC, which includes QGP droplets, and the Ropes tune of Pythia 8, which allows for the merging of strings into higher-tension fields. These observables should also be capable of isolating event activity extremes dominated by perturbative physics at the other end of the spectrum.

Given the complexity of the physics picture and its need to be studied from various angles, these measurements are unlikely to single-handedly confirm or reject the “big hypothesis” of whether QGP can be formed in pp collisions. Nevertheless, they will be able to provide valuable insight into the underlying physics processes in hadronic and partonic interactions and their deeper understanding. Furthermore, the measurements have the potential to significantly discriminate between different phenomenological models.

7.2 Highlights of the $S_O^{(p_T=1.0)}$ measurement

Chapter 5.7.2 introduces measurements of the neutral, weakly decaying K_S^0 , Λ , and $\bar{\Lambda}$ as a function of transverse spherocity $S_O^{(p_T=1.0)}$ in pp collisions at $\sqrt{s} = 13$ TeV using the ALICE detector at the LHC. This observable describes the geometrical shape of the charged particles produced in the collision and provides a simple, albeit effective discriminator between pencil-like events, dominated by a di-jet coming from a single hard partonic scattering, and isotropic events, where particles are produced from multiple sources involving lower momentum transfers.

These measurements provide the first ever experimental results of these strange particles as a function of an event shape observable. Moreover, they develop and utilise the so-called unweighted spherocity, a modification from its traditional form S_O , which is easier to compare with phenomenological models. The results are fully corrected and come with a thorough investigation of their experimental uncertainties.

In this dissertation, the following results are presented: transverse momentum p_T spectra, average $\langle p_T \rangle$, ratios of p_T spectra to pions, baryon-to-meson ratios, and integrated production yields as a function of $S_O^{(p_T=1.0)}$ in high-multiplicity events. The following findings can be highlighted:

1. Figure 5.20 shows that $S_O^{(p_T=1.0)}$ has interplays with the event multiplicity:
 - If high multiplicity is determined at forward-rapidity (V0M), $S_O^{(p_T=1.0)}$ varies mostly the mid-rapidity multiplicity, leading to more constant increases and decreases in particle spectra between jetty and isotropic events (Fig. 5.27).
 - If high multiplicity is determined at mid-rapidity ($N_{\text{SPD}_{\text{Trkts}}}$), $S_O^{(p_T=1.0)}$ varies mostly the $\langle p_T \rangle$, leading to similar yields but significant hardening and softening between jetty and isotropic events (Fig. 5.27 and Fig. 5.29).
2. Generally, the isotropic spectra are closer to the $S_O^{(p_T=1.0)}$ -unbiased spectra than the jetty are, suggesting that average high-multiplicity collisions are somewhat isotropic, whereas jetty events are outliers.
3. Ratios to pions in Fig. 5.31 reveal that relatively to pions, K_S^0 and Λ are consistently suppressed in jetty events and enhanced in isotropic events in $N_{\text{SPD}_{\text{Trkts}}}$ classes. This effect is different in the V0M class.
4. The baryon-to-meson ratios $(\Lambda^0 + \bar{\Lambda}^0)/(2K_S^0)$ in Fig. 5.32 display an enhancement of Λ in intermediate p_T in isotropic events. This is expected with observing radial flow, however, its other characteristic features, such as shift of the peak, are not seen.
5. Ratios of integrated yields to pions in Fig. 5.33 show a characteristic strangeness enhancement behaviour for the $N_{\text{SPD}_{\text{Trkts}}}$ class and the effect increases with increasing the mass and strangeness content. This result is the first observation ever of strangeness enhancement that occurs with (mostly) constant multiplicity. In the V0M class, the effect is not observed.
6. Results are compared with selected MC predictions, mostly favouring Pythia 8 Ropes and EPOS LHC over Pythia 8 Monash, although with varying degrees of success.

7.3 Highlights of the R_T measurements

Chapter ?? presents measurements of the production of K_S^0 , Λ , and $\bar{\Lambda}$ in pp collisions at $\sqrt{s} = 13$ TeV and their dependence on the relative underlying event activity R_T . This quantity controls the magnitude of the underlying event, which consists of many softer particles produced through multiple partonic interactions (MPIs) and other sources, and is unrelated to the primary hard partonic scatterings and its fragmentation. It acts as a dial, selecting events resembling ee collisions ($R_T \rightarrow 0$, $n_{\text{MPI}} \rightarrow 1$)

and AA collisions ($R_T \rightarrow \infty$, $n_{\text{MPI}} \rightarrow \infty$). Particle production is studied in three azimuthal regions based on the highest- p_T track, which serves as a proxy for the primary scattering axis: Towards, Away, and Transverse.

Moreover, to distinguish between the contributions to the underlying event from softer (MPIs) and harder (wide-angle radiation ISR/FSR) interactions, this approach is extended to divide the Transverse region further into Transverse-min and Transverse-max based on the number of particles, and define the classifiers $R_{T,\text{min}}$ and $R_{T,\text{max}}$, accordingly. At the time of conducting this measurement, the $R_{T,\text{min}}$ observable is expected to be the cleanest probe of the $\langle n_{\text{MPI}} \rangle$, as the harder contributions are captured in the $R_{T,\text{max}}$ quantity.

The measurements presented in this dissertation are the first ever experimental results of K_S^0 , Λ , and $\bar{\Lambda}$ as a function of R_T , as well as the first use ever of $R_{T,\text{min}}$ and $R_{T,\text{max}}$ on identified particles. Considerable effort has been required to experimentally use and understand these observables, particularly in the choice of tracks used, the treatment with Bayesian unfolding, and the quantification of systematic uncertainties.

The following results are focused on: p_T spectra, $\langle p_T \rangle$, baryon-to-meson ratios, and self-normalised yields. Key outcomes of this study are:

1. Figure ?? implies a much steeper increase of $\langle p_T \rangle$ of K_S^0 with $R_{T,\text{max}}$ than $R_{T,\text{min}}$.
2. The transverse momentum spectra in Fig. ?? and Fig. ?? show:
 - In the Toward region, the spectra in different R_T events converge at high p_T to the R_T -integrated values, corresponding to the dominance of jet.
 - In the Transverse region, p_T spectra continue hardening with increasing R_T , similarly to the picture in multiplicity measurements.
 - In contrast, the Transverse-min region seems to plateau.
3. The baryon-to-meson ratios in Fig. ?? reveal typical radial flow features in high- R_T events, including the enhancement of intermediate- p_T baryons, their depletion at low- p_T , as well as shifts of the peaks. Moreover,
 - In the Toward region, the effect is the largest. This is due to the mixing of jet- and UE-related particle production, which shows largely different $(\Lambda^0 + \bar{\Lambda}^0)/(2K_S^0)$ ratios.
 - The effect is comparable among the Transverse, Transverse-min, and Transverse-max regions, with only small hints of being slightly bigger in the Transverse-min case. This suggests that the harder and softer components of UE affect the relative production of Λ and K_S^0 in a similar fashion.

4. The self-normalised yields in Fig. ?? of K_S^0 and Λ are consistent with π and p , respectively, in the Toward and Away regions, although more experimental precision is needed. In the Transverse region, the effect of auto-correllation is apparent for the charged particles. Moreover, the K_S^0 and Λ yields rise more slowly in the Transverse-min cases than the Transverse-max.
5. The experimental data favour the Pythia 8 Ropes predictions. EPOS LHC does not display the right amount of sensitivity to the azimuthal region and the Pythia 8 Monash tune underestimates the effects of radial flow. This implies that Colour Reconnection is somewhat insufficient to describe the flow-like behaviour.
6. Generally, higher values of $R_T/R_{T,\min}/R_{T,\max}$ observables need to be reached in order to isolate the different behaviours of softer MPIs and harder ISR/FSR and further test the MC predictions.

7.4 Conclusions and Outlooks

One paragraph wrap-up? One paragraph about possible improvements One paragraph about other areas for exploration

Part IV

Appendices

Appendix A

List of Acronyms

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Appendix B

Mathematical Derivations

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Appendix C

Complementary Material

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Appendix D

Scientific Publications

Author contributions

Paper I: Title paper 1

I participated in developing the theory and wrote the simulation software. I participated in writing the manuscript.

Paper II: Title paper 2

I participated in developing the theory and writing simulation software. I participated in writing the manuscript.

Paper I

S. Doctor and B. someone

An Exact Ewald Summation Method in Theory and Practice

The Journal of Physical Chemistry A, 2020, 124(19), pp. 3943-3946

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Paper II

S. Doctor, B. someone, C. another and D. another

Grand canonical simulations of ions between charged conducting surfaces using exact
3D Ewald summations

Physical Chemistry Chemical Physics, 2020, 22(24), pp. 13659-13665

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