

Jet-like two-particle correlations in p-Pb collisions

Emilia Leogrande



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Jet-like two-particle correlations in p–Pb collisions

Jet-achtige twee-deeltjes correlaties in p–Pb botsingen

(met een samenvatting in het Nederlands)

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We are pathetic. We are stars.

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Introduction

The Large Hadron Collider (LHC) at CERN, Geneva, is the world's biggest particle collider of all times. Built in about a decade, since 2009 it has been able to collide protons and heavy ions at unprecedented energies. The LHC is the ideal experimental ground, at present, to explore the subatomic structure of matter, to investigate the properties of the fundamental forces and to advance the human understanding of the building blocks of our universe.

Among the four largest experiments installed on the LHC ring, ALICE is the one built with the main purpose of studying the system formed when heavy ions collide. The high energy densities reached in such collisions, in fact, are able to recreate in the lab what it is theorized to have been the status of matter, so-called Quark Gluon Plasma (QGP), few microseconds after the Big Bang.

Moreover, a new window in the exploration of the subatomic world has been opened with the collisions of a proton beam with a lead-ion beam in the LHC. By analyzing the data of such collisions, intriguing results have emerged, which have questioned some of the statements of the physics community. *Is the QGP formed also in such small colliding systems? Could the short-lived state of matter approach thermal equilibrium fast enough to form a QGP-droplet ready to evolve as in heavy-ion collisions?* Since the first proton-lead collisions (a small pilot run in September 2012 and a full run in January-February 2013), theorists and experimentalists are working together to give answers to many arisen questions. However, a final word is not yet available and further investigations are needed to unravel the puzzle

of the physics of small systems. The p–Pb run at higher energy, scheduled for December 2016, is eagerly awaited.

The interpretation of physics results is further complicated by the composite nature of colliding particles: protons and neutrons are constituted of elementary particles, also known as *partons*. In each proton–proton (or proton–neutron) collision, multiple interactions among partons can occur. The effect of such multiple parton interactions (MPI) is reflected onto the final state of particle production and affects many physical observables. It is evident that this phenomenon plays the role of a background for many physics channels of discovery, but it is also important on its own right to improve the knowledge about the mechanism of particle production. The first evidences of MPI, indeed, came from the study of charged-particle multiplicity distributions.

Experimental studies on MPI are even more crucial in the regime of small transverse momentum of the produced particles, where theoretical limitations make calculations only approximate. In this regime, each parton–parton collision produces minijets, i.e. jets that, because of their low energy, are not identifiable as cone structures.

In this thesis, the statistical ‘two-particle correlations’ method is used to reconstruct the minijets and, through them, to derive an observable which is proportional to the number of MPI at low transverse momentum in p–Pb collisions. Thus, this work aims at constraining theoretical models with experimental findings on particle production at low momentum and at casting some light on the puzzling realm of small colliding systems, namely on the possible effects of collectivity on the minijet peaks.

The thesis is divided in three sections, as follows.

The first section reviews the phenomenological framework, relevant for the physics covered by this thesis. Chapter 1 regards the physics of hadrons and heavy ions, from the Standard Model of particle physics, with a special focus on the theory of strong interactions (QCD), to the phase diagram of hadronic matter, with a deeper insight into the Quark Gluon Plasma (QGP)

Introduction

and how to recreate it in the lab. Finally, a brief overview of the major results concerning small colliding system, such as p–Pb at LHC, is presented. Chapter 2 focuses on the physics of multiple parton interactions (MPI), from their first experimental signature in multiplicity distributions to the implementation in modern Monte-Carlo models, in particular PYTHIA.

The second section collects the information about the experimental setup. Chapter 3 describes the LHC accelerating chain, with a closer look at the parameters during the p–Pb data-taking period, and the ALICE detector, with a specific focus on the subdetectors used in the present analysis. Chapter 4 analyzes the classification of events according to centrality, discussing specifically the biases arising in such event classification in the case of p–Pb collisions.

The third section includes my PhD work. Chapter 5 presents the analyzed datasets, from their recording to the selection of events and tracks, specific for the present analysis. Chapter 6 describes the analysis procedure and contains the definition of the main observables. Chapter 7 summarizes the systematic uncertainties. Finally, Chapter 8 presents and discusses the physics results.

Part I

PHENOMENOLOGY

1

The physics of hadrons and heavy ions

This chapter provides an introduction first to the theoretical framework and then to the experimental background in which the work of this thesis takes place.

The Standard Model of particle physics is summarized in Sec. 1.1, with a particular focus on the theory of Quantum ChromoDynamics in Sec. 1.2.

Then, the evolution of a system of QCD matter in terms of thermodynamical variables is explored in Sec. 1.3, with a deeper insight into the Quark Gluon Plasma (QGP) phase and the way to recreate it with heavy-ion collisions.

Finally, a brief review of the puzzling results concerning small systems, being those the topic of this work, is given in Sec. 1.4.

1.1 The Standard Model

Four kinds of fundamental forces, through which particles interact, are known to exist: gravitational, electromagnetic, strong and weak. The three latter are described in an unified way in the Standard Model (SM), via the exchange of specific particles of spin 1 called *bosons*. Each interaction is me-

1. The physics of hadrons and heavy ions

diated by a different kind of boson: the photon (γ) for the electromagnetic, the $W^\pm$ and Z for the weak and 8 gluons (g) for the strong processes.

The model contains also a description of all the elementary constituents of matter, point-like particles of spin 1/2 called *fermions*, divided in two species: *leptons* and *quarks*. These are organized in three families, each containing a pair of leptons and of quarks, summarized in Fig. 1.1.

Family	Quarks			Leptons		
	Name	Charge	Mass	Name	Charge	Mass
1	u	$2/3\ e$	$1.5 - 3.3\ \text{MeV}/c^2$	e^-	$-e$	$0.511\ \text{MeV}/c^2$
	d	$-1/3\ e$	$3.5 - 6.0\ \text{MeV}/c^2$	ν_e	0	$< 2\ \text{eV}/c^2$
2	c	$2/3\ e$	$1.27^{+0.07}_{-0.11}\ \text{GeV}/c^2$	μ^-	$-e$	$106\ \text{MeV}/c^2$
	s	$-1/3\ e$	$104^{+26}_{-34}\ \text{MeV}/c^2$	ν_μ	0	$< 0.19\ \text{MeV}/c^2$
3	t	$2/3\ e$	$171.2 \pm 2.1\ \text{GeV}/c^2$	τ^-	$-e$	$1.78\ \text{GeV}/c^2$
	b	$-1/3\ e$	$4.2^{+0.17}_{-0.07}\ \text{GeV}/c^2$	ν_τ	0	$< 18.2\ \text{MeV}/c^2$

Figure 1.1: Quarks and leptons. Table from [1].

All charged fermions experience the electromagnetic force. In addition, leptons and quarks interact via the weak force. Moreover quarks, also equipped with the so-called *color* charge, are allowed to interact via the strong force. A summary of the different kinds of interactions, of their ranges and of the respective particles involved, is given in Fig. 1.2. In nature, quarks are grouped in colorless objects called *hadrons*, which are divided in *baryons* and *mesons* if they are composed of 3 valence quarks (or antiquarks) or a valence quark-antiquark pair respectively ¹.

The last piece of the Standard Model, theorized by Brout, Englert [4] and Higgs [5], has been recently² discovered at CERN [6, 7]. This particle of mass around $126\ \text{GeV}/c^2$ called *Higgs boson* is responsible for giving mass to the vector bosons of the SM.

The SM, in spite of its theoretical consistency and the many experimen-

¹Recent results by LHCb Collaborations have shown the existence of exotic particles composed of five quarks, called pentaquarks [2]

²CERN announced the discovery the 4th July 2012.

1.2 Quantum ChromoDynamics

Force	Strength	Gauge Boson(s)	Applies on
Strong force	1	8 Gluons (g)	Quarks, gluons
Electromagnetic force	$\sim 10^{-2}$	Photon (γ)	All charged particles
Weak force	$\sim 10^{-7}$	$W^\pm, Z^0$	Quarks, leptons
Gravitation	$\sim 10^{-39}$	Gravitons	All massive particles

Figure 1.2: Fundamental forces. Table from [3].

tal confirmations, is not the ultimate theory. In fact, it does not include the gravitational interaction, which exists among all massive particles. Besides, the SM has many free parameters that cannot be predicted from first principles, e.g. the Higgs boson mass itself. In fact, one of the biggest open points concerns the exact value of the Higgs boson mass (the smallness of its mass is linked to the difference of strength between the weak and the gravitational interaction. This is known as hierarchy problem [8]). Among the new theories Beyond Standard Model (BSM) [9], supersymmetry (SUSY) [10] would be able to solve the hierarchy problem, together with other puzzles of the SM, such as neutrino oscillations [11].

1.2 Quantum ChromoDynamics

Within the framework of the SM stands Quantum ChromoDynamics (QCD), the theory of strong interactions. QCD is a quantum field theory with symmetry group $SU(3)_C$, where C refers to the color, which plays the equivalent role of the electric charge in Quantum ElectroDynamics (QED). Unlike the QED force carriers (the photons), which are electrically neutral, the QCD force carriers (the 8 gluons) have non-zero color charge. Thus, self-interactions between gluons are possible. Quarks and gluons, the building blocks of QCD, are the elementary constituents of hadrons³.

The main parameter of QCD is the coupling constant α_S . It characterizes the strength of the strong force. In spite of its name, this quantity

³In this context, it is often referred to quarks and gluons indifferently as *partons*.

1. The physics of hadrons and heavy ions

varies with the momentum transfer Q of the interaction, according to the DGLAP (Dokshitzer - Gribov - Lipatov - Altarelli - Parisi) equations [12], as shown in Fig.1.3.

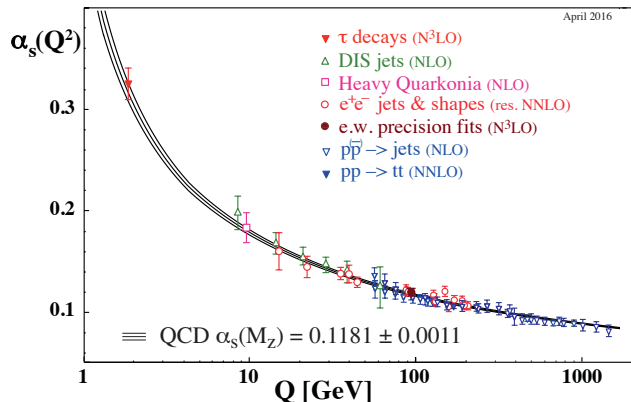


Figure 1.3: Summary of measurements of α_s as a function of the respective energy scale Q . Figure taken from [13].

Two regimes can be distinguished:

- towards small momentum transfer Q , which means large probed length scale, the coupling increases. The potential energy between partons increases linearly with the distance. This is a necessary but not sufficient condition for explaining the experimentally observed phenomenon known as **confinement**, namely the fact that colored objects are not found in nature as isolated. However, given the trend of potential energy with distance, confinement can be explained in the picture of the so-called *string model* [14]: as the distance between partons increases, the color tube connecting them (created by the gluon strong field) stretches; its energy, proportional to its length, increases, until it is energetically more favorable the creation of a new quark-antiquark pair connected by shorter, thus less energetic, color strings.
- towards large momentum transfer Q , which means small probed length

1.3 The phase diagram of the hadronic matter

scale, the coupling decreases. In the limit $\alpha_S \rightarrow 0$ quarks and gluons behave as free particles. This condition is known as **asymptotic freedom**.

In order to calculate cross sections of physics processes in hadron interactions, it is necessary to know how the partons are distributed inside the hadron⁴ and the cross sections at the parton level. The QCD factorization theorem [15] states that these two contributions are, indeed, factorizable. The former is included in the parton distribution functions (PDFs) [16], which contain the probability of having a parton of a certain type carrying a fraction of the momentum of the original hadron, given the energy scale Q of the collision. The parton cross sections are calculable in *perturbative* QCD (pQCD) [17–19], providing that the coupling constant α_S is small, i.e. at the short-distances scale that characterizes asymptotic freedom. In the confinement limit, instead, the coupling constant is large. Therefore, the perturbative approach is no longer valid. This is the case for *soft processes* at low transverse momenta, which are dominant in hadron collisions. One possibility to sidestep this difficulty is to use *Lattice QCD* [20], a computational technique which discretizes the space-time, by placing quarks in lattice structures with a fixed finite distance, and allows to calculate numerically integrals and cross sections.

The boundary between non-perturbative and perturbative QCD is not very well defined. Experimental results from RHIC on pion production in pp collisions [21] show good agreement with the perturbative Next-to-Leading Order (NLO) calculations [22] down to $p_T \simeq 1$ GeV/ c .

1.3 The phase diagram of the hadronic matter

At the ordinary conditions of temperature and pressure, quarks and gluons are confined into hadrons. Lattice QCD allows predicting the properties of

⁴The 3(2) quarks which determine the quantum numbers of baryons(mesons) are called *valence quarks*; apart from these, any hadron may contain an indefinite number of *virtual* (or *sea*) quarks, antiquarks, and gluons which do not influence its quantum numbers. They are off-shell particles popping up as an effect of quantum fluctuations.

1. The physics of hadrons and heavy ions

small volumes of hadronic matter, either in thermal equilibrium or under changing thermodynamical conditions. In this way, one can reconstruct a *phase diagram* for the hadronic matter. This is shown in Fig. 1.4.

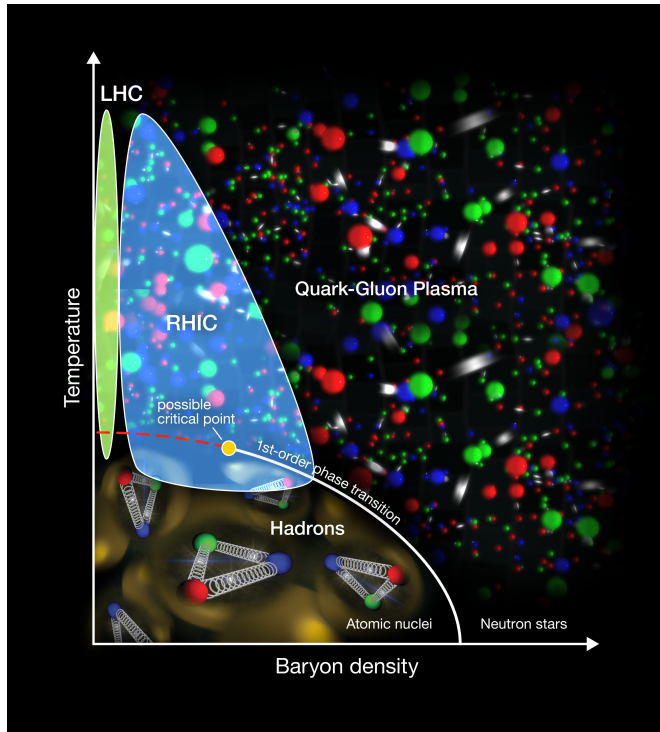


Figure 1.4: Qualitative representation of the phase diagram of the hadronic matter. Image: Brookhaven National Laboratory.

When the temperature is high, $T \gg \Lambda_{\text{QCD}}^5$, due to asymptotic freedom quarks and gluons are deconfined in a plasma-like state, commonly referred to as Quark Gluon Plasma (QGP) [23]. At lower temperatures, the coupling is strong and the partons are bound into hadrons behaving like a gas. The

⁵ $\Lambda_{\text{QCD}} \simeq 200 \text{ MeV}$ [13] is one of the main parameter of the theory of strong interactions.

1.3 The phase diagram of the hadronic matter

transition between QGP and hadron gas state depends on the net baryon density. When the latter is large, a first order transition occurs. At very small or zero baryon density, instead, the transition is a smooth *crossover* [24]. The point of the transition curve in which it changes nature, from first order to crossover, is a point of criticality, which is still under debate as Lattice QCD calculations are only approximate at non-zero baryon density [25]. On the other hand, reliable results are available for the Equation of State (EoS) of the QCD matter for baryon-free systems. As shown in Fig. 1.5, both the energy density ϵ and the specific heat C_V (normalized respectively to T^4 and T^3) present a steep increase with temperature, which is due to the change in the number of degrees of freedom occurring at the phase transition between the confined and the deconfined state. The temperature at which the transition takes place is located at $T_C \sim 155$ MeV [24]. Also shown as a dashed line is the limit for a non-interacting gas of quarks and gluons.

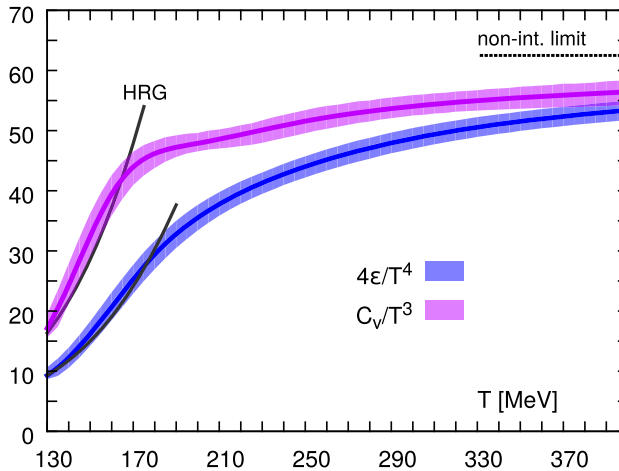


Figure 1.5: Equation of State from Lattice QCD at zero baryon density. Figure taken from [26].

The conditions of high temperature and vanishing net baryon density

1. The physics of hadrons and heavy ions

are thought to be those of the early universe few microseconds after the Big Bang and can be reproduced by colliding heavy ions in particle accelerators. This is discussed more in detail in the next paragraph. The region of very low temperature and high baryon density, instead, corresponds supposedly to the core of neutron stars [27].

1.3.1 QGP in the lab: heavy-ion collisions

Heavy-ion collisions were performed in Brookhaven, by AGS (Alternating Gradient Synchrotron) with Au–Au up to $\sqrt{s_{\text{NN}}} = 11.5$ GeV and by RHIC (Relativistic Heavy Ion Collider) with gold, copper and uranium ions, at center-of-mass energies from 7.7 to 200 GeV per nucleon pair [28].

At CERN, the first lead-ion collisions were performed with the SPS (Super Proton Synchrotron) up to $\sqrt{s_{\text{NN}}} = 158$ GeV. The highest nucleon–nucleon center-of-mass energies with these ions have been reached by LHC with $\sqrt{s_{\text{NN}}} = 2.76$ TeV in Run I [29] and $\sqrt{s_{\text{NN}}} = 5.02$ TeV in Run II [30].

At low collision energies ($\simeq 5 - 10$ GeV per nucleon), the Lorentz-contracted nuclei stop each other, leading to the formation of a *fireball* of high baryon density. At high enough energies, instead, ($\simeq 100 - 200$ GeV per nucleon), the participant nucleons are slowed down but can still have momentum to escape from the collisional region. In this way, a large amount of energy is deposited in a small volume, which results in high energy density matter with low net baryon content [31]. The bigger the collision energy, the closer one gets to a vanishing baryon content at midrapidity.

According to the displacement of the two nuclei at the time of collision, the initial state is characterized by QCD matter arranged in a pattern of geometrical eccentricities. Such matter evolves according to these phases:

- **Pre-equilibrium**

In this early stage, the constituents collide until they reach a local thermal equilibrium. The duration of this process cannot be measured directly, but it is found to be compatible to $\simeq 1$ fm/ c , from comparison of data and models which assume such short thermalization times [32]. The large energy density fluctuations produce fluctuations in

1.3 The phase diagram of the hadronic matter

the initial shape of the collisional region.

- **QGP**

The fluctuating geometrical eccentricities generate pressure gradients. These translate in momentum anisotropies, as a consequence of the anisotropic push due to the pressure. At first, QGP was thought to be a weakly-interacting system [23]. However, with RHIC measurements [28] it became more evident that at the reached energy densities ($10^3 \text{ GeV}/fm^3$ in head-on collisions) the mean free path is small compared to the system size. Thus, the system turns out to be a strongly interacting one and it behaves like a fluid. In fact, it can be successfully described by hydrodynamical models, which take as input the Lattice-QCD EoS⁶. Even though early measurements [28] were compatible with an ideal fluid behaviour, with LHC results [29] it emerged that viscous effects, although small, cannot be neglected. The QGP is characterized by two types of viscosity: a *bulk* and a *shear* viscosity. The latter is particularly important because it degrades the medium's ability to convert initial transverse pressure anisotropies into final transverse momentum anisotropies. However, being the value of the shear viscosity small, traces of initial fluctuations survive to the final hadronic stage. For a detail discussion, the reader is referred to [33]. In its expansion, the QGP cools down until it reaches the critical temperature for the phase transition from deconfined to confined matter.

- **Hadronization**

Right after the critical temperature for the phase transition, the partons in the plasma undergo the hadronization phase, which consists of two subsequent stages: a) *fragmentation*, i.e. partons of the medium shower and create a spray of quarks and gluons, and b) *coalescence* or *recombination*, i.e. produced (shower or thermal) partons which

⁶The success of hydrodynamical models in reproducing data is also a confirmation of the reliability of Lattice-QCD EoS and its prediction of a change in the number of degrees of freedom.

1. The physics of hadrons and heavy ions

are close in phase space recombine to form hadrons. This status of matter becomes too dissipative to be described by a macroscopical hydrodynamical theory. In fact, this stage is successfully modelled by a microscopic kinetic description [34].

- **Chemical Freeze-out**

Inelastic collisions cease and relative hadronic abundances are frozen (except for decays of resonances). Provided that these hadrons are created from a thermalized medium of quarks and gluons, the baryon and antibaryon abundances are studied in the context of a thermal statistical model [35].

- **Kinetic Freeze-out**

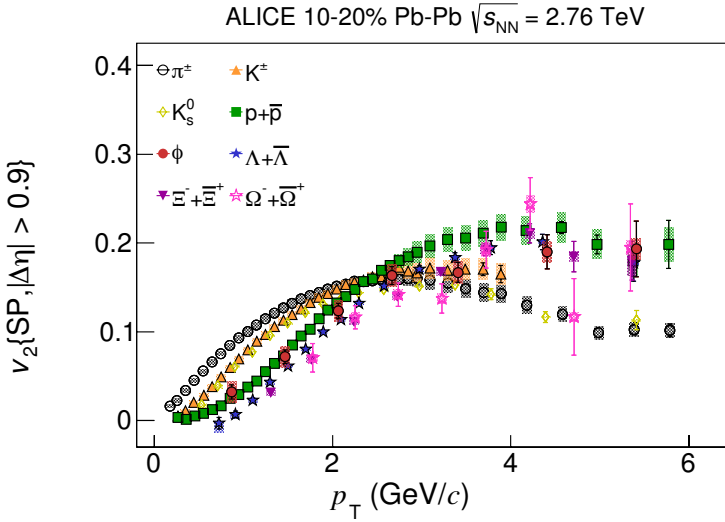
Elastic collisions cease and particles propagate to the detectors in a free streaming mode. Thus, they offer a snapshot of the situation at the freeze-out temperature.

Kinetic freeze-out is reached in about 10 fm/ c after the collision, therefore the QGP is not directly measurable. However, final-state hadrons preserve information on the whole evolution of the system. In particular, the hot QCD matter created when colliding heavy ions can be seen as composed of *soft*, i.e. low-momentum, and *hard*, i.e. high-momentum, partons. The former constitute the dominant contribution and behave as a relativistic fluid. The latter are produced at the very early stages of the collisions and traverse the QCD medium, losing energy in the process. When these partons hadronize, soft and hard particles are formed, which are used as probes to measure the properties of the QGP.

In the realm of soft probes, the anisotropy of particle emission in phase space, i.e. the response of the medium to the initial spatial anisotropy, can be quantified by the so-called *elliptic flow*. This is the second-order coefficient v_2 of the Fourier decomposition which describes the azimuthal particle distribution. Figure 1.6 shows the p_T -differential v_2 for several particle species and for $p_T < 2$ GeV/ c reveals a smaller value for heavier particles. This mass ordering is due to a combined effect of elliptic flow and

1.3 The phase diagram of the hadronic matter

radial flow, the latter being an isotropic transverse expansion which affects the p_T -spectra. Elliptic flow v_2 measurements suggest that the QGP behave as an almost perfect liquid, i.e. with small value of shear viscosity [36].



ALI-PUB-82653

Figure 1.6: p_T -differential elliptic flow for different particle species. Figure taken from [36].

As regards the hard probes stemming from early hard scatterings, those lose energy via gluon radiation and multiple collisions within the medium. This can lead to a suppression in the number of particles, which is estimated via the nuclear modification factor R_{AA} , namely the ratio between the particle yield in heavy-ion collisions divided by the particle yield in pp collisions, normalized to the average number of binary nucleon–nucleon collisions. This is plotted in Fig. 1.7 as a function of the transverse momentum, showing that the suppression is stronger at LHC, meaning that the QGP is hotter, than at RHIC energies [37].

1. The physics of hadrons and heavy ions

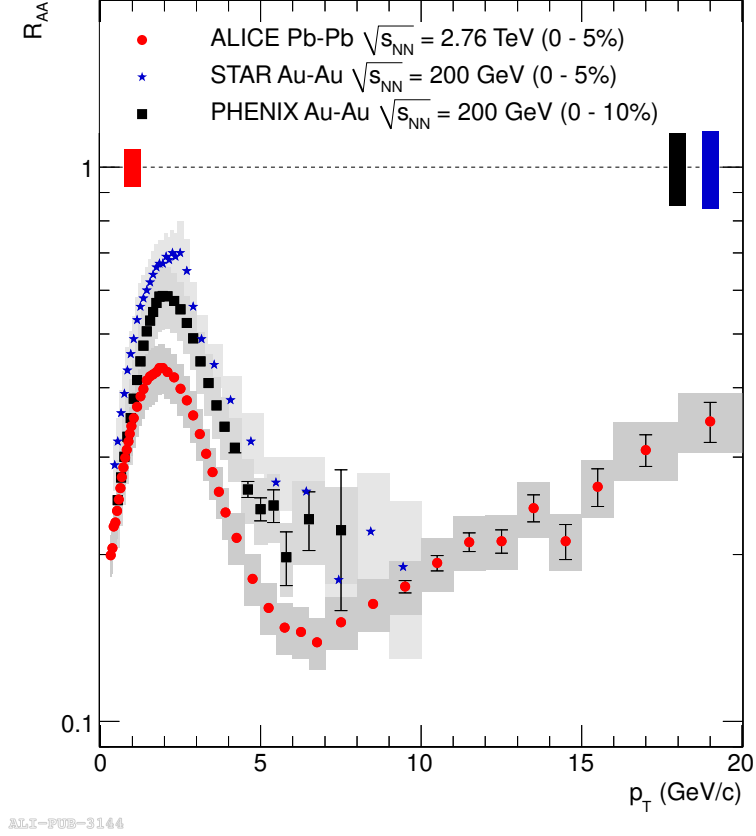


Figure 1.7: Nuclear modification factor versus transverse momentum in ALICE, STAR and PHENIX. Figure taken from [37].

Many other observables are useful to measure the QGP properties. For an extensive review, the reader is referred to [38].

1.4 Phenomenology of small systems: p–A collisions

As just shown for the R_{AA} , measurements of probes of the QGP formed in heavy-ion collisions are usually compared with measurements of the same observables in proton–proton collisions, in which, due to small matter density, no QGP is expected to be formed.

Similarly, intermediate systems, such as d–Au collisions at RHIC [39] and p–Pb collisions at LHC [40], are valuable as references for the heavy-ion data [41]. In fact, such systems were thought to be still too small to create the energy density required to form the QGP. Therefore, the only expected effects should be those not due to the presence of a hot deconfined medium. Those so-called *cold nuclear matter* (CNM) effects include the nuclear modification of parton distribution functions (nPDF) [42, 43], which take into account the initial longitudinal distribution of partons inside the nuclei. In addition, in the kinematic region where QCD is non-perturbative (small x), high gluon densities are reached, which may lead to a saturation in the form of a color glass condensate (CGC) state [44].

The measured charged-particle density in p–Pb collisions [45] is well described by either two-component models, which combine perturbative QCD and soft processes and include nPDF as initial distributions, or saturation models, that, through coherence effects, are characterized by reduced number of soft gluons available for the particle production (details on the models in [45]).

Measurements of probes sensitive to energy loss tend to discard the possibility of the formation of a medium. The nuclear modification factor R_{pPb} [46], for instance, which is the ratio between the particle yield in p–Pb and in pp collisions scaled by the average number of binary nucleon–nucleon collision, is equal to unity up to 30 GeV/ c . Similarly, the dijet transverse momentum imbalance [47] is reproduced with p–Pb simulations of embedded pp PYTHIA events, not requiring any QGP. Moreover, the suppression of J/ψ in the forward direction only [48] is also compatible with effects of cold nuclear matter.

1. The physics of hadrons and heavy ions

In the regime of soft probes, a large correlation in the difference in rapidity and azimuthal angle for pairs of particles, elongated in $\Delta\eta$ and peaked at $\Delta\varphi = 0$ and π (so-called *double ridge*, shown in Fig. 1.8), was observed in high-multiplicity p–Pb collisions [49]. Similar feature appears in Pb–Pb events, where it is understood to be due to a collective expansion of the hot medium, but also saturation models [44, 50] that do not invoke collective phenomena are in good agreement with some of the measurements. However, the elongated structure at $\Delta\varphi = 0$ emerges also in high-multiplicity pp events [51]. The origin of these structures in small systems is still puzzling: whether they arise from initial-state (such as CGC) or from final-state effects (such as rescatterings or collective phenomena in high-density systems) or from the interplay of the two is still under vivid debate.

Toward a deeper comprehension of the phenomenon, the double ridge has been extracted from correlations of identified particles [52]. The p_T -differential second-order coefficient v_2 of its Fourier decomposition, reported in Fig. 1.9, shows a mass ordering, reminiscent of Pb–Pb results [53]. Several authors explain the results in the context of hydrodynamics [54–58].

A similar mass ordering is also found when studying the p_T -spectra for identified particles [59]. The spectra become progressively hard with multiplicity, the higher the mass species. In Pb–Pb collisions, this phenomenon is interpreted as an effect of the collective expansion of the system in a locally thermalized medium. However, other non-QGP effects are able to explain the data: *color reconnection* (see Sec. 2.2.1), which consists of final-state interactions between partons from different hard scatterings, mimics collective phenomena.

Finally, another measurement that shows the ambiguity of p–Pb results is the correlation between the multiplicity of the collision and the average transverse momentum [60]. In Fig. 1.10 $\langle p_T \rangle$ is compared among the three collision systems pp, p–Pb and Pb–Pb. At low multiplicity, $\langle p_T \rangle$ from p–Pb data behaves like the one from pp data, the latter being successfully explained by color reconnection. On the other hand, at high multiplicity the trend is more similar, although with a less pronounced saturation, to Pb–Pb events, where the rescattering with particles of the medium is thought to be responsible for the redistribution of the spectrum.

1.4 Phenomenology of small systems: p–A collisions

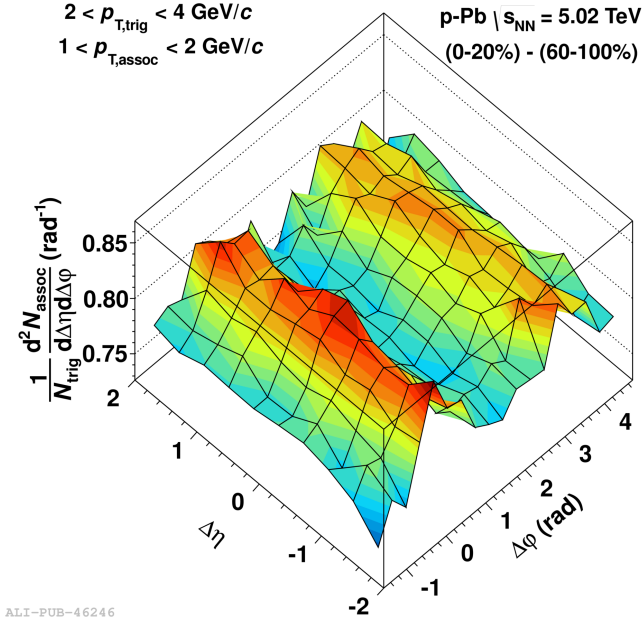


Figure 1.8: Double ridge from two-particle correlations in high multiplicity events after subtraction of two-particle correlations in low multiplicity events. Figure taken from [49].

With this range of puzzling results, it has become clear that further investigation as regards p–Pb collisions, and small systems in general, is needed. One step in this direction is made with the work presented in this thesis, which, while studying the jet-like structure emerging on top of the double ridge, is also able to cast some light on the physics phenomena from which the double ridge itself arises.

1. The physics of hadrons and heavy ions

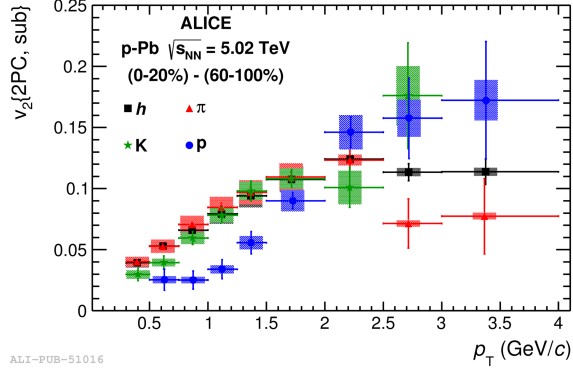


Figure 1.9: Second-order Fourier coefficient v_2 as a function of p_T for pions, kaons and protons. Figure taken from [52].

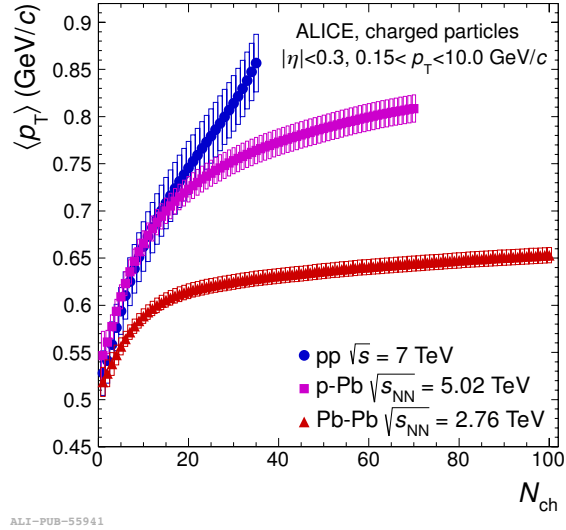


Figure 1.10: Average transverse momentum as a function of charged particle multiplicity for pp, p-Pb and Pb-Pb collisions. Figure taken from [60].

2

The physics of Multiple Parton Interactions

Due to the composite structure of hadrons, a single nucleon-nucleon collision can be pictured as a collision between two groups of partons. In this view, more than one parton-parton interaction can occur in a single hadron-hadron collision. The probability for these *Multiple Parton Interactions (MPI)* to happen increases with center-of-mass energy. In fact, the higher the energy of the collision, the smaller the longitudinal momentum x which can be probed. At smaller x , the parton density¹ is higher, allowing for more MPI to occur [61, 62].

MPI are one of the main contributions to the so-called *Underlying Event (UE)*, which gathers all the activity accompanying the hardest leading order partonic scattering, i.e. additional MPI, initial/final state radiations (ISR, FSR) and interactions between the constituents of the beam remnants.

Multiple parton interactions can arise both at the hard, i.e. high- p_T , and at the semi-hard, i.e. low- p_T , scale². In the former case, they man-

¹Since the parton distribution function of gluons dominates at high energies and small longitudinal momentum x , at LHC colliding partons are essentially gluons.

²There is no sharp cutoff value at which one can distinguish between hard and semi-hard MPI. The work presented in this thesis is performed in the rather low- p_T range

2. The physics of Multiple Parton Interactions

ifest themselves mainly in the form of DPS (Double Parton Scatterings). Those events are characterized by particular topologies, e.g. 4 jets with pair-balanced p_T . Many theoretical and experimental studies are ongoing in this research topic, whose state of the art is summarized in [63]. In this thesis, however, the focus is made on the latter case, where MPI contribute to soft and semi-hard physics structures, in particular *minijets*, which arise from 2→2 scatterings with small momentum transfer and consequently few particles in the final state. Given the inability to describe the very soft regime with perturbation theory, the relevance of experimental measurements containing information on semi-hard MPI is crucial.

In general, UE activity, and in particular MPI, affect the multiplicity of the event. In fact, one of the first signatures of multiple parton-parton interactions emerged from the analysis of charged-particle multiplicity distributions. It was realized that multiple rather than single negative binomial fits are needed to describe those distributions when collision energy becomes high enough [64]. The inclusion of MPI turned out to be one among the models able to justify the need for multiple fits, thus to properly describe the particle production in high-energy hadron collisions. This process from the signature to the formulation of MPI is reviewed in Sec. 2.1. Naturally, exhaustive Monte-Carlo models need to include an implementation of MPI. In this thesis, the specific case of the PYTHIA framework is described in Sec. 2.2, with a qualitative description of its main features and their evolution from the first perturbative-QCD-inspired model.

2.1 Signature of MPI from multiplicity distributions

From Feynman scaling to KNO-scaling violation

In 1969 [65] Feynman pictured a collision between two particles as an exchange of quantum numbers, e.g. isospin. In this process, the (isospin)

between 0.7 and 5 GeV/c, so that semi-hard MPI can be probed, together with soft physics (low- p_T physics without any semi-hard interaction occurring).

2.1 Signature of MPI from multiplicity distributions

current changes sign while passing from one particle (travelling in the positive z direction) to the other ($-z$), causing its source field³ to radiate. As the collision energy increases, this energy flux gets narrower in the z direction, up to the limit of a δ -function. Fourier-transformed in momentum space, this corresponds to a uniform field energy dp_z . The number of particles with energy E , then, amounts to dp_z/E . In statistical terms, the probability to find a particle of type i , with longitudinal and transverse momentum p_z and p_T respectively, is given by

$$f_i(p_T, p_z) \frac{dp_z}{E} d^2 p_T \quad (2.1)$$

Feynman persuaded himself that the structure function $f_i(p_T, p_z)$ is invariant at different collision energies. This is known as *Feynman scaling*.

From this assumption, he concluded⁴ that the mean multiplicity in high-energy collisions depends logarithmically on the center-of-mass energy⁵:

$$\langle N \rangle \propto \ln \sqrt{s}. \quad (2.2)$$

With the same argument, Koba, Nielsen and Olesen formulated in 1972 [64] the so-called *KNO scaling*. It represents an extension of Feynman's $\sqrt{s}$ -invariance assumption to a function $f^{(q)}(p_{T1}, x_1; \dots; p_{Tq}, x_q)$, which describes a correlator among q particles. With the employment of Eq.2.2 in the calculation of the moments of the multiplicity distribution, the latter is found to be of the form

$$P(n) = \frac{1}{\langle n \rangle} \Psi \left(\frac{n}{\langle n \rangle} \right) + \mathcal{O} \left(\frac{1}{\langle n \rangle^2} \right) \quad (2.3)$$

where $\Psi(z := n/\langle n \rangle)$ is an energy-independent function. The normalized KNO-variable z is invariant for η regions and therefore can be used to test the change in dynamics with the pseudorapidity.

³Fields are sources of currents. When currents change, these fields radiate, as it happens in the Bremsstrahlung mechanism.

⁴He didn't give a mathematical proof. A detailed derivation is given in [66].

⁵Since the maximum reachable rapidity increases with $\ln \sqrt{s}$ as well, this translates into the fact that, when the scaling holds, the rapidity density of particles is constant, assuming that the rapidity distribution is flat.

2. The physics of Multiple Parton Interactions

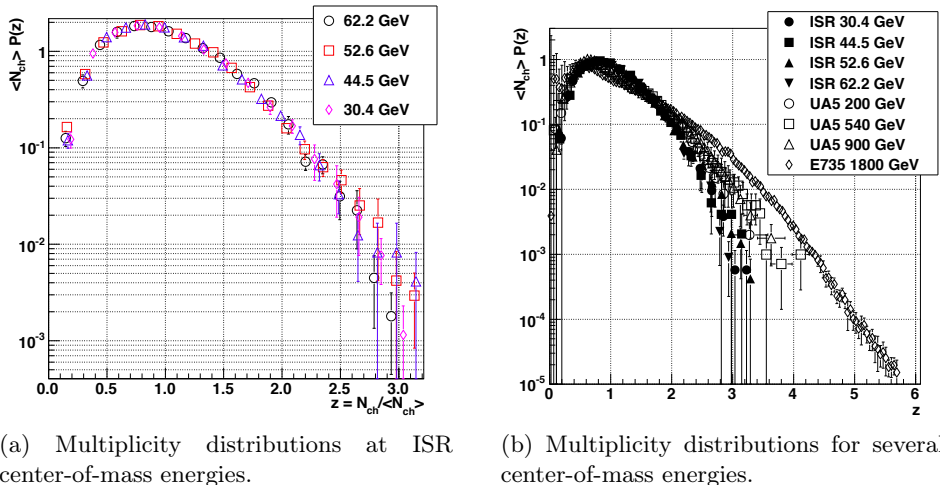


Figure 2.1: $\sqrt{s}$ -dependence of multiplicity distribution of NSD events in full phase space in KNO variables. Figure taken from [66].

To test the KNO scaling at the experimental level, several multiplicity distributions from ISR⁶ at different center-of-mass energies from 30.4 to 62.2 GeV are plotted in Fig. 2.1a, scaled by the mean multiplicity, as a function of the KNO-variable z . They fall nicely on one curve, confirming the KNO (and thus the Feynman) scaling. However, when including distributions by the UA5 experiment [67] from $\sqrt{s} = 200$ to 900 GeV, the KNO scaling does not seem to hold strictly, as shown in Fig. 2.1b. With the distribution at $\sqrt{s} = 1.8$ TeV from the E735 experiment [68], the violation is more pronounced at all z . Although it is evident that the data points from the lowest- and highest-energy sets deviate from each other, no definitive statement can be formulated for the intermediate energies. A study of the moments of the distribution is required, which goes beyond the scope of this thesis. A complete mathematical formulation is available in [66].

⁶The Intersecting Storage Rings (ISR), the very first hadron collider, was operating at CERN between 1971 and 1984. It collided protons on protons, antiprotons and α particles.

2.1 Signature of MPI from multiplicity distributions

A further study has been performed by applying fit functions to multiplicity distributions. In fact, it has been proposed [69] to use Negative Binomial Distributions (NBDs). Their mathematical form is the following:

$$P_{p,k}^{NBD}(n) = \binom{n+k-1}{n} (1-p)^n p^k \quad (2.4)$$

which describes the probability for n failures and $k-1$ successes before the k -th success in a binomial trial with success probability p . NBD distributions follow KNO scaling if k is constant as a function of center-of-mass energy. The physical explanation for the application of the NBD distribution to particle production is understood in the context of the *clan model* [70]. It describes the underlying production by cascades of particles, e.g. by decay or fragmentation. A clan contains all particles that stem from the same *ancestor*. As can be seen in Fig. 2.2a, NBD fits are in good agreement with UA5 multiplicity distributions in full phase space up to $\sqrt{s} = 540$ GeV, although they start to fail in describing the data already from $\sqrt{s} = 900$ GeV. On the other hand, the fit agrees well if using a combination of two NBDs, as shown in Fig. 2.2b.

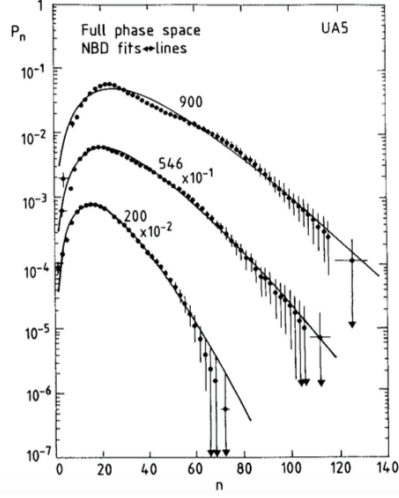
The two-component approach

The first interpretation of the need for a double NBD fit came out of a systematic study performed by Giovannini and Ugoccioni [72]. They assumed two different kinds of events, namely soft and semi-hard, i.e. respectively without and with clusters of particles with transverse energy of at least 5 GeV, as the origins of the two NBD functions. In this context, the multiplicity distribution assumes the form of a weighted sum of NBD functions:

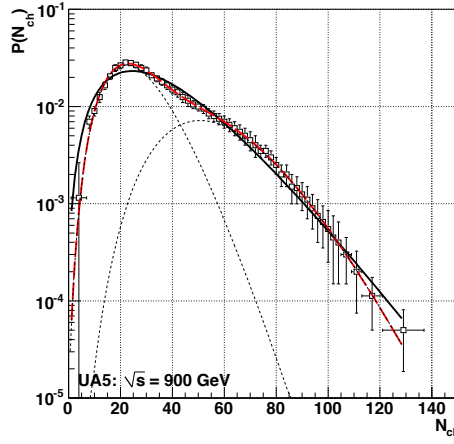
$$P(n) = \alpha_{soft} \cdot P_{\langle n \rangle_{soft}, k_{soft}(n)}^{NBD} + (1 - \alpha_{soft}) \cdot P_{\langle n \rangle_{semi-hard}, k_{semi-hard}(n)}^{NBD} \quad (2.5)$$

The soft component is found to hold the KNO scaling, with the violation being due exclusively to the semi-hard component.

2. The physics of Multiple Parton Interactions



(a) Multiplicity distributions fitted with a single NBD at $\sqrt{s} = 200, 540, 900$ GeV in full phase space. Figure taken from [71].



(b) Multiplicity distribution at $\sqrt{s} = 900$ GeV, fitted with a combination of two NBDs. Figure taken from [66].

Figure 2.2: UA5 multiplicity distributions with fit functions.

2.1 Signature of MPI from multiplicity distributions

Interpretation in the MPI framework

An alternative approach [73, 74] ascribes the origin of the two NBDs to two parton-parton interactions occurring in the same event, not requiring the interplay of more events at the same time.

As first step, the part of the multiplicity distribution that obeys the KNO scaling is subtracted. This is easily done by comparing the distribution, plotted as a function of the KNO variable z , to a KNO fit valid at ISR collision energies. A new variable can be defined, namely $z' = n/\langle n_1 \rangle$, where $\langle n_1 \rangle$ is the average multiplicity of the KNO-scaling part of the total multiplicity distribution⁷.

In the left side of Fig. 2.3, multiplicity distributions with $\sqrt{s}$ from 200 GeV to 1.8 TeV are plotted. Each of them has been previously scaled by $\langle n_1 \rangle$, and they are shown as a function of the KNO-scaling variable $n/\langle n_1 \rangle$. At relatively low multiplicity, the curves fall on top of each other and they agree with the fit of ISR data in KNO variables (solid line). This region is interpreted as being originated from events with single parton-parton collisions only. The fraction which departs from the fit, violating the KNO scaling, shows a peak at around $n/\langle n_1 \rangle = 2$ and is interpreted as the evidence of double (or higher order) parton scattering.

The onset of double parton scattering is estimated to be between $\sqrt{s} = 100$ and 200 GeV, while triple parton scatterings appear at energies slightly below 546 GeV. At the latter energy, though, the triple parton component can be neglected with good approximation and the multiplicity distribution, once subtracted the KNO-scaling part from it, can be used as an estimate of the double parton contribution only. This is further subtracted from higher-energy multiplicity distribution in order to get the triple (or higher order) parton scattering contribution. The decomposition in single, double and higher order parton scatterings contributions for the multiplicity distribution at 1.8 TeV is shown in the right side of Fig. 2.3.

With the data-driven decomposition procedure described above, cross-sections are derived for single (σ_1), double (σ_2) and higher order ($\sigma_3 = \sigma_{\geq 3}$)

⁷ $\langle n_1 \rangle$ rises with $\sqrt{s}$ [74].

2. The physics of Multiple Parton Interactions

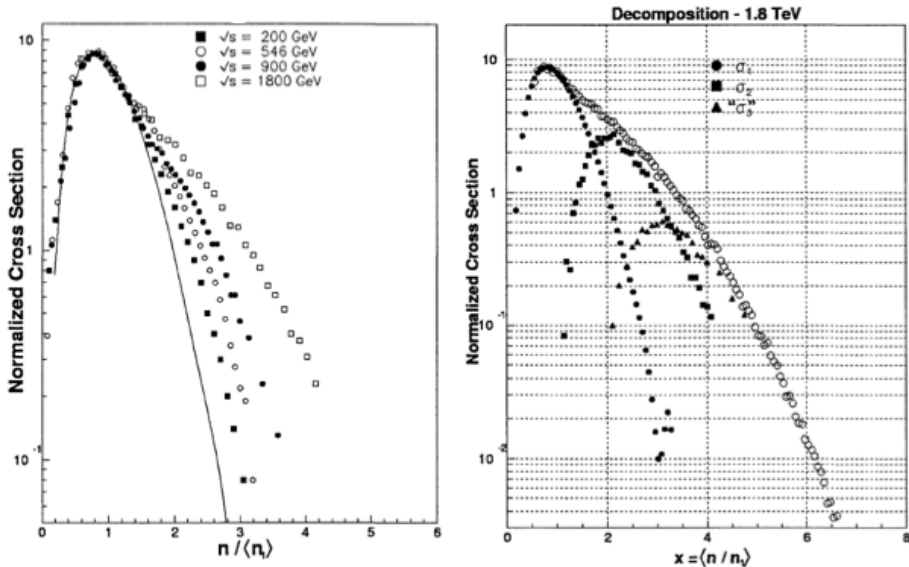


Figure 2.3: Left: Multiplicity distributions scaled by the average multiplicity of the KNO-scaling part of the total multiplicity for $\sqrt{s}$ from 200 GeV to 1.8 TeV. Figure taken from [75]. Right: Multiplicity distribution at $\sqrt{s} = 1.8$ TeV decomposed in three contributions: single, double and triple (or higher order) parton scattering. Figure taken from [76, 77].

parton interactions separately. As shown in Fig. 2.4, in this study the growth of inelastic cross-section with $\sqrt{s}$ has been attributed mainly to multiple (i.e. higher than one) parton interactions, while the single parton contribution does not seem to vary with collision energy.

2.2 MPI in PYTHIA

In the following a basic qualitative description of the MPI framework in Pythia is given. The main ingredients are introduced and their evolution from the first perturbative QCD-inspired model to the modern implemen-

2.2 MPI in PYTHIA

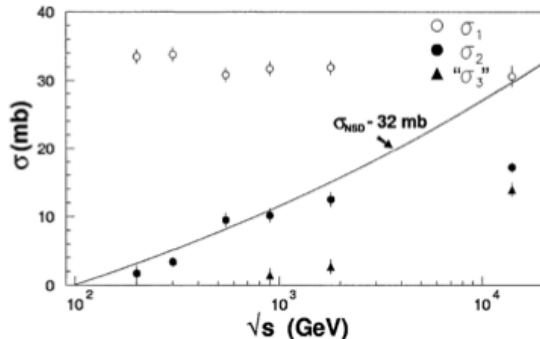


Figure 2.4: Energy dependence of single (empty circles), double (full circles) and higher order (triangles) parton interaction cross-sections. Figure taken from [77].

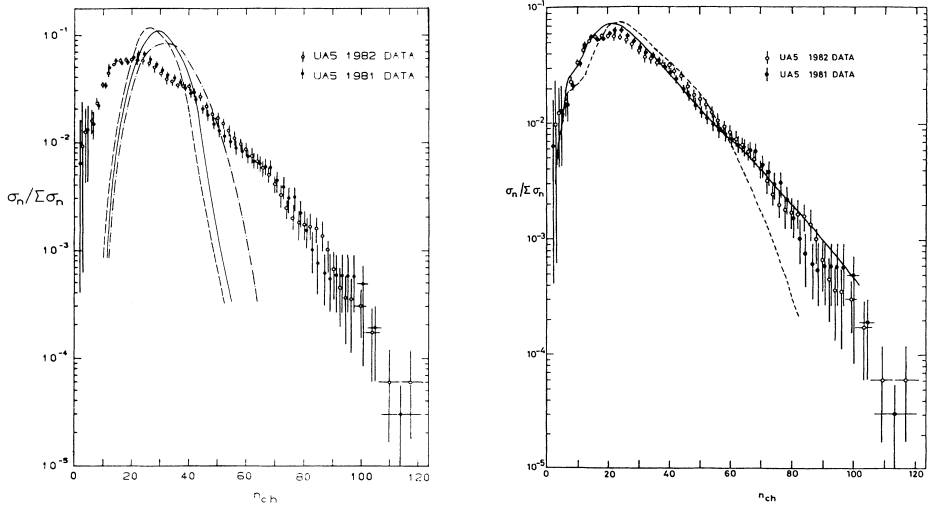
tation in PYTHIA is outlined. The content of this section is a compendium of [78–86], which the reader is referred to for a more exhaustive overview.

2.2.1 From pQCD-inspired model to modern framework

With the models for particle production in hadron-hadron collisions available before 1987, the charged-particle multiplicity distributions at high energies were not even remotely reproduced. Figure 2.5a shows two sets of data points from proton-antiproton collisions at $\sqrt{s} = 540$ GeV with the UA5 experiment, corresponding to 1981 and 1982 data taking periods, together with lines from theoretical models [87]. The dashed line refers to a model where only low- p_T physics is simulated: particles are produced via exchange of a very soft gluon between the two colliding hadrons. Once the hard scatterings are introduced in this model, the distribution becomes the one drawn as a solid line. Neither of them fits the data points, not even after the addition of the initial and final state radiation, as plotted with the dashed-dotted line.

Endorsing the idea of MPI affecting the multiplicity distributions, the first model was built [87] to include this contribution. The model aims at a

2. The physics of Multiple Parton Interactions



(a) Dashed line: low- p_T events only. Solid line: hard interactions included. Dash-dotted line: initial/final state radiation included.

(b) Dashed line: multiple-interaction model with fixed impact parameter. Solid line: multiple-interaction model with double-Gaussian matter distribution.

Figure 2.5: UA5 multiplicity distributions (two sets of data points for 1981 and 1982) with theoretical models (lines). Figure taken from [87].

2.2 MPI in PYTHIA

unified description of high- and low- p_T phenomena within one single framework. Of course, the main difficulty is that while hard MPI are describable in the perturbative regime, semi-hard MPI involve non-perturbative physics. However, by extending perturbation theory to rather low p_T , but still some distance above Λ_{QCD} , most events contain multiple *perturbatively calculable* interactions [79].

The integrated cross section for QCD $2 \rightarrow 2$ processes above some $p_{T\text{min}}$ is divergent in the limit $p_{T\text{min}} \rightarrow 0$:

$$\sigma_{\text{int}}(p_{T\text{min}}) = \int_{p_{T\text{min}}}^{\sqrt{s}/2} \frac{d\sigma}{dp_T} dp_T \propto \frac{1}{p_{T\text{min}}^2} \quad (2.6)$$

and exceeds the total hadron-hadron cross section σ_{tot} below few GeV/c .

This is shown in Fig. 2.6, for 1.8 TeV proton-antiproton collisions at Tevatron and for prediction at 14 TeV proton-proton collisions at LHC. The integrated parton-parton cross sections (falling curves) become greater than the total hadron-hadron cross sections (horizontal curves), respectively at p_T around 2.5 and 5 GeV/c at Tevatron at LHC. This is well above Λ_{QCD} , so that the breakdown of perturbative theory cannot be claimed as the cause.

The first consideration needed to solve the $\sigma_{\text{int}} > \sigma_{\text{tot}}$ paradox is that the interaction cross section is an inclusive observable. Thus, if two parton interactions happen in the same event, those count twice in the σ_{int} , but only once in σ_{tot} . Hence, one identifies the ratio between the two cross sections as the per-event average number of hard interactions above a certain $p_{T\text{min}}$:

$$\langle n \rangle(p_{T\text{min}}) = \frac{\sigma_{\text{int}}(p_{T\text{min}})}{\sigma_{\text{tot}}}. \quad (2.7)$$

With this definition, the divergence shown in Fig. 2.6 is attributed to the number of hard scatterings, rather than to the cross section itself. However, a correction must be applied to avoid that the number of hard scatterings diverges, when $p_T \rightarrow 0$. A first estimate of the effective $p_{T\text{min}}$ cutoff would be $\hbar/d$, from Heisenberg's uncertainty principle, being d the color screening distance. The latter is not known by first principle, though, then an empiri-

2. The physics of Multiple Parton Interactions

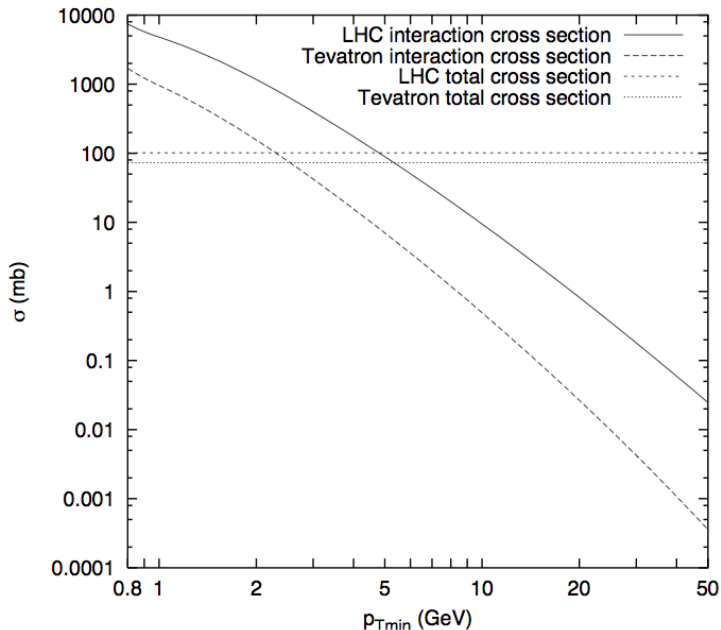


Figure 2.6: Integrated parton-parton cross section as a function of p_T for 1.8 TeV proton-antiproton collisions at Tevatron (dashed falling line) and prediction for 14 TeV proton-proton collisions at LHC (solid line). The horizontal lines represent the total hadron-hadron cross sections (dotted for Tevatron, dashed for LHC). Figure taken from [79].

cal cutoff is applied, at around⁸ 2 GeV/ c [79]. The cutoff value must have an energy dependence: higher energies translate into smaller x probed, where the parton densities are higher, meaning a decrease of the color screening distance.

However, the introduction of p_{Tmin} may still not be sufficient to have a correct physical picture. Indeed, one has to take into account energy-

⁸Approximately the same value has been found when tuning on data the required cutoff for the infrared divergence of the jet cross section, namely the divergence occurring when the energy of gluon emission tends to zero..

2.2 MPI in PYTHIA

momentum conservation. This is not guaranteed while assuming that different parton-parton collisions take place independently of each other, as they can still have $p_T > p_{T\min}$ but use up more momentum than available from the original hadron. This issue is solved in the following way. One arranges the scatterings in falling sequence of p_T values: $\sqrt{s}/2 > p_{T1} > \dots > p_{Tn} > p_{T\min}$. To take into account the $i - 1$ preceding interactions, parton distributions are not evaluated at x_i for the i -th scattered parton from a hadron, but at the rescaled value

$$x'_i = \frac{x_i}{1 - \sum_{j=1}^{i-1} x_j}. \quad (2.8)$$

In this way, the parton distribution function for the i -th scattering becomes zero for x values above the total fraction $\sum_{j=1}^{i-1} x_j$ used by the preceding $i-1$ scatterings. Thus, the number of possible hard scattering is limited not only by the $p_{T\min}$, but also by the energy of the parent hadron.

A further evolution is to consider the impact parameter dependence of the collisions [79]. The smaller the latter, in fact, the higher the probability of having MPI, i.e. the higher the mean value of the probability distribution:

$$\langle \tilde{n}(b) \rangle = k\mathcal{O}(b) \quad (2.9)$$

where $\mathcal{O}(b)$ is the overlap function between the two colliding hadrons, estimated as the convolution of hadronic matter distributions⁹. Such impact parameter dependence is able to explain the so-called *jet-pedestal effect* [82], consisting in a higher UE activity in events with jets with respect to events without jets. In fact, events with hard scales are more likely to happen in central, i.e. small b collisions.

Figure 2.5b shows the same charged-particle multiplicity distributions as in Fig. 2.5a together with theoretical curves for a multiple-parton interaction

⁹Several different matter distributions have been used. Among those are spherical, Gaussian and double Gaussian, the latter also though of as intermediate steps towards a representation of three cores, reflecting the presence of three valence quarks. In more recent implementations [81], the matter distribution is modelled with a momentum dependence, such that the wave-function for small- x partons is spatially broader than for large- x ones.

2. The physics of Multiple Parton Interactions

model ($p_{Tmin} = 2 \text{ GeV}/c$) with fixed impact parameter (dashed line) and with variable impact parameter from overlapping double Gaussian matter distributions (solid line). The data are much better reproduced, validating the role of MPI, with their impact parameter dependence, in high-energy hadronic collisions.

The above-listed features of the early framework are the basis for the modern PYTHIA¹⁰ [78] implementation of MPI. The basic picture has been refined with remarkable improvements [79]. For instance, the scaling of the parton distribution functions according to Eq. 2.8 to account for the energy-momentum conservation is modified by differentiating between valence quarks and sea partons. In addition, MPI is combined with ISR and FSR, i.e. a common p_T -ordered sequence of interactions or branchings is defined, allowing radiations not only for the partons involved in the hardest scattering. The interleaved evolution is described by the *master evolution equation* of interaction probability P in PYTHIA8:

$$\begin{aligned} \frac{dP}{dp_T} = & \left(\frac{dP_{\text{MPI}}}{dp_T} + \sum \frac{dP_{\text{ISR}}}{dp_T} + \sum \frac{dP_{\text{FSR}}}{dp_T} \right) \times \\ & \times \exp \left(- \int_{p_T}^{p_{T^{i-1}}} \left(\frac{dP_{\text{MPI}}}{dp_{T'}} + \sum \frac{dP_{\text{ISR}}}{dp_{T'}} + \sum \frac{dP_{\text{FSR}}}{dp_{T'}} \right) dp_{T'} \right) \end{aligned} \quad (2.10)$$

Here p_T is the transverse momentum of the i -th interaction or branching, with the exponential factor expressing that no further scattering occurs in between the $(i-1)$ -th and i -th interaction or branching. At each lower p_T step, the sums run over an increased number of partons (two more in case of MPI, one more in case of radiation). The meaning of the master equation is a *conditional probability*: given that something has occurred at higher p_T scales, one can estimate the probability of some processes at lower and lower p_T . In a sense, then, it expresses a resolution ordering.

A more recent refinement of the model [80] includes the possibility of rescattering for a parton which has already undergone an interaction. In the simplest case, the subsequent interaction takes place between a parton coming from a previous scattering and a parton from the beam remnant

¹⁰The latest version is PYTHIA8.2.

2.2 MPI in PYTHIA

(single rescattering). In the more complex case, both partons come from previous scattering (double rescattering). The inclusion of rescattering implicates a rescaling of the parton distribution functions, in order to take into account the momentum and flavor changes. The original MPI probability in Eq. 2.10 is generalized to

$$\frac{dP_{\text{MPI}}}{dp_{\text{T}}} \rightarrow \frac{dP_{uu}}{dp_{\text{T}}} + \frac{dP_{u\delta}}{dp_{\text{T}}} + \frac{dP_{\delta u}}{dp_{\text{T}}} + \frac{dP_{\delta\delta}}{dp_{\text{T}}} \quad (2.11)$$

where the uu component stands for the original MPI probability, the $u\delta$ and δu for the single rescattering and the $\delta\delta$ for the double rescattering probabilities. In this way, they are commonly interleaved with the MPI, ISR and FSR processes.

Finally, the introduction of MPI models in hadronic collisions has made the reformulation of color connections among partons necessary. The experimental puzzle has been the rise of $\langle p_{\text{T}} \rangle$ as a function of the number of charged particles n_{ch} , observed for the first time by UA1 [88] and confirmed in many experiments since then, including at LHC [60, 89, 90]. Higher charged-particle multiplicities stem predominantly from events with high MPI activity. If the particle production from each MPI is independent, then the $\langle p_{\text{T}} \rangle$ should not depend on the number of MPI, hence on n_{ch} . On the other hand, if each subsequent MPI affects less and less the final multiplicity, still providing the same p_{T} , the rise of average transverse momentum with multiplicity could be explained. This experimental fact strongly suggests that some form of hadronization phenomenon is at play, connecting partons from different MPI systems. In a naive picture of each MPI system as separate from the others in color space, e.g. each MPI final state of two quarks color-connected to each other only, one ignores that the incoming partons have net non-zero color charge. One should take into account that one (for a quark as initiator) or two (for a gluon) color strings should connect the initiators to the beam remnants. Still, the general picture is somehow incomplete, as especially at high multiplicity, i.e. at high number of MPI, there is non-negligible phase-space overlap between color strings connecting different initiators and the remnants. The straightforward solution is to allow for different MPI systems to be color-connected.

2. The physics of Multiple Parton Interactions

In the original *color reconnection* (*CR*)¹¹ scheme [85], one defines to which MPI system each parton belongs according to the p_T scale. The connection is then made between partons of the lower- p_T system and partons of some higher- p_T one, in such a way as to find the smallest total string length. This procedure is carried on iteratively: once the fate of the lowest- p_T system has been decided, the second lowest is considered and so on. At the end of the connection chain, the actual merging is performed in the opposite direction, i.e. color dipoles are built from the hardest system and lower- p_T partons are progressively inserted, causing the splitting of the original dipole into two new dipoles.

Those dipoles in the above-described model are constructed in the leading-color (LC) approximation: each quark (antiquark) carries a color (anticolor) charge, hence can be linked to a single unique other parton, while gluons carry a color and an anticolor charges, thus forming transverse kinks on strings with quark and antiquark as endpoints. The newest implementation of CR in PYTHIA8 [86] includes subleading color effects. Although the probability of such contributions is dampened by the number of colors¹² as $1/N_C^2$, the suppression is compensated as the number of partons increases, i.e. when MPI occur. The main feature of this model is the introduction of so-called *color junctions*, which allow for more color combinations to happen. Without digging into the formalism, it is worthy just to mention that this aspect allows to reconcile the measured baryon/meson ratio [91], since it increases the probability of forming baryons.

¹¹The commonly used notation CR is a bit misleading. It refers to both *color connections* and *color reconnections*. The former are correction factors used to solve color-space ambiguities, such as the color correlations between different MPI initiators or the subleading color connections. The latter consist in explicit exchange of color during physical processes, such as gluon exchange or interactions of strings. This thesis work, although employing the CR notation in line with the common convention, deals only with the former case.

¹²In a previous model, an attempt was made to exclude the subleading color possibilities by introducing an infinite number of colors ($N_C \rightarrow \infty$). However this envisaged that any color coming into a hard scattering had to be compensated by an anticolor shared by many string endpoints in the remnant.

Part II

EXPERIMENTAL SETUP

3

The ALICE Experiment at the LHC

ALICE is one of the four biggest experiment installed on the Large Hadron Collider (LHC) at CERN. Its name stands for *A Large Ion Collider Experiment*, as it is the only one specifically designed to detect the products of collisions of heavy ions. However, the physics of ALICE spans among many different collision systems: pp and p–Pb collisions are also performed and analyzed, providing a wide pool of interesting results and allowing a rich comparison of observables at different energies and initial conditions.

This Chapter reports details about the LHC and the ALICE experiment. In Sec. 3.1 the collider is described, with a focus on its data-taking conditions during the p–Pb run, being it the topic of this work. Further, in Sec. 3.2, details on the experimental apparatus are given, with a deeper insight into the subdetectors employed for this thesis.

3.1 The Large Hadron Collider

The Large Hadron Collider is the biggest particle collider ever built and it is currently operating, since 2009. It is situated at CERN (Geneva), 175 m underground in its deepest point, and it is 27 km long. The LHC is, in

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fact, only the last ring of a chain of accelerators, which together contribute to let the beams reach the desired energy. The proton and lead beams undergo a slightly different acceleration paths, illustrated in Fig. 3.1. The protons are extracted from a hydrogen bottle with the help of an electric field, then they are accelerated in linear machines and then brought, in sequence, up to 1.4 GeV by the booster, 25 GeV by the Proton Synchrotron (PS), 450 GeV by the Super Proton Synchrotron (SPS). The ion sample is a small lead bar, which is heated to about 500 degrees Celsius to vaporize a small number of atoms, which are ionized with an electrical field in several steps of stripping. The subsequent step is the acceleration in the LEIR (Low-Energy Ion Ring, earlier used as Low-Energy Antiproton Ring and thus known as LEAR), before entering the PS and SPS. The 27-km ring of the LHC is where the final stage of acceleration happens. Once the beams have reached the required energy, they are bent with appropriate configurations of magnetic fields to collide in the four interaction points, where the experiments are placed.

3.1.1 The p-Pb data taking

The p-Pb run was performed in January-February 2013, anticipated by a short pilot run performed in September 2012. The beam energies were 4 TeV for the proton beam and 1.58 TeV per nucleon for the lead beam, resulting in collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The nucleon-nucleon centre-of-mass system moved with respect to the ALICE laboratory system, introduced in the next section, with a rapidity of -0.465, i.e., in the direction of the proton beam. Figure 3.2 shows the evolution of the integrated luminosity during the p-Pb run for the ALICE experiment. At the end of the 20 days of run, the integrated delivered luminosity to ALICE was 32.84 nb^{-1} .

3.2 The ALICE detector

Fig. 3.3 shows a schematic representation of the experimental apparatus, together with the ALICE global coordinate system. Its origin coincides

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The LHC injection complex

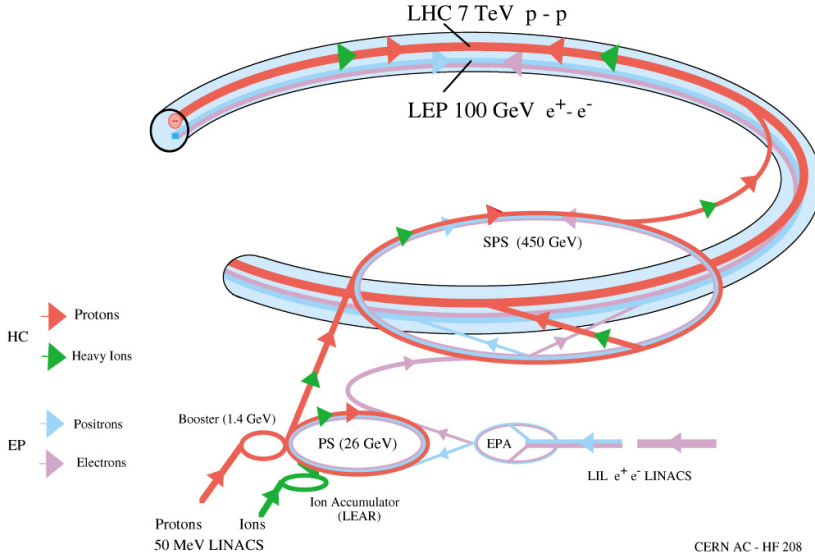


Figure 3.1: The proton (red) and Pb-ion (green) beams accelerating paths at the LHC complex. Figure taken from [92].

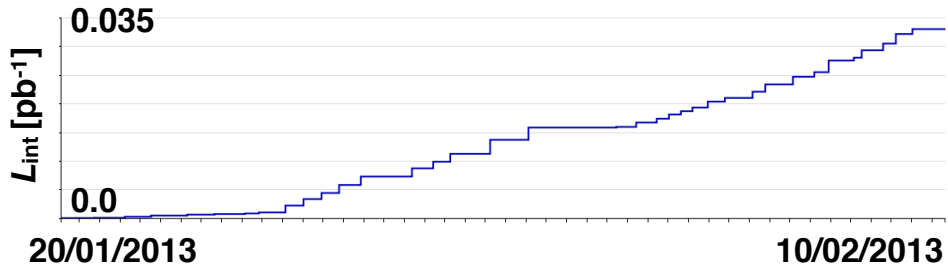


Figure 3.2: Integrated luminosity delivered to the ALICE experiment during the p-Pb run (from 20/01/2013 to 10/02/2013). Figure taken from [93].

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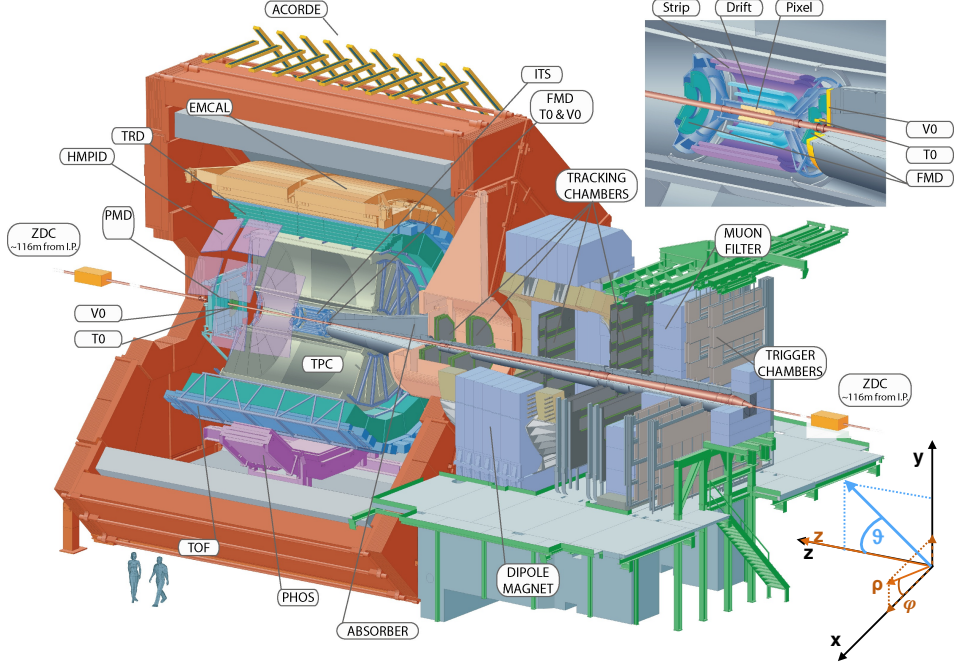


Figure 3.3: The ALICE experiment, with its subdetectors and its global coordinate system (cartesian in black and cylindrical in orange, with the polar angle in light blue).

with the nominal interaction point. The black arrows depict the *cartesian coordinate system*:

- z : the longitudinal coordinate, along the beam direction, points opposite (A-side) to the muon chambers (C-side);
- x and y : the transverse coordinates point respectively toward the center of the LHC and toward the top of the apparatus.

The *cylindrical coordinate system* is sometimes used for convenience, with the same longitudinal coordinate along the beam direction, while $\rho =$

3.2 The ALICE detector

$\sqrt{x^2 + y^2}$ and φ as transverse coordinates. It is drawn in the picture in orange color. In light blue, the polar angle ϑ is also shown. It is the angle between the track and the beam axis. Alternatively to the polar angle, the pseudorapidity is used in high energy physics experiments, which is defined as:

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

so that as the polar angle approaches zero (particles almost in the beam direction) the pseudorapidity tends toward infinity, while as the polar angle approaches 90 degrees (particles transverse to the beam) the pseudorapidity goes to zero. The region where pseudorapidity is around zero is commonly referred to as *midrapidity*.

The ALICE experiment is a complex detector, made up by 13 subdetectors for many different purposes. According to their pseudorapidity coverage, they can be defined as central or forward detectors. A magnet of 0.5 T surrounds the central detectors. Those detectors are usually used for tracking and particle identification measurement, while forward detectors can be used either for trigger or centrality determination or multiplicity/energy measurements, with the exception of muon chambers, which are placed forward and are used for tracking muons, and the PMD used for identifying photons. In the following, only the subdetectors used for this analysis are examined, i.e. the ITS and TPC among the central and the VZERO and ZDC among the forward (backward) ones. For a fully comprehensive description of the detectors, the reader is referred to [94, 95].

3.2.1 ITS

The Inner Tracking System (ITS, see Fig. 3.4) is the closest detector to the interaction point. It is composed of six layers: two of silicon pixel detectors (SPD), two of silicon drift detectors (SDD) and two of silicon strip detectors (SSD), from inward to outward. Those are placed at transverse radii from 3.9 cm to 43 cm and together they cover a pseudorapidity region of $|\eta| < 0.9$.

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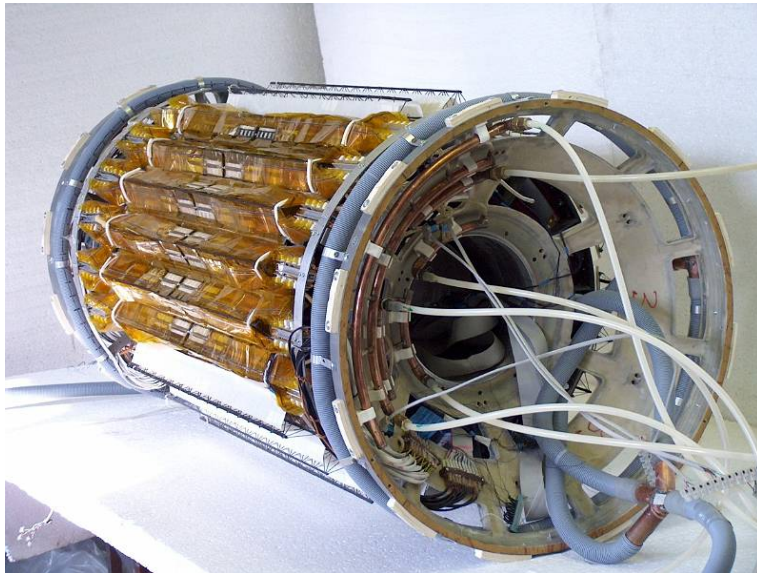


Figure 3.4: ALICE Inner Tracking System. Image: INFN Torino.

Among the tasks of the ITS are the reconstruction of the primary vertex of the collision as well as the reconstruction of secondary vertices of heavy-quark decays (D and B mesons) and hyperons, which it performs with a resolution better than $100\ \mu\text{m}$ in transverse direction. The ITS plays a central role in the tracking, either in combination with other detectors, e.g. improving the momentum resolution of the TPC, or as a stand-alone tracker to reconstruct charged particles that are deflected or decay before reaching the Time Projection Chamber (TPC). The minimum measurable p_{T} at nominal field with the two innermost layers is about $35\ \text{MeV}/c$. The relative momentum resolution achievable with the ITS is better than 2% for pions with transverse momentum between $100\ \text{MeV}/c$ and $3\ \text{GeV}/c$. Moreover, the ITS contributes to the particle identification through the measurement of the specific energy loss (dE/dx). The material budget is low, the total thickness in terms of radiation length X/X_0 is less than 8% [96].

The SPD consist of silicon diodes with a thickness of $200\ \mu\text{m}$. The

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sensor matrix includes 256×160 cells measuring $50 \mu\text{m}$ ($r\varphi$) by $425 \mu\text{m}$ (z) each. The binary read out of the SPD provides information to form the tracks and thus to measure charged-particle multiplicity, but the lack of energy-loss information prevents particle identification. In addition, it can be used as a $L0$ trigger (the lowest of the trigger levels), by reading the logic *OR* signal from chips, indicating that at least one pixel in that chip has produced a signal.

The SDD consist of a $300 \mu\text{m}$ -thick layer of homogeneous high-resistivity silicon. Thanks to the analogic read out, it allows the particle identification along with the tracking measurement.

The SSD consist of sensors equipped on both sides with silicon microstrips. These are arranged under a stereo angle of 35 mrad , allowing for a two-dimensional measurement of the track position, together with an energy-loss estimation for particle identification.

Table 3.1 recaps the spatial precision in transverse ($r\varphi$) and longitudinal (z) direction for the three silicon detectors.

Table 3.1: Spatial precision in transverse ($r\varphi$) and longitudinal (z) direction for the SPD, SDD and SSD of the ALICE ITS.

	SPD	SDD	SSD
Spatial resolution in $r\varphi$ (μm)	12	35	20
Spatial resolution in z (μm)	100	25	830

3.2.2 TPC

The ALICE Time Projection Chamber (TPC, see Fig. 3.5) is the largest in the world, with a volume of 90 m^3 , filled with a gas mixture of $\text{Ne}/\text{CO}_2/\text{N}_2$ ¹. It is located around the ITS, up to a radius of 2.5 m , and it has a length of 5 m . In its center, at $z = 0$, the electrode grid is placed, which produces a drift field of 100 kV towards the multi-wire proportional chambers at $z =$

¹Since the data taking period of Run2, successive to the work of this thesis, a new gas mixture used: Ar/CO_2 [97].

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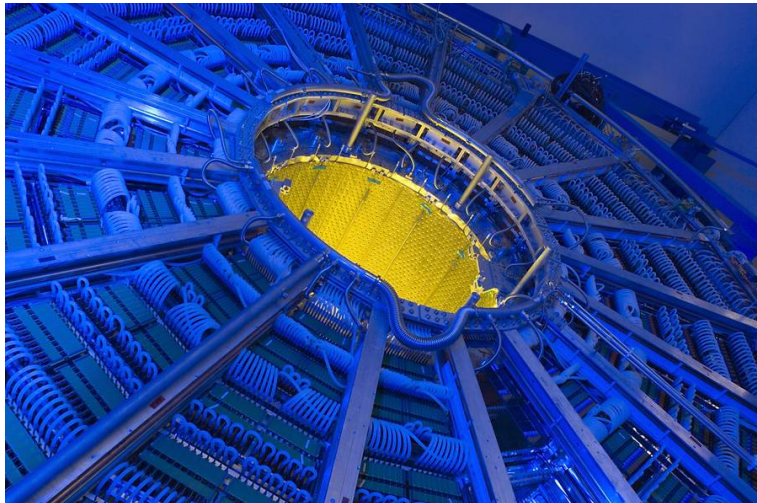


Figure 3.5: ALICE Time Projection Chamber. Image: ALICE Collaboration.

± 2.5 m, from where the signal is read out. Its coverage in pseudorapidity for full-length tracks is $|\eta| < 0.9$. It provides, in combination with the other central barrel detectors, the measurement of charged particles, i.e. their momentum, particle identification, and production vertex. A maximum of 159 clusters can be measured for a track. The position resolution for the inner/outer² radii is $1100/800 \mu\text{m}$ in the transverse plane and $1250/1100 \mu\text{m}$ along the beam axis. Particles with a p_T from about $200 \text{ MeV}/c$ (at nominal field) up to $100 \text{ GeV}/c$ can be measured. For tracks with a momentum below $4 \text{ GeV}/c$, the momentum resolution is better than 2.5%. For tracks with more than 140 clusters, the TPC reaches an energy resolution of 5.5%. The material budget of the ITS and TPC in X/X_0 units is less than 11%.

One of the major challenges with the ALICE TPC is that it has a

²The radial dependence of the track density led to two different types of readout chambers and a radial segmentation of the readout plane. The active area varies radially from 84.1 to 132.1 cm and from 134.6 to 246.6 cm for the inner and outer chamber, respectively.

3.2 The ALICE detector

long drift time (about $90 \mu s$), which means that in that time the detector is busy and cannot collect event. Moreover, at a nominal luminosity of $3 \cdot 10^{30} cm^{-2} s^{-1}$ for pp events, around 30 collisions pile up and have to be disentangled from the triggered event. This is possible by discriminating on the vertices position in the reconstruction phase.

3.2.3 VZERO

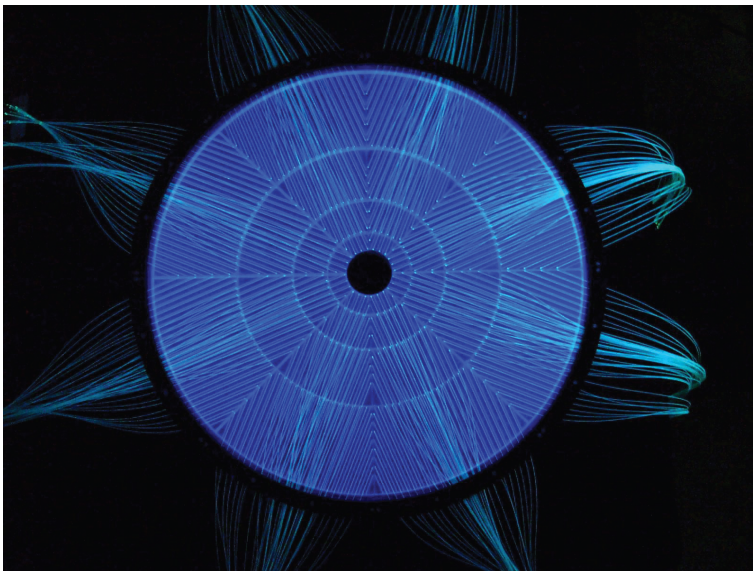


Figure 3.6: VZERO-A detector. Image: ALICE Collaboration.

The VZERO is a forward detector consisting of two arrays of scintillator counters, called VZERO-A and VZERO-C, which are installed on either side of the ALICE interaction point, at $+340$ cm and -90 cm respectively. They cover the pseudorapidity ranges $2.8 < \eta < 5.1$ (A-side) and $-3.7 < \eta < -1.7$ (C-side) and are segmented into 32 individual counters distributed in four rings (see Fig. 3.6). With a time resolution of about 1 ns, beam-gas events that occurred outside of the nominal interaction region can be

3. The ALICE Experiment at the LHC

identified. Moreover, the VZERO detector is used for other purposes: it provides minimum-bias triggers for the central barrel detectors and it can be used as a centrality estimator (see more extensive discussion in Chap.4).

3.2.4 ZDC

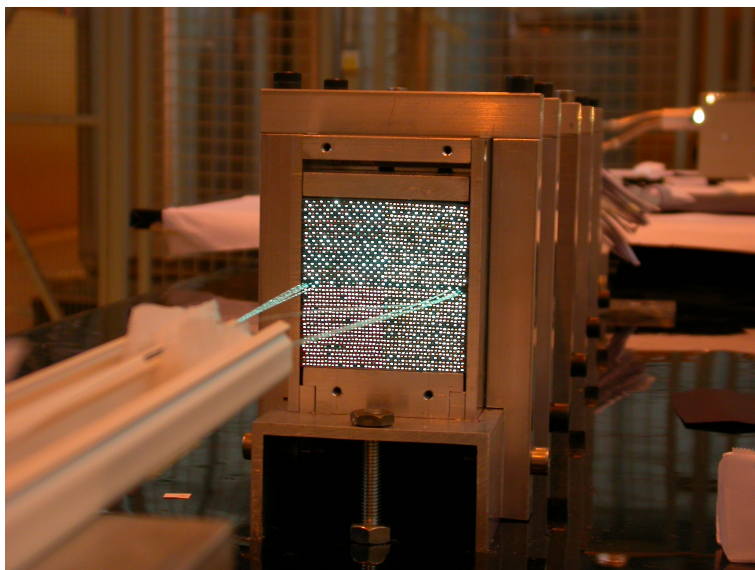


Figure 3.7: Neutron calorimeter of the ZDC detector. Image: INFN Torino.

The Zero Degree Calorimeter (ZDC) consists of two sets of hadronic calorimeters, located at 116 m on either sides of the interaction point, and two electromagnetic calorimeters (ZEM), placed at about 7 m from the IP in the direction opposite to the muon arm (A-side), on both sides of the beam pipe.

The hadronic calorimeters can measure forward neutrons (ZN, with $|\eta| > 8.8$) and protons (ZP, with $6.5 < |\eta| < 7.5$). Protons are spatially separated from neutrons by the magnetic elements of the LHC beam line. Therefore, the ZN is placed between the beam pipes at 0 degrees relative to the LHC

3.2 The ALICE detector

axis, while the ZP is placed externally to the outgoing beam pipe on the side where positive particles are deflected. The hadronic ZDCs are quartz fibres sampling calorimeters: the shower generated by incident particles in a dense absorber (*passive* material) produces Cherenkov radiation in quartz fibres (*active* material) interspersed in the absorber (see Fig. 3.7). Due to the small amount of available space (particularly for the neutron calorimeter, whose transverse size can not be larger than 7 cm), the detectors are very compact. For this reason a very dense tungsten-alloy was used as passive material for ZN, to maximize the containment of the showers.

The energy released by spectators in the zero degree calorimeters is used as an estimate of the centrality (for details see Chap.4).

3. The ALICE Experiment at the LHC

4

Centrality determination in p–Pb collisions

Heavy ions are extended objects and the displacement between their two centers in the transverse plane, namely the *impact parameter*, *quantifies the centrality of the collision*. *Small (large) impact parameters correspond to central (peripheral) collisions*.

Both in Pb–Pb and in p–Pb collisions, the Glauber model is employed to determine geometrical quantities characterizing the event. Its basic features are described in detail in Sec. 4.1. This model derives, from the impact parameter b , the number of participants N_{part} and the number of binary nucleon–nucleon collisions N_{coll} . In particular, one derives the *probability distributions* $\pi_\nu(\nu)$, where ν stands for N_{part} and N_{coll} . These are not directly measurable quantities, though, and they need to be related to measurable observables M via some model. For instance, M could be a multiplicity or an energy distribution. These two cases are revised in Sec. 4.2 and Sec. 4.3, where the V0A charged particle multiplicity distribution and the ZNA energy distribution are used as centrality estimators. The former follows a Negative Binomial Distribution (NBD, Sec. 4.2), while the latter is derived with a Slow Nucleon Model (SNM, Sec. 4.3.1). These models allow one to calculate the conditional probability $P(M|\nu)$ to observe M

4. Centrality determination in p–Pb collisions

for a given ν . From the convolution of the Glauber $\pi_\nu(\nu)$ and the $P(M|\nu)$, one calculates the probability to observe a certain multiplicity or energy:

$$P_{calc}(M) = \sum_\nu P(M|\nu)\pi_\nu(\nu). \quad (4.1)$$

The calculated $P_{calc}(M)$ is compared with the measured $P_{meas}(M)$ and the average ν is extracted for each M -interval, i.e. *event class* or *centrality class*. The details of this procedure are reported for both choices of M in Sec. 4.2.1 and Sec. 4.3.2.

In order to determine unambiguously the geometric parameters, the centrality estimators must vary monotonically as a function of ν . For both Pb–Pb and p–Pb collisions, the *average* of centrality estimators changes monotonically with ν , but fluctuations of M for a fixed ν are present. Those are small for Pb–Pb collisions, but sizable for p–Pb events, compared to the width of the ν distribution¹. For large fluctuations, a centrality classification of the events based on multiplicity may select a sample of nucleon–nucleon collisions which is biased compared to a sample defined by cuts on the impact parameter. A discussion of the biases, i.e. the same value of M may be associated to several ν , is in Sec. 4.2.2. This has been proven to reflect a bias in the fluctuations of the number of multiple parton–parton interactions [98]. The results of this thesis further confirm the bias observation (Sec. 8.5.2).

4.1 The Glauber Monte Carlo

Glauber first presented his model to address the problem of high-energy scattering with composite particles in 1959 [99]. Up to that point, no calculation existed treating the many-body nuclear problem. Two approaches are developed in the Glauber framework, one, the so-called *optical limit*, useful to make the full multiple scattering integral numerically calculable, and the *Glauber Monte Carlo*, exploiting the modern computational powers

¹In Pb–Pb events, the fluctuation around the average value of multiplicity for a fixed N_{part} amounts to 30%. The correspondent value in the case of p–Pb events is about 400%.

4.1 The Glauber Monte Carlo

for a more precise simulation. In spite of the different approaches, the two models find results $(N_{\text{part}}, N_{\text{coll}})$ in good agreement with each other². As regards the fluctuations in N_{part} and N_{coll} , though, with the optical limit approach one is not able to estimate them, while this is possible with the Glauber Monte Carlo.

The basic idea of the Glauber Monte Carlo [99, 100] is that nuclei are composed of uncorrelated nucleons sampled from some density distributions. The colliding nuclei are arranged with a random impact parameter [100]. In this configuration, using the relative distance between two participant-nucleon centroids and the nucleon-nucleon inelastic cross section, the interaction probabilities are computed.

In all calculations of geometric parameters using a Glauber approach, some experimental data must be given as inputs to the model. The two most important are the nuclear charge densities and the energy dependence of the inelastic nucleon-nucleon cross section.

Nuclear charge density The nuclear density for $^{208}_{82}\text{Pb}$ is modelled by a Woods-Saxon distribution for a spherical nucleus

$$\rho(r) = \rho_0 \frac{1}{1 + \exp\left(\frac{r-R}{a}\right)} \quad (4.2)$$

where ρ_0 is the nucleon density in the nucleus center, $R = 6.62 \pm 0.06$ fm the nucleus radius and $a = 0.546 \pm 0.010$ fm the skin depth, inferred from low-energy experimental measurements [101].

Inelastic nucleon-nucleon cross section As low-momentum-transfer processes are involved, the $\sigma_{NN}^{\text{inel}}$ cannot be calculated using perturbative QCD. Thus, this is taken as an experimental input to the model and provides the only beam-energy dependence for the Glauber calculations. At $\sqrt{s} = 5.02$ TeV, the cross section amounts to $\sigma_{NN}^{\text{inel}} = 70 \pm 5$ mb, estimated via interpolation of data collected at $\sqrt{s} = 2.76$ TeV and 7 TeV [102, 103]. Whether the colliding nucleons are protons or neutrons, corrections are negligible. Indeed, the given nucleon-nucleon $\sigma_{NN}^{\text{inel}}$ has been used to estimate

²At least in an inclusive measurement. When dividing the sample up into bins of fractional cross section, a systematic difference in N_{part} is present between the two calculations as the events get more peripheral.

4. Centrality determination in p–Pb collisions

2.1 ± 0.1 b as p–Pb cross section, result which is compatible with the one measured directly via a van der Meer scan in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV[104].

Further step, once given the inputs to the model, is to simulate the collision. For p–Pb collisions, the colliding proton and the Pb nucleus are randomly displaced in the transverse plane. The impact parameter b is selected from the geometrical distribution $dP/db = 2\pi b$ up to a $b_{\text{max}} \simeq 20$ fm $> 2R_{\text{Pb}}$ (b_{max} large enough to simulate collision down to zero interaction probability). The nucleon-nucleus collision is treated as a sequence of independent binary nucleon-nucleon collisions, where the nucleons travel on straight-line trajectories. The collision is assumed to happen if the transverse separation between them is less than the distance corresponding to the inelastic nucleon-nucleon cross section, i.e. $d < \sqrt{\frac{\sigma_{\text{NN}}^{\text{inel}}}{\pi}}$. The cross section is assumed to be independent on the already happened nucleon–nucleon collisions, i.e. the same value is considered for successive collisions, neglecting any deflection or energy loss.

The number of collisions N_{coll} and the number of participants N_{part} are determined by counting, respectively, the binary nucleon collisions and the nucleons that experience at least one collision. In the case of p–Pb events, $N_{\text{coll}} = N_{\text{part}} - 1$.

In the present work, N_{part} and N_{coll} are used, which were extracted from Glauber distributions. Second order corrections to the Glauber model were implemented by Gribov [105], who introduced the event-by-event fluctuations in the configuration of the upcoming proton, conveyed by variations of the nucleon-nucleon cross section. As a result, the Glauber-Gribov N_{part} distributions are much broader than the Glauber N_{part} distribution, while, by construction, the total inelastic p–Pb cross section is unaltered by the proton fluctuations. Both forms (Glauber and Glauber-Gribov) agree very well when compared to the measured multiplicity distribution, so that with this experimental observable it is not possible to discriminate between the two [98].

4.2 Centrality from charged-particle multiplicity distributions

4.2 Centrality from charged-particle multiplicity distributions

As previously stated, the centrality estimator must have a monotonical dependence on the number of participants (or number of binary collisions). Previous experiments [106, 107] have proven that the rapidity-integrated multiplicity of produced charged particles in hadron-nucleus collisions is proportional to N_{part} : $N_{\text{ch}}^{h-A} = N_{\text{ch}}^{pp} N_{\text{part}}/2$. However, experimentally with the detectors at colliders one can measure multiplicity in limited pseudorapidity regions. The pseudo-rapidity densities exhibit an η -dependent scaling, with an approximate N_{part} scaling at $\eta = 0$ and an approximate $N_{\text{part}}^{\text{target}} = N_{\text{part}} - 1$ in the ion-going direction [108, 109]. Similar study by ALICE [110] has confirmed that the primary charged-particle pseudorapidity density at $\eta = 0$ scaled by the mean number of participants is consistent with the corresponding value in pp collisions interpolated to the same center-of-mass energy. Therefore, the linear dependence on N_{part} (N_{coll}) of charged-particle multiplicity in certain pseudorapidity intervals justifies its use as a centrality estimator. The VZERO – A hodoscope detects charged particles produced at forward rapidity in the Pb-going direction, thus it is sensitive to the ion fragmentation. The V0A³ multiplicity distribution is used in the following as centrality estimator.

Multiplicity distributions have been found (see Sec 2.1) to follow a Negative Binomial Distribution (NBD).

The NBD is then sampled $\langle N_{\text{ancestors}} \rangle$ times for each Glauber Monte-Carlo event, with a model for particle production based on $\langle N_{\text{ancestors}} \rangle = f \cdot N_{\text{part}} + (1 - f) \cdot N_{\text{coll}}$, where f is a proportionality factor [111].

4.2.1 NBD-Glauber fit

The Glauber $P(N_{\text{part}})$ probability is convoluted with the conditional probability $P(n|N_{\text{part}})$ from the NBD-fit of the multiplicity distribution per N_{part} ,

³Conventionally, VZERO – A refers to the detector, while V0A to the multiplicity estimator.

4. Centrality determination in p–Pb collisions

to give the $P_{calc}(n)$ through Eq.4.1. Experimentally, the multiplicity distribution $P_{meas}(n)$ is obtained from the signal in the VZERO – A detector, in the direction of the remnants of the Pb nucleus. This distribution is fit with the NBD-Glauber, excluding the region at low (< 10) V0A amplitude. The measured distribution (data points) and the calculated probability (red line) are shown on top of each other in Fig. 4.1. They agree well, considering that the measured distribution is sensitive to experimental parameters such as noise and gain of the detector read-out. This distribution is sliced in percentiles of cross section. A binning of 5% width has been used for the final results of this thesis.

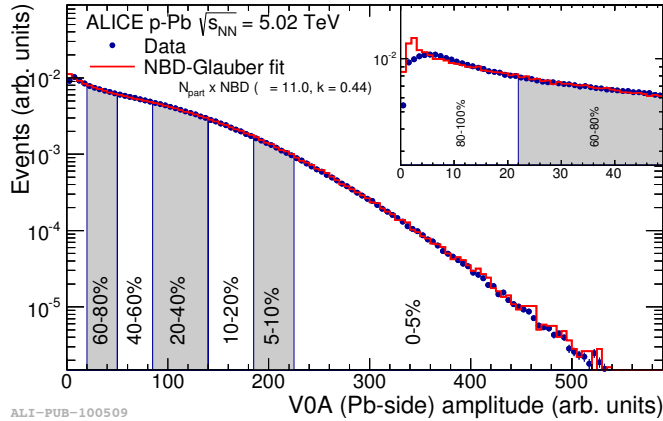


Figure 4.1: V0A multiplicity distribution in arbitrary units for p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV in ALICE. Data points (black dots) are fitted with a NBD+Glauber distribution (red line) with $\mu = 11.0$ and $k = 0.44$ (values from a χ^2 minimization procedure). Figure taken from [98].

4.2.2 Biases on centrality

The method used to select centrality classes according to multiplicity in p–Pb collisions is the same as the one used for Pb–Pb events. However, major differences arise from the two collision systems, one of which is illus-

4.2 Centrality from charged-particle multiplicity distributions

trated in Fig. 4.2. The multiplicity measured by V0A is shown as a function of the number of participants extracted from Glauber, for p–Pb collisions on the left and for Pb–Pb collisions on the right. The first difference concerns the multiplicity and N_{part} ranges spanned: as expected, Pb–Pb collisions give rise to a much more energetic QCD matter, with higher number of participants and final particles. As a consequence, while in Pb–Pb collisions the fluctuations of the two quantities are smaller than the total range, in p–Pb collisions they are of the same size as the multiplicity range. This means that to a fixed N_{part} a large range of multiplicity corresponds, leading to a bias in the selection of events based on multiplicity. This is hereafter referred to as *multiplicity bias*.

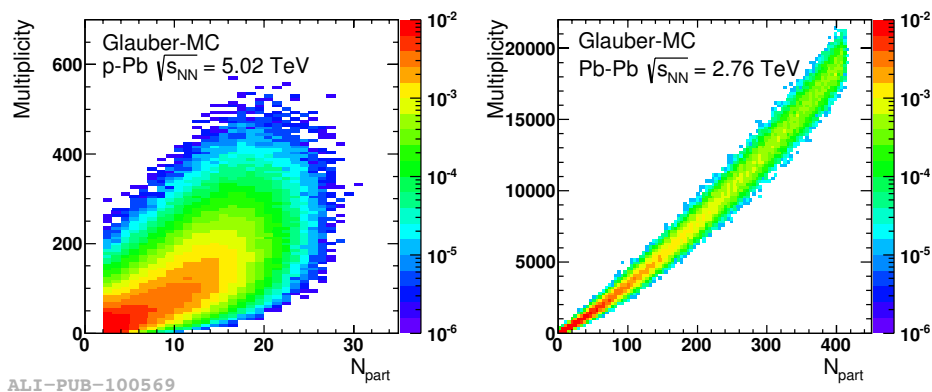


Figure 4.2: V0A multiplicity as a function of N_{part} from Glauber Monte Carlo in p–Pb (left) and Pb–Pb (right) collisions. Figure taken from [98].

To quantify the bias, the Glauber model itself can be employed. In principle, in an unbiased sample, the average multiplicity from the NBD-Glauber distribution (μ) should be equal to the average measured multiplicity per average number of ancestors ($\langle \text{Multiplicity} \rangle / \langle N_{\text{ancestors}} \rangle$). Their ratio is supposed to be 1 in each centrality class. Figure 4.3 plots the ratio, from pure Glauber Monte-Carlo simulation and measured in Pb–Pb and p–Pb collisions, as a function of centrality classes. The Glauber Monte-

4. Centrality determination in p–Pb collisions

Carlo events are classified according to the impact parameter b . For Pb–Pb and p–Pb events, centrality classes are defined with several estimators in different pseudorapidity regions. V0M is the logical AND between V0A and V0C. CL1 is a central-rapidity estimator, requiring a signal in the outer layer of the SPD. More details and numerical values can be found in Sec. 5.5. The ratio is unity overall by definition for the impact-parameter-based selection on Monte Carlo (yellow triangles). A Pb–Pb V0M-based selection (blue triangles) is also plotted, showing that the ratio is 1, except at the most peripheral classes. As concerns the p–Pb estimators-based selection (all other markers in the plot), the ratio differs from unity, becoming greater than 1 for central events and lower than 1 for peripheral events. The flatness of the ratio from the Monte Carlo is expected, as no multiplicity bias is introduced when selecting classes according to the impact parameter. In Pb–Pb collisions, where the multiplicity fluctuations are small, the ratio deviates from unity only for the most peripheral collisions. Contrary, in p–Pb, the most central (peripheral) collisions have on average much higher (lower) multiplicity per participant. The amount of fluctuation is also dependent on the estimator. It has been shown that the estimators with the largest bias correspond to those with the largest width of the NBD distribution [98].

To interpret this result, one can consider the correspondence between the number of produced particles (multiplicity) and the number of *ancestors*, explained in the clan model. Hence, the multiplicity fluctuation corresponds to a fluctuation in the number of ancestors, i.e., in recent Monte-Carlo generators, in the number of MPI. As an example, results for the PYTHIA event generator coupled to the p–Pb Glauber calculation are presented in the following. PYTHIA is used for generating N_{coll} independent pp collisions. Figure 4.4 shows the centrality dependence of the ratio of $\langle n_{\text{hard}} \rangle$, which is the number of semi-hard parton–parton scatterings (see Eq. 2.7), and $\langle N_{\text{coll}} \rangle$. The trend is similar to Fig. 4.3, suggesting that the multiplicity bias is an expression of the bias on the number of MPI.

Another kind of bias exists for events with hard scatterings and arises from the different kinematic ranges spanned in p–Pb and pp collisions, as explained in the following. In Fig. 4.5, the multiplicity distribution mea-

4.2 Centrality from charged-particle multiplicity distributions

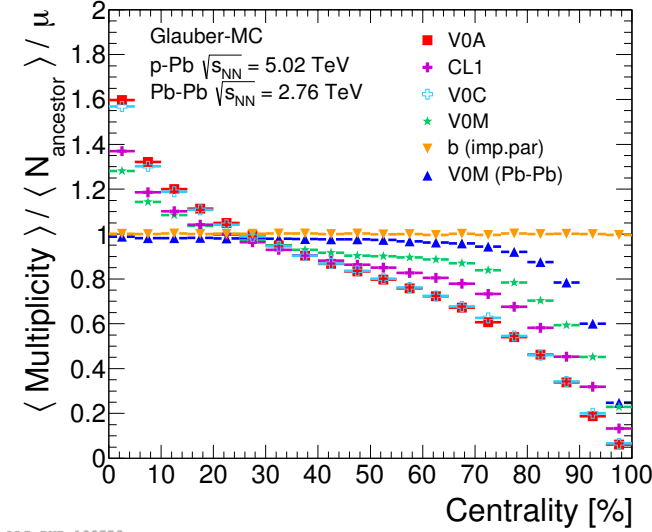


Figure 4.3: Ratio between the measured multiplicity per ancestor ($\langle \text{Multiplicity} \rangle / \langle N_{\text{ancestors}} \rangle$, with $\langle N_{\text{ancestors}} \rangle = \langle N_{\text{part}} \rangle$) and the NBD-Glauber mean multiplicity μ as a function of centrality classes, for Pb-Pb (blue triangles) and p-Pb (all other markers). The yellow triangles are obtained with an impact-parameter-based selection on Monte-Carlo events. All the other markers are relative to several estimators-based selection, as described in the legend. Figure taken from [98].

sured with the midrapidity estimator CL1 is compared between p-Pb (blue) and pp (red) events. The area on the left of the vertical dotted lines corresponds to fractions (80%–100% and 60%–100%) of the p-Pb distribution, which in turn corresponds to fractions (80% and 97%, respectively) of the pp cross-section. The high-multiplicity tail of the pp distribution is likely to arise from the fragmentation of partons that have been generated in a hard scattering with a large momentum transfer. This tail falls out of the 80-100% p-Pb multiplicity class. Therefore, the peripheral p-Pb event class suffers from a bias towards the bulk of soft collisions. This is known as *jet-*

4. Centrality determination in p-Pb collisions

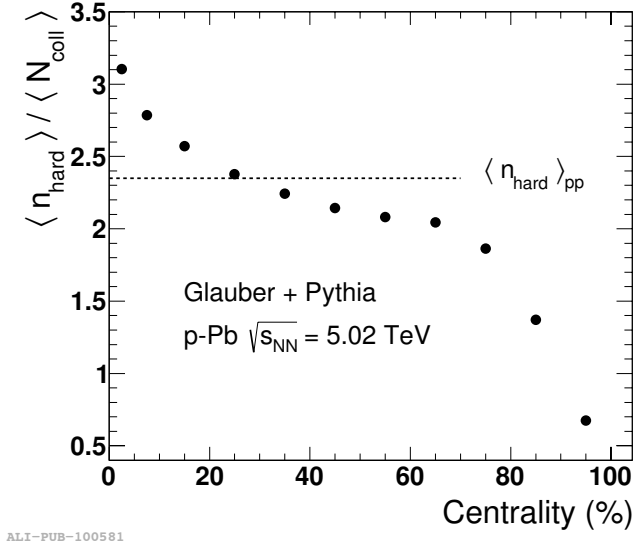


Figure 4.4: Centrality dependence of the ratio between the average number of semihard scatterings and the average number of binary nucleon-nucleon collisions for a p-Pb system created with N_{coll} independent pp collisions with the Glauber+PYTHIA model. The dotted line represents the offset of the average number of semihard scatterings in pp collisions. Figure taken from [98].

veto bias. In other words, what happens is that events containing jets are shifted towards more central classes.

Finally, another bias involves the peripheral events. The MPI probability is dependent on the nucleon-nucleon impact parameter ($\langle b_{\text{NN}} \rangle$). The latter, as shown in Fig. 4.6 for a HIJING simulation, rises significantly for $N_{\text{part}} < 6$. As a consequence, the average number of MPI for peripheral events is reduced, enhancing the effect of the multiplicity bias for these events. Since the impact parameter is a purely geometrical feature, the so-called *geometric bias* is independent on the centrality estimator.

As a summary, three different biases are present:

4.3 Centrality from zero-degree energy

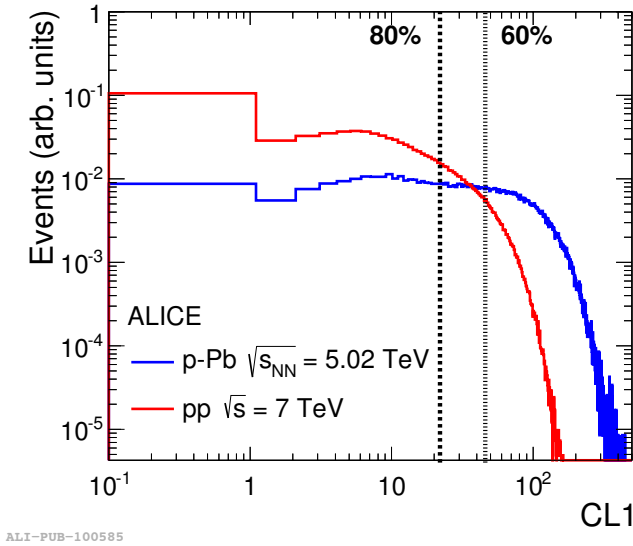


Figure 4.5: Midrapidity multiplicity distribution for pp (red) and p-Pb (blue). Dotted lines delimit the 80% and 60% lowest classes of p-Pb events. Figure taken from [98].

- **multiplicity bias:** affects in particular central and peripheral events and is estimator-dependent;
- **jet-veto bias:** affects peripheral and mid-central events, involves only midrapidity estimators and is p_T -dependent;
- **geometric bias:** affects peripheral events and is estimator-independent.

4.3 Centrality from zero-degree energy

The zero-degree calorimeters (ZDCs) detect the *slow* nucleons produced in the collisions. It has been shown [112] that their multiplicity, hence the energy deposit in the calorimeter, is monotonically related to N_{coll} and can

4. Centrality determination in p-Pb collisions

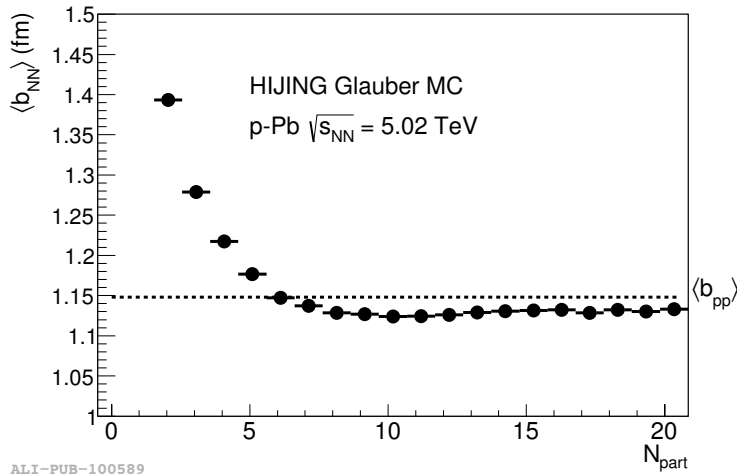


Figure 4.6: Average nucleon-nucleon impact parameter as a function of N_{part} from a HIJING-Glauber Monte Carlo p-Pb simulation. The dotted line represents the mean impact parameter in pp collisions. Figure taken from [98].

then be used as centrality estimator. In fact, experiments at lower energies show that the features of the emitted nucleons, i.e. angular momentum and multiplicity distributions, depend only weakly on the projectile energy. One can conclude that the emission of slow particles is mainly governed by nuclear geometry.

4.3.1 Slow Nucleon Model

Emitted nucleons are classified as *black* or *gray*, according to their velocity. Black particles, with $\beta \lesssim 0.25$ in the nucleus rest frame, are produced by nuclear evaporation processes. Gray particles, with $0.25 \lesssim \beta \lesssim 0.7$, are the knocked-out nucleons. The relative fraction of black and gray neutrons is evaluated as $\langle N_{black\ n} \rangle / \langle N_{gray\ n} \rangle = 0.9$, while for protons $\langle N_{black\ p} \rangle =$

4.3 Centrality from zero-degree energy

$0.65\langle N_{\text{gray p}} \rangle$ holds [113]. Since there are no available models to describe slow nucleon emission at LHC energies, a heuristic approach has been developed, known as Slow Nucleon Model (SNM) [98], relying on the weak dependence of slow nucleon properties on the collision energy.

The average number of slow protons is found to be linearly dependent on N_{coll} , in either case of black and gray ones:

$$\langle N_{\text{slow p}} \rangle \propto N_{\text{coll}} \quad (4.3)$$

The average number of slow neutrons is estimated through the following formula:

$$\langle N_{\text{slow n}} \rangle = \alpha(\gamma \langle N_{\text{slow p}} \rangle) + \left(a - \frac{b}{c + \gamma \langle N_{\text{slow p}} \rangle} \right). \quad (4.4)$$

where the proportionality factor $\gamma = 1.71$ is obtained through a minimization procedure [114]. The first term, then, is proportional to N_{coll} and it describes the emission of the gray neutrons. Black neutrons, instead, are expected to saturate and this is described by the second term.

The number of emitted nucleons results trivially proportional to the energy detected in the beam direction.

4.3.2 SNM-Glauber fit

The Glauber $P(N_{\text{coll}})$ probability is convoluted with the conditional probability $P(e|N_{\text{coll}})$ where e is the ZN energy distribution from the SNM, to give the $P_{\text{calc}}(e)$ through Eq.4.1. The experimental distribution $P_{\text{meas}}(e)$ is given by the neutron energy spectrum measured in the Pb-remnant-side (ZNA) of the calorimeter. Figure 4.7 shows the experimental energy distribution (data points) fitted with the SNM-Glauber (red line). The zoom-in shows that at low energy the peaks for single neutron emission (from 1 to 4) appear. The main features of the measured distribution are very well described by the SNM, in spite of it being a heuristic model. The discrepancies occur at high energy, where the pile-up in the data becomes non negligible, and at low energy (< 4 neutrons), since the model assumes that

4. Centrality determination in p–Pb collisions

for $N_{\text{coll}} = 0$ no neutron impacts on the detector apart from the effect of nuclear evaporation, while in fact there is a contribution due to participant neutrons and EM emission.

Further, the spectrum is sliced in percentiles of the hadronic cross section to define the event classes.

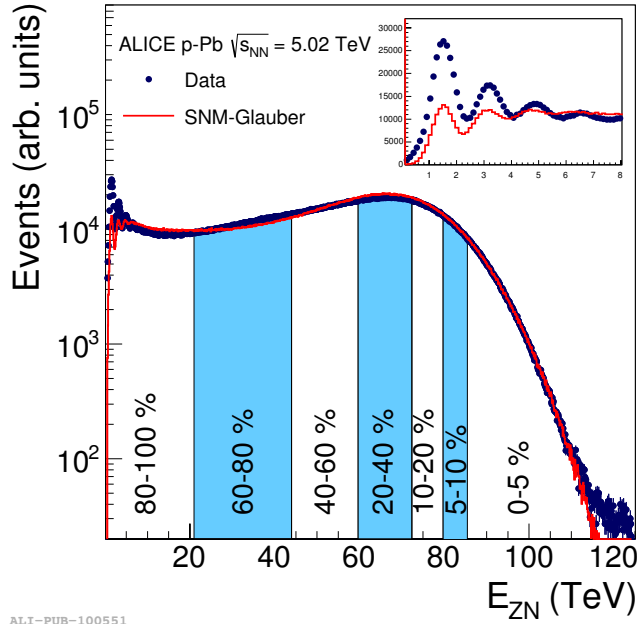


Figure 4.7: Slow nucleon energy spectrum measured by ZNA (Pb-going direction) in p–Pb collisions. Data points (blue markers) are fitted with an SNM-Glauber distribution (red line). Figure taken from [98].

The ZNA-based selection does not suffer from any bias, apart from the geometrical bias, which is estimator-independent. However, the N_{coll} determination by the SNM-Glauber model is model-dependent. To avoid this problem, while keeping the lack of bias of the ZNA estimator, the so-called *hybrid method* has been developed. It consists of selecting centrality classes

4.4 $\langle N_{\text{coll}} \rangle$

based on the energy deposit in the ZNA and in each of those classes determining the N_{coll} according to specific assumptions on the scaling of particle multiplicity in regions of phase space different from the one covered by the ZNA.

Three sets of $\langle N_{\text{coll}} \rangle$ are estimated, based on the following assumptions:

- (a) $N_{\text{coll}}^{\text{mult}}$: the charged-particle multiplicity at midrapidity is proportional to the number of participants N_{part} :

$$\langle N_{\text{part}} \rangle_i^{\text{mult}} = \langle N_{\text{part}} \rangle_{\text{MB}} \left(\frac{\langle dN/d\eta \rangle_i}{\langle dN/d\eta \rangle_{\text{MB}}} \right)_{-1 < \eta < 0}$$

$$\langle N_{\text{coll}} \rangle_i^{\text{mult}} = \langle N_{\text{part}} \rangle_i^{\text{mult}} - 1$$
- (b) $N_{\text{coll}}^{\text{high-}p_{\text{T}}}$: the yield of high- p_{T} particles at midrapidity is proportional to the number of binary collisions N_{coll} :

$$\langle N_{\text{coll}} \rangle_i^{\text{high-}p_{\text{T}}} = \langle N_{\text{coll}} \rangle_{\text{MB}} \frac{\langle S \rangle_i}{\langle S \rangle_{\text{MB}}}, \text{ where } S \text{ is the charged-particle yield with } 10 < p_{\text{T}} < 20 \text{ GeV}/c$$
- (c) $N_{\text{coll}}^{\text{Pb-side}}$: the charged-particle multiplicity in the Pb-going direction is proportional to the number of participants belonging to the target nucleons $N_{\text{part}}^{\text{target}} = N_{\text{part}} - 1 = N_{\text{coll}}$:

$$\langle N_{\text{coll}} \rangle_i^{\text{Pb-side}} = \langle N_{\text{coll}} \rangle_{\text{MB}} \frac{\langle S \rangle_i}{\langle S \rangle_{\text{MB}}}, \text{ where } S \text{ is the raw signal in the innermost ring of the VZERO - A detector.}$$

The three sets of values do not differ more than 9% in all centrality classes. Being the latter (c) estimated in the Pb -going direction, thus more sensitive to the fragmentation of the ions, it is used in the final results of this work (see Sec. 8.5.2).

4.4 $\langle N_{\text{coll}} \rangle$

In this section, a summary of the $\langle N_{\text{coll}} \rangle$ extracted in 5%-wide centrality bins and further used in this analysis is collected in Tab.4.1. The two sets of values represent:

- $\langle N_{\text{coll}} \rangle^{V0A}$ extracted from the NBD-Glauber method in V0A centrality classes

4. Centrality determination in p–Pb collisions

- $\langle N_{\text{coll}} \rangle^{\text{Pb-side}}$ extracted from the hybrid model with inner V0A-ring multiplicity scaling in ZNA centrality classes

Table 4.1: Average values of N_{coll} extracted with the NBD-Glauber for V0A centrality classes ($\langle N_{\text{coll}} \rangle^{V0A}$) and with the hybrid method with V0A multiplicities $\langle N_{\text{coll}} \rangle^{\text{Pb-side}}$ in ZNA centrality classes.

Centrality classes	$\langle N_{\text{coll}} \rangle^{V0A}$	$\langle N_{\text{coll}} \rangle^{\text{Pb-side}}$
0–5%	15.60	13.28
5–10%	13.60	12.25
10–15%	12.50	11.65
15–20%	11.50	11.09
20–25%	10.70	10.52
25–30%	9.89	9.92
30–35%	9.13	9.29
35–40%	8.32	8.59
40–45%	7.55	7.86
45–50%	6.72	7.10
50–55%	5.91	6.33
55–60%	5.04	5.60
60–65%	4.20	4.91
65–70%	3.58	4.27
70–75%	3.06	3.66
75–80%	2.57	3.13
80–85%	2.15	2.66
85–90%	1.82	2.22
90–95%	1.52	1.86
95–100%	1.25	1.48

Part III

ANALYSIS

5

Data, event and track selection

This chapter contains technical details concerning the data treatment for the present thesis. Sections 5.1 and 5.2 collect information about datasets and Monte Carlo productions. Data recorded by ALICE undergo a selection based on trigger and vertex requirements. Those are described in Sec. 5.3. Once passed the selection, the tracks belonging to these events are requested to fulfill some experimental conditions, examined in Sec. 5.4.

As this thesis aims at studying the multiplicity dependence of two-particle correlations, events are divided in multiplicity bins¹. Several detectors are used as estimators. All the details are given in Sec. 5.5.

5.1 ALICE p–Pb data

ALICE has recorded more than 120 million minimum bias proton-lead collision events in January-February 2013. Events with good quality included in the ALICE data-taking period LHC13c amount to around 100 million minimum bias. The corresponding runs are listed below:

¹Apart from this chapter, the term *centrality* is replaced by *multiplicity* throughout this thesis. That is because, unlike in ion collisions, it is hard to give a meaningful connotation to the *centrality of a hadron-ion collision*.

5. Data, event and track selection

- 195677, 195675, 195673, 195644, 195635, 195633, 195596, 195593, 195592, 195568, 195567, 195566, 195531, 195529.

Raw data are reconstructed and calibrated in as many iterations as needed until no more signs of miscalibration are found in the properties of events and tracks. Two iterations (pass2) have been performed for the above-mentioned dataset.

Reconstructed data have been analyzed with the *AliROOT* software [115] and with extensive use of the LEGO train framework [116].

5.2 Monte Carlo p–Pb simulations

Monte Carlo simulations are used to estimate the detector performance and to correct the data. Moreover, for a validation of the correction procedure, a Monte-Carlo closure test is performed.

In order to reproduce the same detector configuration as in real data runs, a production 'anchored' to ALICE data-taking LHC13c period is employed, namely the LHC13b2_efix_p1. This production is based on the DPMJET version 3.05 event generator [117], with particle transport through the detector using GEANT3 version 3.21 [118]. DPMJET contains the features of soft multiparticle production according to the Dual Parton Model (DPM) [119]. It has been shown that DPM provides a good phenomenological description of soft processes, although it fails to reproduce some effects, e.g. the double ridge [49].

As a systematic check, the HIJING model [120] has been used as event generator. This Monte Carlo describes hard and semihard parton scatterings, with special attention to the minijets. Negligible differences are found in the final observables when using HIJING with respect to DPMJET for the corrections.

5.3 Event selection

The event selection is based on two main steps. First, the trigger selection defines the kind of desired events, as described in Section 5.3.1. A sub-

5.3 Event selection

sequent selection is made according to the properties of the reconstructed vertex, examined more deeply in Section 5.3.2. Further selections are applied to the remaining events. For instance, pile-up events as described in 7.1 are removed.

Figure 5.1 shows the number of events after each selection step: no selection (0), trigger selection (1), vertex selection (2), pile-up rejection (3). As shown in step 3, about 80 million events pass all the selection criteria and are used in this analysis.

5.3.1 Trigger selection

The trigger selection employed in this analysis is `AliVEvent::kINT7`. It is a minimum bias trigger, which means that it requires the lowest selection criteria to all inelastic collisions. It consists of two stages: online and offline.

The online trigger is based on a coincidence of at least one beam-beam hit in V0A (out of 32 total) and at least one beam-beam hit in V0C (out of 32 total). Beam-beam events are discriminated from beam-gas events with a cut on the arrival timing of the signal. The trigger selection accepts 99.2% for non-single-diffractive (NSD²) events and has a negligible contamination from single-diffractive (SD) and electromagnetic (EM) events [121].

The offline trigger consists of a measurement of 'average' times in V0A and V0C, by averaging corresponding cell-by-cell timing. These values may be different from the times of hits measured online, since the former are measured with time slewing corrections depending on the charge in each cell. Then, a cut is applied on these average times, to select only events compatible with beam-beam timing. A further cut is applied on timing of neutrons hitting the ZDC on Pb-going direction (namely the ZNA). This time should

²Three types of diffractive events exist:

1. Single Diffractive (SD): $hh \rightarrow h + gap + X$, with *gap* denoting an empty rapidity region, and *X* anything that is not the original beam particle;
2. Double Diffractive (DD): $hh \rightarrow X + gap + X$, i.e. both hadrons fragment;
3. Non Diffractive (ND): inelastic events without diffractive dissociation.

The Non Single Diffractive (NSD) category comprises 2. and 3.

5. Data, event and track selection

be compatible with neutrons originated from beam-beam collisions. V0 and ZDC offline selection rejects about 0.5-1.0% of events triggered online.

The number of events which pass the trigger selection on the data sample analyzed in this analysis is shown in step 1 in Fig. 5.1.

5.3.2 Vertex selection

Among the triggered events, only those events that have at least one reconstructed vertex are analyzed.

The vertex is reconstructed in two phases. The first estimate of its position is made using only the information from the SPD space points³. Then, once all the tracks have been fully reconstructed as described in the following section, the vertex is fitted again using the information from the whole set of reconstructed tracks, together with the average position and spread of the beam-beam interaction region estimated for the particular run [95]. The vertex reconstruction algorithm is fully efficient for events with at least one reconstructed primary charged particle in the common TPC and ITS acceptance ($|\eta| < 0.9$). Furthermore, the longitudinal position of the vertex must fulfill the $|z_{vtx}| < 10$ cm condition with respect to the nominal interaction point at $z_{vtx} = 0$.

The number of events which pass the vertex selection on the data sample analyzed is shown in step 2 in Fig. 5.1.

5.4 Track selection

In this work, only primary charged particles are analyzed. Those are either prompt particles from the primary vertex or their decay products, except those from weak decay of strange particles. Particle tracks are reconstructed via the Kalman filter method [95] with the two main ALICE tracking detectors: the ITS and the TPC. In particular, the default cuts for each reconstruction step are presented in the following.

³Centers of gravity of charge clusters deposited in the sensitive elements of the detectors (pad rows in the TPC, silicon sensors in the ITS).

5.4 Track selection

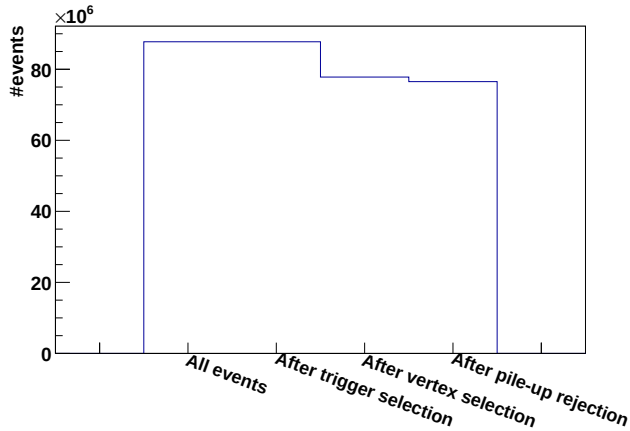


Figure 5.1: Number of events at each selection step.

- The first phase of the track reconstruction involves the determination of the space points in the TPC, starting from the clusters in the two outer rows inward to the inner ones. The following cuts are applied:
 - minimum 20 hit pad-rows in the TPC out of a maximum of 159
 - χ^2 per point of the momentum fit in the TPC less than 4
- The reconstruction continues with the prolongation of the TPC track into the ITS, making use of the hits in the silicon subdetectors. One requires:
 - at least 2 matching hits in the ITS out of 6, one of those being necessarily one in the first three layers
 - χ^2 per point of the momentum fit in the ITS less than 36
- Then, the distance of closest approach (DCA) between the reconstructed TPC-ITS track and the vertex has to fulfill some requirements:
 - DCA in longitudinal direction D_z smaller than 2 cm

5. Data, event and track selection

- p_T -dependent cut on the DCA in transverse direction: $D_{xy} < 0.018 \text{ cm} + 0.035 \text{ cm } p_T \text{ (GeV)}^{-1}$

- Finally, tracks are propagated outwards.

In the presented analysis, kinematic cuts are applied to the fully reconstructed tracks:

- pseudorapidity within combined ITS-TPC acceptance $|\eta| < 0.9$
- cut on transverse momentum $p_T > 0.2 \text{ GeV}/c$.

In the AliROOT framework, sets of predefined cuts are called *filterbits* and are identified with a number: the one above described is called filterbit 96. This set is optimized to have a strong rejection of tracks from secondary particles. In particular, this is achieved through the requirement of tight DCA cuts, at the expense of a less uniform acceptance in azimuthal angle.

As comparison, looser cuts are applied which result in a quite uniform acceptance in azimuthal angle. On the other hand, the rejection of tracks from secondary particles is less efficient. Those cuts, also referred to as filterbit 768 (or *hybrid tracks*), are summarized below:

- TPC
 - p_T -dependent minimum number of TPC clusters: $N = 70 + 1.5 \cdot p_T \text{ (GeV}/c)$ for $p_T < 20 \text{ GeV}/c$, $N = 100$ for $p_T > 20 \text{ GeV}/c$
 - χ^2 per point of the momentum fit in the TPC less than 4
 - fraction of shared TPC clusters less than 0.4
- ITS
 - χ^2 per point of the momentum fit in the ITS less than 36
- DCA
 - $D_z < 3.2 \text{ cm}$
 - $D_{x,y} < 2.4 \text{ cm}$.

5.4 Track selection

Another available set is filterbit 32: this is slightly different from filterbit 96, in that it requires at least one ITS cluster being necessarily in the SPD, rather than in one of the first three layers.

Figure 5.2a and 5.2b show respectively the φ and η single particle distribution for three above-mentioned sets of track cuts. As expected, the number of particles is larger with the looser DCA cuts, i.e. filterbit 768 (green curve), since more secondaries and primaries with larger DCA are included. On the other hand, with the two tighter DCA cuts, i.e. filterbit 96 (red) and 32 (blue), one finds more pronounced azimuthal inefficiencies. However, the latter are reduced with the inclusion of the third layer of the ITS in the track cuts. This is quantifiable in the difference between filterbits 96 and 32.

5.4.1 Efficiency and contamination corrections

The reconstruction efficiency ϵ is defined as:

$$\epsilon(p_T, \eta, z_{vtx}, \text{mult}) = \frac{N_{\text{reconstructed}}(p_T, \eta, z_{vtx}, \text{mult})}{N_{\text{generated}}(p_T, \eta, z_{vtx}, \text{mult})} \quad (5.1)$$

where $N_{\text{generated}}$ and $N_{\text{reconstructed}}$ are the number of particles, respectively at the generator level and after the reconstruction, extracted from the Monte Carlo. The contamination correction, instead, removes secondary tracks, i.e. originating from weak decays or from interactions with the detector materials. The contamination factor is defined as:

$$\lambda(p_T, \eta, z_{vtx}, \text{mult}) = \frac{N_{\text{secondaries}}(p_T, \eta, z_{vtx}, \text{mult})}{N_{\text{total}}(p_T, \eta, z_{vtx}, \text{mult})} \quad (5.2)$$

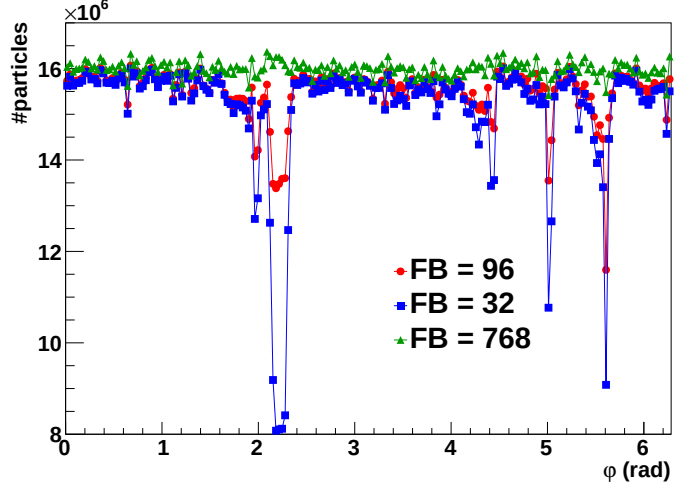
where N_{total} is the sum of primaries and secondaries.

The two corrections are applied in the form of a weighting factor, which is calculated from:

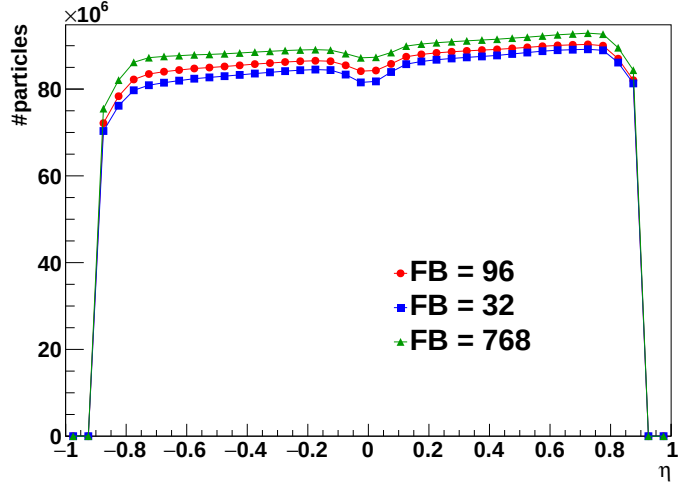
$$F(p_T, \eta, z_{vtx}, \text{mult}) = \frac{1 - \lambda(p_T, \eta, z_{vtx}, \text{mult})}{\epsilon(p_T, \eta, z_{vtx}, \text{mult})} \quad (5.3)$$

These corrections are found to be independent of multiplicity, therefore in the correction procedure the factor is multiplicity-integrated. The weighting

5. Data, event and track selection



(a) φ distribution.



(b) η distribution.

Figure 5.2: Single particle distributions for several filterbits: 96 (default for this analysis, red), 32 (blue), 768 (green).

5.5 Classification of events in multiplicity classes

factor is stored in a THn ROOT object, as it is dependent on p_T , η and z_{vtx} . The correction is applied 'on-the-fly', i.e. while constructing the single particle distributions and the two-particle correlation functions.

Figure 5.3 shows on the left the efficiency and on the right the contamination as a function of η (top panels) and p_T (bottom panels). At large $|\eta|$, i.e. greater than 0.9, the efficiency drops and the contamination gets higher. This is expected, since those values are outside the TPC acceptance. As a function of the transverse momentum, in the p_T range 0.2-1 GeV/ c the efficiency is 68–80% while the contamination decreases from 4 to 1%, while for $p_T > 1$ GeV/ c the efficiency is around 80% and the contamination becomes less than 1%. The z_{vtx} -dependence reflects the η -dependence of the efficiency: it drops at large values of the longitudinal coordinate of the vertex.

5.5 Classification of events in multiplicity classes

This thesis aims at examining the multiplicity dependence of the two-particle correlations. As already explained in Chapter 4, unlike in Pb–Pb, the concept of centrality for events from p–Pb collisions is not easily defined. The experimentally driven procedure consists in measuring the multiplicity of produced particles and classifying events into 'classes'. For this analysis, classes are defined according to the charge deposition in the VZERO-A detector. The A-side is preferred, being it the direction of flight of the Pb beam, thus more sensitive to the fragmentation of Pb ions. The events are classified in 5% percentile ranges of the multiplicity distribution, denoted as '0–5%' to '95–100%' from the highest to the lowest multiplicity.

Other detectors are employed to study the variation of the final observables with the η -gap between the tracking detectors and the detectors used for multiplicity estimation. For instance, the CL1 estimator requires a signal in the outer layer of the SPD ($|\eta| < 1.4$), while the ZNA uses the neutron calorimeter of the ZDC in the A-side ($|\eta| > 8.8$). The estimation for the 95–100% class with ZNA is not reliable, due to the limited efficiency of the detector in case of very low multiplicity.

5. Data, event and track selection

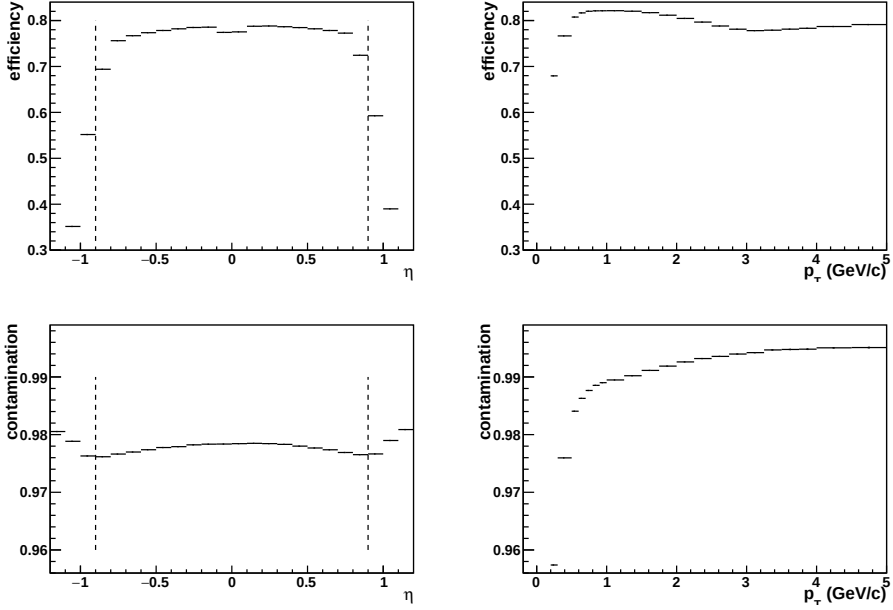


Figure 5.3: Top row: efficiency vs η (left) and p_T (right). Bottom row: contamination vs η (left) and p_T (right). η -projections are integrated in $0.2 < p_T < 5.0$ GeV/ c . p_T -projections are integrated in $|\eta| < 0.9$. The dashed lines delimit the η -region in which the present analysis is performed.

Table 5.1 shows the corrected average values of multiplicity at midrapidity ($|\eta| < 0.9$, with $p_T > 0.2$ GeV/ c) for each class corresponding to the different multiplicity estimators. These are denoted as $\langle N_{\text{ch}} \rangle^{\text{estimator}}$. According to the estimator, different dynamic ranges of multiplicities are spanned, from the larger one with CL1, followed by the one with V0A, to the smaller one with ZNA. These values are also shown in Fig. 5.4 as a function of the multiplicity percentile for V0A, CL1 and ZNA.

5.5 Classification of events in multiplicity classes

Table 5.1: Corrected per-class average charged particle multiplicity at midrapidity ($|\eta| < 0.9$, with $p_T > 0.2$ GeV/ c) with V0A, CL1, ZNA estimators.

Class	V0A	CL1	ZNA
0–5%	68.8	77.3	45.7
5–10%	54.9	59.3	43.2
10–15%	48.4	51.2	41.5
15–20%	43.6	45.4	39.9
20–25%	39.8	40.7	38.3
25–30%	36.4	36.8	36.5
30–35%	33.3	33.2	34.6
35–40%	30.5	29.9	32.5
40–45%	27.7	26.9	30.2
45–50%	25.1	24.2	27.9
50–55%	22.7	21.4	25.6
55–60%	20.4	18.9	23.1
60–65%	18.1	16.5	20.9
65–70%	15.8	14.1	18.8
70–75%	13.6	11.8	16.7
75–80%	11.4	9.5	14.8
80–85%	9.4	7.3	13.1
85–90%	7.6	5.5	11.4
90–95%	6.0	3.7	10.1
95–100%	4.5	2.3	—

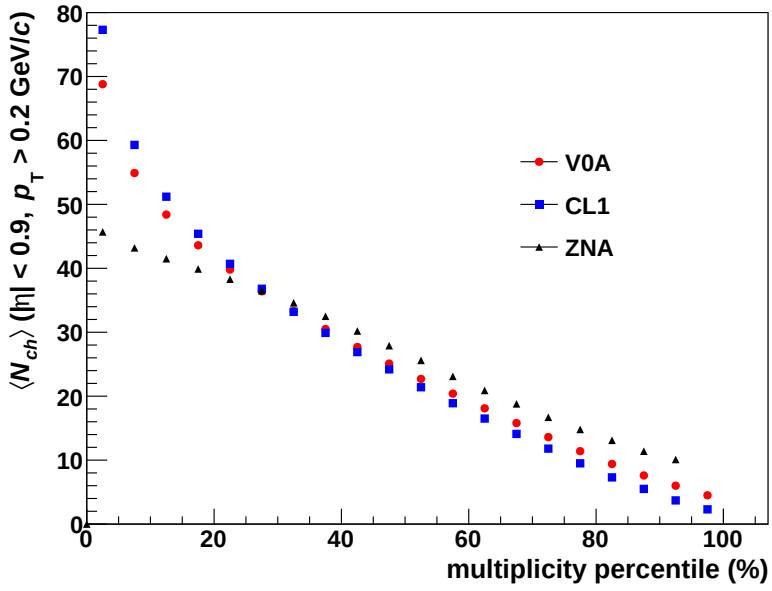


Figure 5.4: Average charged particle multiplicity $\langle N_{ch} \rangle$ as a function of multiplicity percentile for V0A (red circles), CL1 (blue squares) and ZNA (black triangles).

6

Analysis Procedure

The two-particle correlations method is a statistical procedure based on pairing particles produced in the collisions and on studying them as a function of their spatial separation. According to the latter, one may disentangle the contributions from different kinds of phenomena. For instance, two particles with a small separation in azimuthal angle and pseudorapidity have a high probability of belonging to a minijet, i.e. stemming from the same $2 \rightarrow 2$ scattering. But, if their separation in pseudorapidity is large, collective-like effects could be invoked as common source of the two particles. This work investigates the multiplicity dependence of the minijets, emerging from two-particle correlations once removed the long pseudorapidity-range component, in p-Pb events. Furthermore, this method provides an estimate of the number of semi-hard scatterings from which minijets originate. These sources turn out to be an alternative way to characterize the underlying MPI activity.

In this chapter, the analysis method is described. How to extract the two particle correlation functions and their different structures are explained in Sec. 6.1. The subsequent corrections, i.e. mixed event, efficiency and contamination from secondaries, are further described in Sec. 6.1.1. In Sec. 6.2, the subtraction of correlations extraneous to minijets is explained and the $\Delta\eta$ value chosen as a limit between short- and long-range is justified.

Ultimately, the final observables of this study are defined mathematically and presented as extracted experimentally in Sec. 6.3, with particular emphasis on the link between sources of semi-hard scatterings and minijets in PYTHIA.

6.1 Two-particle correlations

The two-particle correlation method consists of relating pairs of particles produced in an event, according to their properties. Both particles are chosen according to a requirement for their transverse momentum: the so-called 'trigger' particles must fulfill $p_T > p_{T\text{cut},\text{trig}}$, while for the 'associated' particles $p_T > p_{T\text{cut},\text{assoc}}$ holds. In each event, one has to consider all possible combinations of pairs. It follows that all trigger particles act as associated ones as well. Being in general $p_{T\text{cut},\text{trig}} \geq p_{T\text{cut},\text{assoc}}$, all associated particles whose momentum is lower than $p_{T\text{cut},\text{trig}}$ ($p_{T\text{cut},\text{assoc}} < p_T < p_{T\text{cut},\text{trig}}$) cannot play the role of triggers and are denoted $N_{\text{pure assoc}}$ in the following. Given a number of trigger particles N_{trig} in one event, the number of possible combinations is:

$$N_{\text{pairs}} = \frac{N_{\text{trig}} \cdot (N_{\text{trig}} - 1)}{2} + N_{\text{trig}} \cdot N_{\text{pure assoc}}. \quad (6.1)$$

The $1/2$ factor accounts for double counting.

This combinatorics is performed for each event and then summed over all the events in a given multiplicity class. In order for the pair yield, i.e. number of pairs, to be compared among different multiplicity classes, it is normalized to the number of trigger particles per each multiplicity class.

The per-trigger pair yield is studied as a function of the difference in azimuthal angle $\Delta\varphi = \varphi_{\text{trig}} - \varphi_{\text{assoc}}$ and in pseudorapidity $\Delta\eta = \eta_{\text{trig}} - \eta_{\text{assoc}}$ between the two particles. Since tracks are accepted in the combined ITS-TPC acceptance ($|\eta| < 0.9$), the possible values lie within $|\Delta\eta| < 1.8$. The 2π -wide interval in $\Delta\varphi$ is represented shifted in $[-\pi/2, 3\pi/2]$.

The region $-\pi/2 < \Delta\varphi < \pi/2$ is known as 'near side', while $\pi/2 < \Delta\varphi < 3\pi/2$ is known as 'away side'.

6.1 Two-particle correlations

The per-trigger pair yield for particles from the same event is then defined as:

$$S(\Delta\eta, \Delta\varphi) = \frac{1}{N_{trig}} \frac{d^2 N_{assoc, same}}{d\Delta\eta d\Delta\varphi}. \quad (6.2)$$

Pairs are also required to have a separation of $|\Delta\varphi_{min}^*| > 0.02$ rad and $|\Delta\eta| > 0.02$, where $\Delta\varphi_{min}^*$ is the minimal azimuthal distance at the same radius between the two tracks within the active detector volume after accounting for the bending in the magnetic field. This cut removes the problem of track splitting (when one single track is reconstructed as two close tracks) and track merging (when two tracks are so close that cannot be resolved). Then the correction to the yield is made via the mixed event technique, explained in the next paragraph.

Furthermore, correlations induced by secondary particles from neutral-particle decays are suppressed by cutting on the invariant mass (m_{inv}) of the particle pair. In this way, pairs are removed which are likely to stem from a γ -conversion ($m_{inv} < 0.04$ GeV/ c^2), a K^0 decay ($|m_{inv} - m(K^0)| < 0.02$ GeV/ c^2) or a Λ decay ($|m_{inv} - m(\Lambda)| < 0.02$ GeV/ c^2). The corresponding masses of the decay particles (electron, pion, or pion/proton) are assumed in the m_{inv} calculation.

6.1.1 Event mixing

The event mixing is a common technique in many analyses. In particular, for correlation studies, it is useful to remove effects of unwanted correlations, e.g. due to limited acceptance or dead areas in the detectors. These two cases are explained with simple toy models. Finally, it is shown how the mixed event correction is applied to the two-dimensional two-particle correlation distributions.

Limited η acceptance

As illustration, one can assume for simplicity that both trigger and associated particles have a flat distribution, limited in pseudorapidity. A template distribution is shown in the left panel of Fig. 6.1. The $\Delta\eta$ distribution is

obtained from the convolution of two of such kinds of theta functions. Due to the fact that the η distributions are not infinitely extended, the $\Delta\eta$ distribution of the pairs acquires a triangle shape centered in 0. This is drawn in the right panel of the Fig. 6.1.

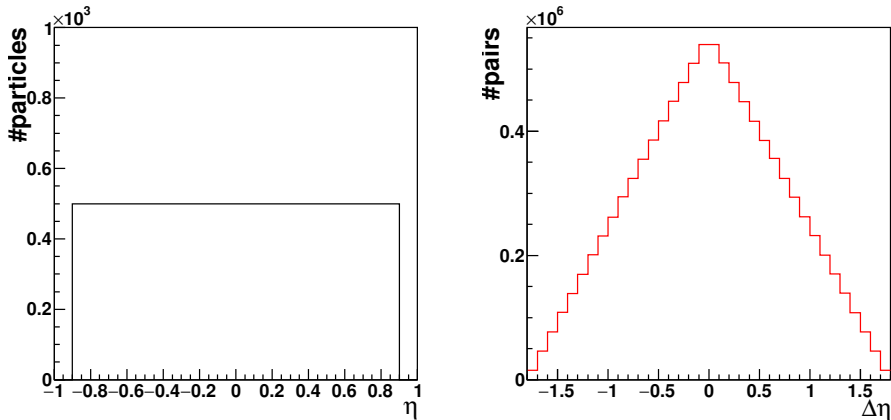


Figure 6.1: Left: flat single particle pseudorapidity distribution in $|\eta| < 0.9$. Right: Pair $\Delta\eta$ distribution from flat single particle η distribution.

Holes in φ distributions

As an example, one can assume that both trigger and associated particles have one hole in an otherwise flat distribution in φ . An example is drawn in the top left panel of Fig. 6.2. This could be due to a dead area in the detector, or any other effect such as a less efficient read-out in the zone, which would affect in the same way both trigger and associated particle distributions. Therefore, since the difference in φ of the positions of the holes in the single (trigger and associated) particle distributions is centered around 0, the effect of the hole is to enhance the near-side peak, as it is shown in the top right panel of Fig. 6.2.

Another peculiar case is shown in the bottom of Fig. 6.2, where two

6.1 Two-particle correlations

holes exist (left panel) whose positions are separated by π , which translate in two peaks mimicking the near- and the away-side peaks (right panel). The correction procedure for this effect is discussed in the next paragraph.

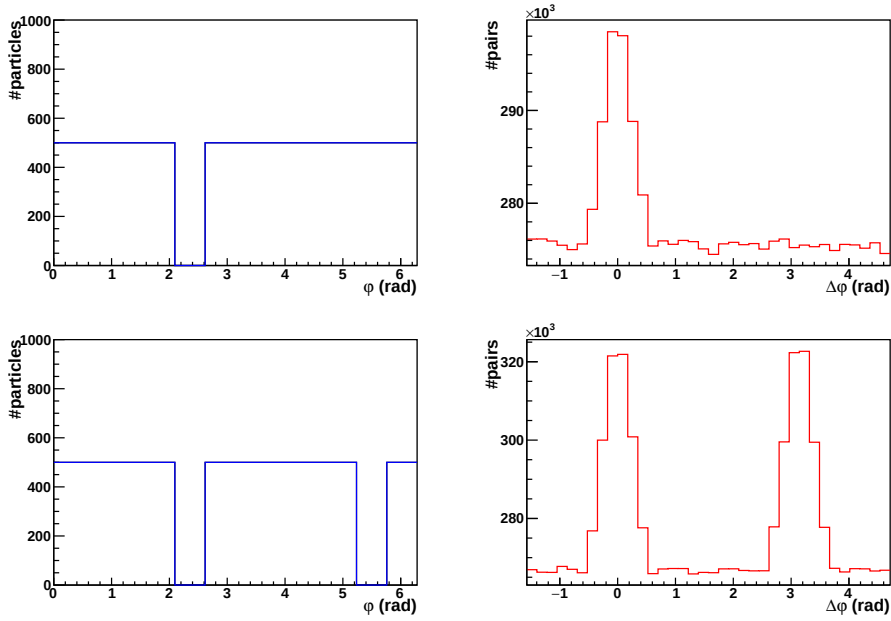


Figure 6.2: Top: flat single particle azimuthal distribution with one hole in φ (left) and pair $\Delta\varphi$ distribution from flat single particle distributions with one hole in φ (right). Bottom: flat single particle azimuthal distribution with two holes in φ (left) and pair $\Delta\varphi$ distribution from flat single particle distributions with two holes in φ (right).

Mixed event correction to two-particle correlations

To mix the events, first one selects the trigger particle from a certain event 'N'. According to its multiplicity and vertex position¹, one looks at the pool of the previous N - 1 events which belong to the same class and z_{vtx} bin. The correlation pairs are then constructed, according to the p_T values as in the same event, by associating one particle from the previous N - 1 events to the original trigger particle. On average, from 5 to 20 events are mixed. Once the correlation is performed, the event 'N' is added to the event pool, available to be mixed with the N+1-th event.

The per-trigger pair yield for particles from the mixed event is then defined as:

$$M(\Delta\eta, \Delta\varphi) = \frac{d^2 N_{assoc, mixed}}{d\Delta\eta d\Delta\varphi} \quad (6.3)$$

and is shown in Fig. 6.3 for $0.7 < p_{T,assoc} < p_{T,trig} < 5$ GeV/c for the highest (0–5%) V0A multiplicity class.

Before applying the mixed event correction to the same event correlation function, a normalization factor is introduced such that no correction factors are applied in case trigger and associated particle have the same direction of motion. In other words, $M(0,0)$ must be equal to 1. That is why a normalization factor α is multiplied to the $M(\Delta\eta, \Delta\varphi)$ distribution:

$$B(\Delta\eta, \Delta\varphi) = \alpha M(\Delta\eta, \Delta\varphi). \quad (6.4)$$

The two-particle correlation distribution, corrected for the mixed event distribution, thus reads as:

$$\frac{S(\Delta\eta, \Delta\varphi)}{B(\Delta\eta, \Delta\varphi)} \quad (6.5)$$

Especially in the lowest multiplicity classes, where the number of particles is lower, statistical effects can perturb the two-particle correlation

¹The longitudinal position of the vertex z_{vtx} slightly affects the acceptance window in pseudorapidity. This is more relevant for the SPD, the innermost tracking detector, than for the TPC. Nevertheless, to avoid any bias, events are mixed only within a 2 cm-wide z_{vtx} range. Moreover, as this analysis is performed as a function of multiplicity, events are mixed within the same multiplicity class.

6.1 Two-particle correlations

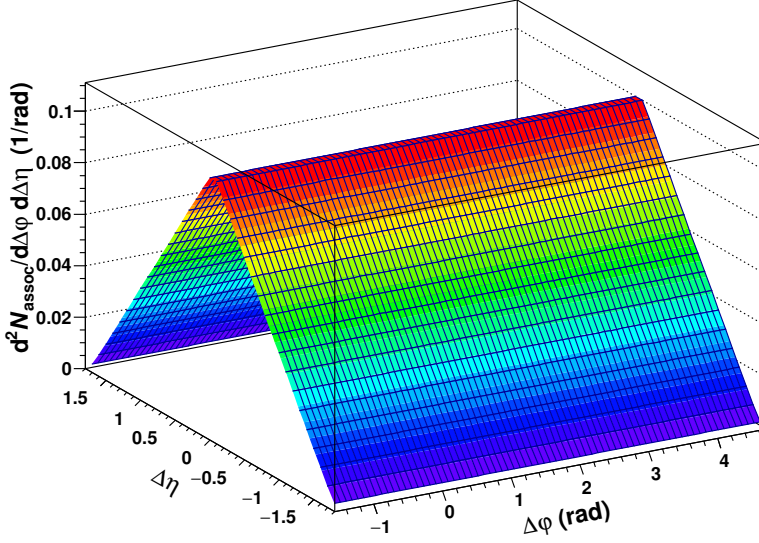


Figure 6.3: Mixed event pair yield M as a function of $\Delta\phi$ and $\Delta\eta$ for $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5$ GeV/ c for the highest (0–5%) V0A multiplicity class.

function. The so-called *wing effect* consists in large fluctuations that appear at the edge of the $\Delta\eta$ region, as shown in the left side of Fig. 6.4 for the 95–100% V0A multiplicity class, and it has to be corrected for. To do that, only the away side is projected into the whole $\Delta\eta$. Under the assumption that the away-side peak is flat in $\Delta\eta$, the projection is fitted with a constant function, as in the right side of Fig.6.4, and the correction factors needed to flatten the distribution are extracted point by point and applied to the measured data in the whole $(\Delta\eta, \Delta\phi)$ range.

Mandatory for the correction chain to be validated is that the correlation functions for Monte-Carlo reconstructed (after corrections) particles reproduce those for Monte-Carlo generated. All Monte-Carlo non-closure effects are to be added to the systematic uncertainty. This study is described in

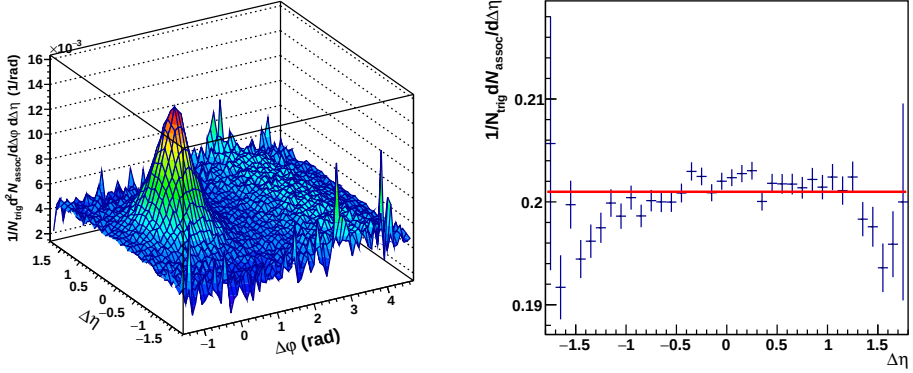


Figure 6.4: Left: $(\Delta\eta, \Delta\varphi)$ correlation function in the 95–100% V0A multiplicity class. Right: $\Delta\eta$ -projection of the away side, with a zero-degree polynomial fit.

Sec. 7.5.

6.1.2 Corrected per-trigger yield

It is interesting now to look at the features of the two-particle correlation after the above-mentioned corrections. Figure 6.5 shows the corrected per-trigger pair yield in the highest (0–5%) V0A multiplicity class for $0.7 < p_{T,assoc} < p_{T,trig} < 5$ GeV/ c . The near side presents the minijet peak around $(\Delta\eta, \Delta\varphi) = (0, 0)$. Based on this peak, the two subregions of short- and long-range correlations are defined as the ones, respectively, containing and excluding the peak. In the present analysis, $\Delta\eta_{cut}$ is set to 1.2, so that the short-range subregion $|\Delta\eta| < \Delta\eta_{cut}$ fully contains the near-side minijet peak in all multiplicity classes. However, the long-range near-side correlation is not flat in $\Delta\varphi$ either: the elongated bump in $\Delta\eta$ is the so-called ‘ridge’, similar to what has been found in heavy ion events and which suggests effects of collective origin. Both the short- and the long-range subregions in the away side show the same structure in $\Delta\varphi$, due to the fact that the

6.1 Two-particle correlations

recoiling jet contribution is elongated in $\Delta\eta$. This is by construction, i.e. the pseudorapidity difference of the pair can assume all values in the η window, and because of the longitudinal momentum distribution of partons in the colliding protons. Hence, it mixes up with the away-side ridge.

As a summary, these features appear in the two-particle correlations:

- near side ($-\pi/2 < \Delta\varphi < \pi/2$)
 - short range ($|\Delta\eta| < 1.2$): minijet peak on top of the ridge
 - long range ($|\Delta\eta| > 1.2$): ridge only
- away side ($\pi/2 < \Delta\varphi < 3\pi/2$)
 - whole range ($|\Delta\eta| < 1.8$): recoiling jet on top of the ridge.

The procedure described so far is quite general and employed by many correlation-type analyses.

However, towards an estimate of the sources of semi-hard scatterings, which is based on the particle counting (see Sec.6.3.2), this analysis requires that the *actual* number of particles traversing the acceptance-limited detectors is counted [122]. Basically, one needs to keep the whole correction settings provided by the mixed event technique, excluding the correction for the limited η acceptance. This is achieved by multiplying the correlation functions by the pair acceptance of a *perfect but pseudorapidity limited* detector $\tilde{B}(\Delta\eta) = 1 - |\Delta\eta|/(2 \cdot \eta_{\max})$, with $\eta_{\max} = 0.9$. In this way, the resulting associated yields per trigger particle count only the particles entering the detector acceptance.

The final two-dimensional two-particle correlation functions are, hence, calculated as:

$$\frac{S(\Delta\eta, \Delta\varphi)}{B(\Delta\eta, \Delta\varphi)} \cdot \tilde{B}(\Delta\eta). \quad (6.6)$$

Fig. 6.6 shows the final, i.e. fully corrected, two-dimensional two-particle correlation distribution for $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ for the highest (0–5%) V0A multiplicity class.

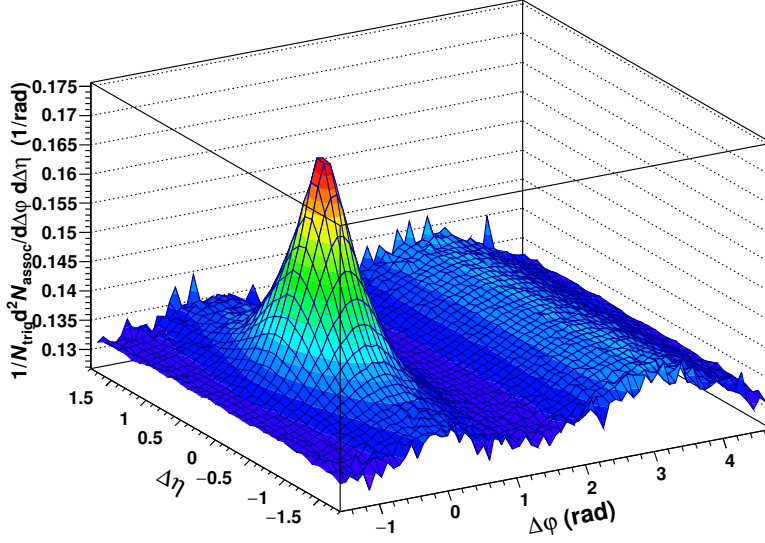


Figure 6.5: Mixed-event corrected per-trigger pair yield S/B as a function of $\Delta\varphi$ and $\Delta\eta$ for $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5$ GeV/ c for the highest (0–5%) V0A multiplicity class.

6.2 Long-range correlation subtraction

Ideally, one would like to remove completely the 'double-ridge' structure, to examine the multiplicity dependence of the jet-like component of the two-particle correlation function.

This seems feasible in the near side, where the minijet peak is limited in the short-range region. In fact, one can estimate the non-jet-related contribution in the long range and subtract it from the short range. As concerns the away side, the recoiling jet is elongated in the whole $\Delta\eta$ range, as already explained, thus another method has to be employed to isolate it from the ridge. The way to proceed is to use the information from the near

6.2 Long-range correlation subtraction

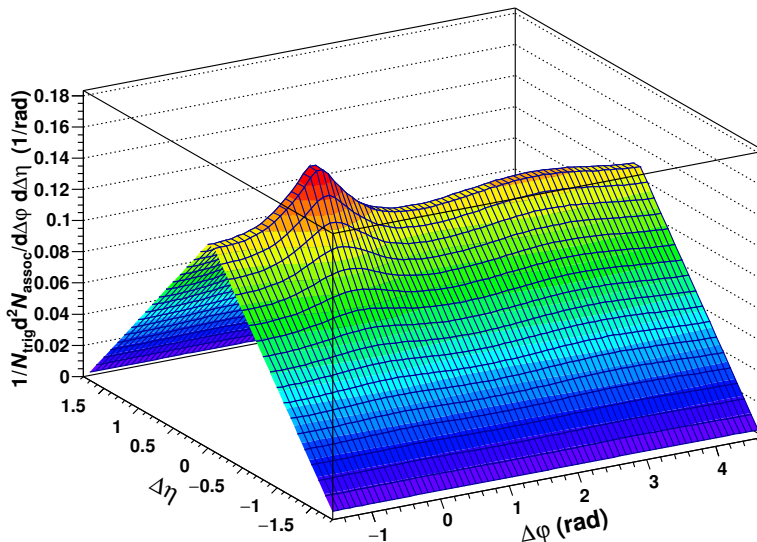


Figure 6.6: Final per-trigger pair yield $S/B \cdot \tilde{B}$ as a function of $\Delta\varphi$ and $\Delta\eta$ for $0.7 < p_{T,assoc} < p_{T,trig} < 5$ GeV/ c for the highest (0–5%) V0A multiplicity class.

side as an estimate for the away side, as described below.

The assumption in this procedure is that any $\Delta\eta$ dependence of this structure under the minijet peak is neglected. This is reasonable as no dependence is observed outside the peak. This allows one to perform the subtraction between the projections onto $\Delta\varphi$ of the 2-dimensional correlation functions. The modulation in azimuth of the ridge structures is well described by a Fourier expansion with non-zero second and third coefficients only:

$$\frac{1}{N_{trig}} \frac{dN_{assoc}}{d\Delta\varphi} = a_0 + 2a_2 \cos(2\Delta\varphi) + 2a_3 \cos(3\Delta\varphi). \quad (6.7)$$

Hence, one can assume that by mirroring the near-side long-range correla-

tion in the away side, i.e. with respect to $\Delta\varphi = \pi/2$, one gets a reasonable estimate of the ridge structure in both sides. In fact, the second order component is correctly estimated, being it even with respect to $\pi/2$. The mirroring does not work for the odd third component, though, which anyway is smaller than the second one. To assess the effect of the third coefficient on the extracted observables, an additional functional form of the third order is subtracted before the symmetrization, as a systematic study. This is explained in full details in Sec. 7.6.

The different contributions of the correlation function with $0.7 < p_{T,\text{assoc}} < p_{T,\text{assoc}} < 5.0$ GeV/ c are shown in Fig. 6.7, on the left for the 0–5% V0A multiplicity, on the right for the 95–100% V0A multiplicity. The blue points are the projection of the short-range correlation in $|\Delta\eta| < 1.2$. The black points are the long-range correlation in $1.2 < |\Delta\eta| < 1.8$ in both sides, scaled for the different $\Delta\eta$ intervals. The green markers represent the results of the symmetrization: the near side does not change by construction and is mirrored around $\Delta\varphi = \pi/2$. The red points in the lower panels show the result of the subtraction of the symmetrized-long-range function from the short-range one. For the highest multiplicity class (left panel), the away-side peak shows a small dip, due to the shape of the symmetrized third order component, which is drawn as a dotted line in the away side. The lowest multiplicity class is not affected by this contribution. No ridge has been indeed found in the low multiplicity correlation functions.

6.3 Final observables

The correlation functions extracted so far, deprived of the ‘double-ridge’ component, contain information about jet fragmentation, contribution of jets to the overall multiplicity and number of sources of semi-hard scatterings.

6.3 Final observables

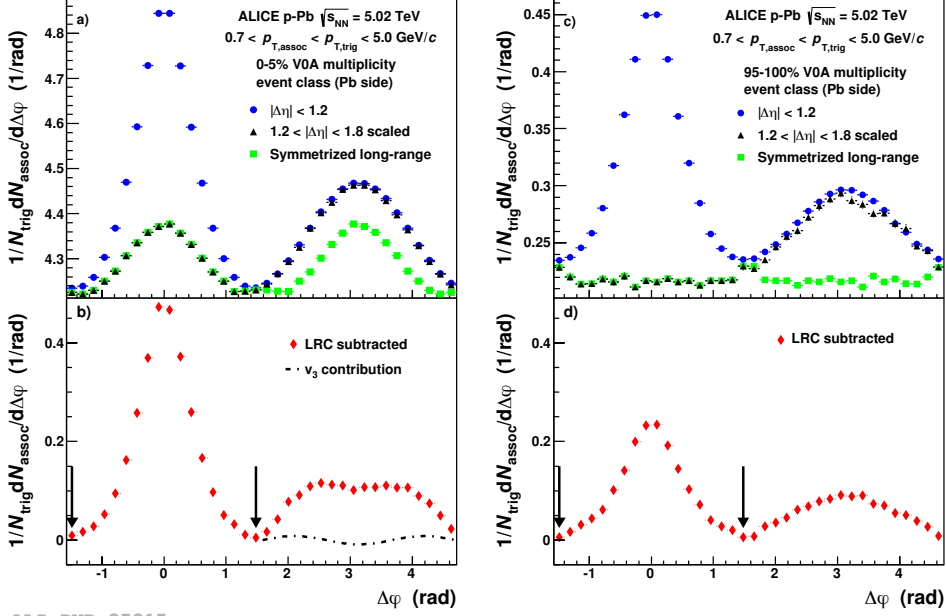


Figure 6.7: $\Delta\phi$ -projections of the correlation functions with $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5.0$ GeV/c in the 0–5% (left) and 95–100% (right) V0A multiplicity classes. The blue points show the short-range correlation, the black points represent the long-range correlation scaled for the different $\Delta\eta$ interval, the green points are the result of the symmetrization of the near-side ridge in the away side. The red points in the lower panels are the result of the long-range correlation subtraction. For the 0–5% multiplicity class, a dotted line is drawn to show the third order component of the Fourier expansion in the away side.

6.3.1 Per-trigger yields

As regards the amount of particles populating the jets, the best estimate is given by the so-called average per-trigger yields:

$$\langle N_{\text{assoc}} \rangle = \int_{\text{peak}} \frac{1}{N_{\text{trig}}} \frac{dN_{\text{assoc}}}{d\Delta\varphi} \quad (6.8)$$

where the integral is performed below the peak in the near side and in the away side, accordingly, and above the baseline. In fact, after the long-range correlation subtraction, the combinatorial background from uncorrelated particles should have been removed. However, due to minor differences between the real efficiencies and the ones estimated from the Monte Carlo and to a slight dependence on η of the single particle distribution, a small residual baseline is present. It is evaluated by subtracting the long-range from the short-range correlation in the whole $\Delta\varphi$ -range and then by fitting the remaining away side with a constant function. It amounts to 0.003, hardly visible in Fig. 6.7, and it has been further subtracted. The integral is calculated with the bin counting method. A fit procedure has been used for a systematic check and is discussed in Sec. 7.7. Since the away-side peak is slightly larger than the near side and extends further than $\pi/2 < \Delta\varphi < 3\pi/2$, the limits of integration for the away side are defined within $|\Delta\varphi| > 1.48$. The remaining region $|\Delta\varphi| < 1.48$ is assigned to the near side. The integration regions are indicated by two drawn arrows in Fig. 6.7. Those limits have been varied by ± 0.09 for the systematic uncertainty estimation. The average value of the per-trigger yields in the near and away side are denoted as $\langle N_{\text{assoc, nearside}} \rangle$ and $\langle N_{\text{assoc, away side}} \rangle$ respectively, and vary with event multiplicity.

6.3.2 Uncorrelated seeds

The average number of produced particles depends on the fragmentation properties of the partons as well as on the number of parton scatterings per event.

In this analysis, the average number of particles per event is the number

6.3 Final observables

of trigger particles. Hence, with the goal to reduce the dependence on fragmentation properties and thus to get an estimate of the number of *independent* parton scatterings, the average number of *all* trigger particles is divided by the number of *correlated* trigger particles. In fact, the denominator accounts for particles that belong to the same minijet pair of the trigger particle, meaning that they stem from the same semi-hard parton scattering. This is the reason why it can be estimated as the sum of the number of particles in the minijets, i.e. the per-trigger *associated* yields with *symmetric* p_T bins, and the trigger particle itself. This ratio, called average number of *uncorrelated seeds*, is defined as:

$$\langle N_{\text{uncorrelated seeds}} \rangle = \frac{\langle N_{\text{trig}} \rangle}{\langle N_{\text{correlated triggers}} \rangle} = \quad (6.9)$$

$$= \frac{\langle N_{\text{trig}} \rangle}{\langle N_{\text{assoc,nearside}} \rangle + \langle N_{\text{assoc,awayside}} \rangle + 1}. \quad (6.10)$$

Previous studies on proton-proton collisions simulated in PYTHIA [122] have shown that the uncorrelated seeds are linearly correlated to the number of multiparton interactions in a certain p_T range, independent of the η range explored. Figure 6.8 shows on the left the correlation between the average number of uncorrelated seeds and the number of MPI extracted from the Monte Carlo, with a $p_T > 0.7 \text{ GeV}/c$ selection for both trigger and associated particles, in $|\eta| < 1$ and $|\eta| < 10$. The ratios between data and linear fit are shown in the right panel. The linear trend is not affected by the choice of the pseudorapidity window and it is valid up to p_T thresholds of the order of GeV. The $p_T = 0.7 \text{ GeV}/c$ as the default cut for the present work is sufficiently high ($p_T > \Lambda_{\text{QCD}}$) to reduce contributions of decay products from short-lived particles, e.g. resonances and strings, but low enough to probe the bulk of particle production [122].

Thus, in case p-Pb collisions behave as a superposition of pp events, also in this analysis the average number of uncorrelated seeds is supposed to give a good estimate of the number of independent sources of minijet production, i.e. the number of MPI.

6. Analysis Procedure

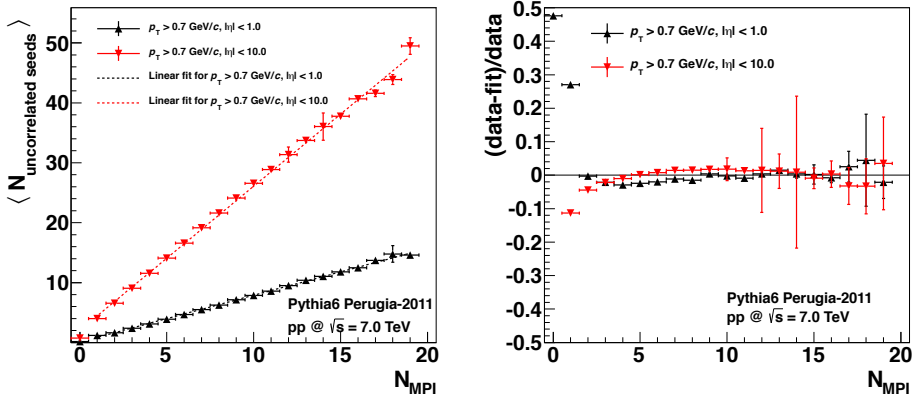


Figure 6.8: Left: correlation between average number of uncorrelated seeds and number of MPI extracted from PYTHIA, with p_T cut of $0.7 \text{ GeV}/c$ in two different pseudorapidity windows: $|\eta| < 1$ (red points) and $|\eta| < 10$ (black points). Right: ratio between data and linear fit. Figure taken from [122].

7

Systematic uncertainties

In this section, a collection of the contributions to the total systematic uncertainty is given. Table 7.1 collects the several contributions and the total systematic uncertainty for the per-trigger near-side and away-side yields and for the uncorrelated seeds. Most of the systematics on the yields do not propagate to the uncorrelated seeds, as they get reduced in the ratio (see Eq. 6.10).

7.1 Pile-up

The mean number μ of proton-nucleon interactions per bunch crossings in the analyzed sample is 0.0497, which translates into a pile-up fraction from the same bunch crossing:

$$\frac{P(n > 1)}{P(n > 0)} = 1 - \frac{\mu e^{-\mu}}{1 - e^{-\mu}} \simeq 0.03. \quad (7.1)$$

The out-of-bunch pile-up, i.e. from different bunch crossings, does not affect the SPD, whose integration time is smaller than the bunch spacing in p-Pb (200 ns). It can affect in particular the TPC and therefore a systematic study has been performed with the multi-vertexer method. It consists in classifying vertexes with different 'BCid' (bunch crossing identification

numbers), using information from the tracks reconstructed in the TOF. All the tracks with the same BCid are assigned to a unique vertex. If there is more than one vertex in the event, those vertexes are excluded¹ which have a number of contributors (tracks used to define the vertex) < 5 . The out-of-bunch pile-up contribution to the final results amounts to 1%.

7.2 Event generator

For this analysis, efficiencies are extracted with the DPMJET Monte Carlo. A comparison of the results obtained with the HIJING Monte-Carlo [120] efficiencies reveals negligible differences.

7.3 Track cuts

As described in Sec. 5.4, other than the standard track cuts for this analysis (filterbit 96), two other sets of cuts are used (filterbits 32 and 768) to estimate the dependence of the final results on the chosen track cuts. Results from each of the latter two sets are compared to the default set, by dividing per-trigger yields and uncorrelated seeds obtained with the different filterbits. One example for the per-trigger near-side yield is shown in Fig. 7.1, where the top panel shows the yields when using the different sets of cuts and the bottom panel shows the ratio with the default set of cuts. Two systematic variations are then estimated, one from the ratio between filterbits 96 and 32 (called σ_1), and one from the ratio between filterbits 96 and 768 (called σ_2), by fitting them with a zero-th order polynomial. Both $\sigma_{1,2}$ amount to slightly less than 1%.

7.4 Tracking efficiency uncertainty

The detector description in Monte-Carlo simulation is not as perfect as in reality. Therefore, discrepancies between data and Monte Carlo naturally

¹The cut on the number of contributors is optimized in order to have reliable vertices. It depends on the colliding system (for pp collisions it was set to 3).

7.4 Tracking efficiency uncertainty

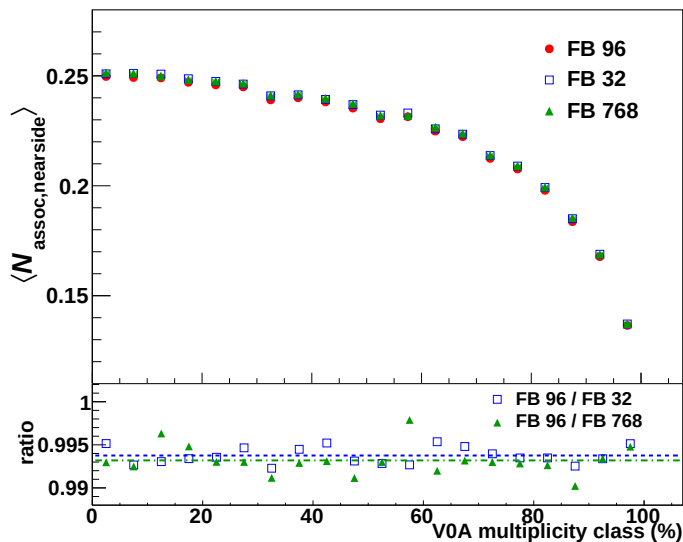


Figure 7.1: Top: Per-trigger near-side yield as a function of V0A multiplicity class, obtained with the default set of track cuts (filterbit 96), shown in red, and with two other sets of cuts (filterbit 32, shown in blu, and filterbit 768, shown in green). Bottom: ratio between the near-side yields with the default and with the changed cuts (32, shown in blu, 768 shown in green). Also shown as horizontal lines are the fit of the ratios with a zero-th order polinomial.

arise. Moreover, these discrepancies can depend on the cuts applied to the tracks. The dependence of the differences between data and simulations on the cut variation, the so-called tracking efficiency uncertainty, is estimated [123] to be of the order of 3% for p-Pb collisions at $\sqrt{s} = 5.02$ TeV.

7.5 Monte-Carlo closure test

Figure 7.2 shows the two $\Delta\varphi$ -projections of the Monte-Carlo generated and the Monte-Carlo reconstructed correlation functions, together with their ratio in the bottom panel. The differences are negligible. This results in a negligible contribution from Monte-Carlo non-closure effects to the systematic uncertainty for the yields and the uncorrelated seeds.

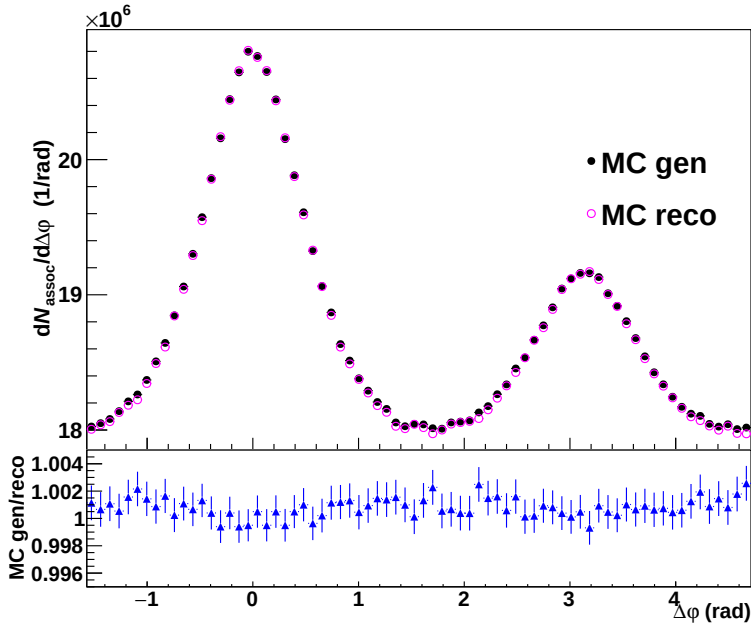


Figure 7.2: $\Delta\varphi$ -projections of the Monte-Carlo generated (black) and Monte-Carlo reconstructed (magenta) correlation functions (top) and their ratio (bottom).

7.6 v_3 effect

The mirroring of the near-side long-range correlation into the away side is the procedure adopted in this analysis to remove the ridge-like structure and study only the minijet contribution. However, it is not a perfect method, in the sense that only the even components are properly subtracted. The odd coefficient of the Fourier decomposition, namely the v_3 , which is the only non-negligible one, is subtracted from the near side but is added to the away side, resulting in twice its contribution to the total away-side correlation function.

Ideally, one would estimate the Fourier coefficients for each multiplicity class. The method employed in the two-particle correlation analysis to quantify the ridge [49], though, consists in a subtraction between high and low multiplicity event classes. Following the same method, it is then possible to estimate the v_3 for each multiplicity class, apart from the lowest one: the correlation function for the 95–100% multiplicity class is subtracted from all the other multiplicity classes and what is left from the subtraction is fitted with a Fourier decomposition. To quantify the subtracted contribution from the 95–100% class, in this work a *conservative* value of 0.01 has been chosen, of the same order of magnitude as the one measured by CMS at low multiplicity [124]. This is added to the v_3 values estimated for all multiplicity classes.

Once evaluated the value of v_3 for each multiplicity class, a functional form

$$2v_3^2 \cos(3\Delta\varphi) \tag{7.2}$$

is subtracted from the two-particle correlations and the standard analysis method is applied as described above. Figure 7.3 shows the $\Delta\varphi$ -projections of the two-particle correlations with (green points) and without (red points) subtraction of Eq. 7.2, in six multiplicity classes: 0–5% (top left), 5–10% (top right), 45–50% (middle left), 50–55% (middle right), 90–95% (bottom left), 95–100% (bottom right). No difference, as expected, is found in the low multiplicity classes, where no relevant ‘ridge’ structure is present. No difference either, by construction, is seen in the near-side peak in high multiplicity classes. The higher influence is in the away side at high multiplicity,

7. Systematic uncertainties

confirming the hypothesis on the origin of the away-side peak shape. In fact, by subtracting the third order component, one restores the gaussian shape of the peak. The effect of this subtraction on the amplitude of the away-side

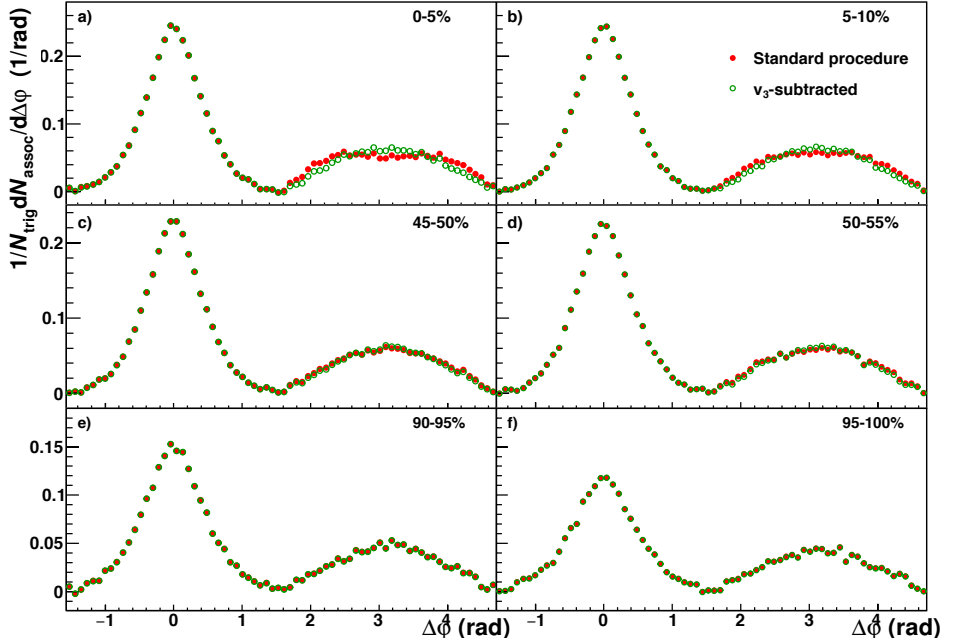


Figure 7.3: Two-particle correlation distributions as a function of azimuthal angular difference $\Delta\varphi$ for $p_{T\text{cut,trig}} = p_{T\text{cut,assoc}} = 0.7 \text{ GeV}/c$. The six panels show six different multiplicity classes: 0–5% (top left), 5–10% (top right), 45–50% (middle left), 50–55% (middle right), 90–95% (bottom left), 95–100% (bottom right). The different colors are associated to different procedures: red full circles for the standard, green empty circles for the third-order component subtraction.

peak goes from 0 to 4%, respectively from low to high multiplicity classes. This effect is taken into account only in the systematic uncertainties related

7.7 Yields from fit

to the away side and it is not used as a correction of the extracted yields.

7.7 Yields from fit

An alternative procedure to the bin counting has been employed to extract the yields. It is the same method used in [122]. The $\Delta\varphi$ -projections of the two-particle correlations are fitted with a composite function. On top of a constant for the combinatorial background, three Gaussians are summed: one for the away-side peak and two for the near side, as the latter presents an enhanced tail region, for which a superposition of two Gaussians with different widths is needed:

$$f(\Delta\varphi) = C + A_1 \exp\left(-\frac{\Delta\varphi^2}{2\sigma_1^2}\right) + A_2 \exp\left(-\frac{\Delta\varphi^2}{2\sigma_2^2}\right) + A_3 \exp\left(-\frac{(\Delta\varphi - \pi)^2}{2\sigma_3^2}\right). \quad (7.3)$$

The several contributions to the fit function are drawn separately in Fig. 7.4, together with the total function in red. The three Gaussians are drawn on top of the constant blue line for the uncorrelated background. One clearly sees that the fit well reproduces the data in the near side, while it fails to describe decently the away side. This is due to the long-range subtraction procedure, which gives as a result a non-Gaussian shaped peak. In fact, this is the stronger motivation in favor of the bin counting rather than the fit: no assumption on the peak shape is needed.

The difference in the yields obtained with the two methods is included in the systematic uncertainties and amounts to 5%.

7.8 Summary

The systematics across multiplicity classes are mostly correlated. The systematics are also multiplicity dependent, except the one from the mirroring of the v_3 . The total value is calculated with the formula of error propagation:

$$\sigma_{\text{TOT}} = \sqrt{\sum_i \sigma_i^2} \quad (7.4)$$

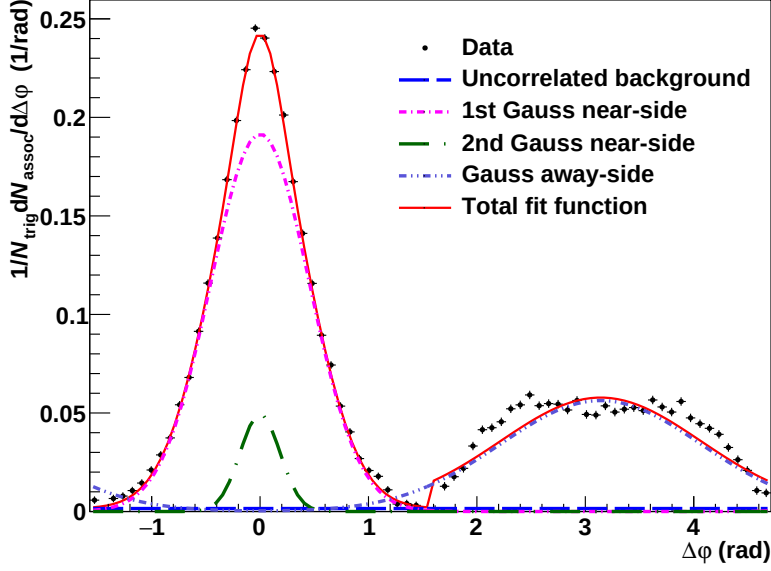


Figure 7.4: The $\Delta\varphi$ -projection of the two-particle correlation function in 0–5% multiplicity class is shown, together with the total fit function (red line) and the separate contributions: a constant for the combinatorial background (blue line), two Gaussians for the near-side peak (magenta and green curves), one Gaussian for the away-side peak (purple curve).

where σ_i indicates the singular contributions.

7.8 Summary

Table 7.1: Systematic uncertainties of this analysis. The several sources are listed, with their impact on the final observables, i.e. near-side yield, away-side yield and uncorrelated seeds. The last line contains the total uncertainties.

Sources	$\langle N_{\text{assoc,nearside}} \rangle$	$\langle N_{\text{assoc,awayside}} \rangle$	$\langle N_{\text{uncorrelated seeds}} \rangle$
Pile-up	1%	1%	negligible
Event generator	negligible	negligible	negligible
Track cuts	1%	1%	negligible
Tracking efficiency	3%	3%	3%
Monte Carlo closure	negligible	negligible	negligible
v_3	0%	0–4%	0–1%
Integration method	5%	5%	1%
Total	6%	6–8%	3%

8

Results

This chapter contains the results of this thesis.

In Sec. 8.1 the $\Delta\varphi$ -projections of the two-dimensional correlation functions are shown for several multiplicity classes and multiplicity estimators. In Sec. 8.2, similar $\Delta\varphi$ distributions, this time after long-range correlation subtraction, are shown. From the comparison with the previous ones, one can monitor the effect of the subtraction on the correlation functions for each multiplicity class. The average per-trigger near-side and away-side yields are shown as a function of multiplicity, for several p_T cuts for trigger and associated particles and for the different multiplicity estimators, in Sec. 8.3. Moreover, yields are compared from correlation functions with and without long-range subtraction. This comparison is the focus of Sec. 8.4. Finally, the average number of uncorrelated sources of particle production, namely the *uncorrelated seeds*, are presented in Sec. 8.5.

8.1 $\Delta\varphi$ per-trigger yields

The two-dimensional two-particle correlations have been projected over the whole $\Delta\eta$ range. Figure 8.1 shows the per-trigger yield as a function of $\Delta\varphi$, extracted using a p_T cut of 0.7 GeV/ c for both trigger and associated particles, with different multiplicity estimators: V0A, CL1 and ZNA. From

the top left to the bottom right, the six panels represent six multiplicity classes: 0–5% and 5–10% (top row), 45–50% and 50–55% (middle row) and 90–95% and 95–100% (bottom row). As a function of the azimuthal angular difference, one recognizes the two peaks at $\Delta\varphi = 0$ and $\Delta\varphi = \pi/2$ on top of the uncorrelated background. The two peaks contain the contribution from both the minijet particle multiplicity and the collective-like structure.

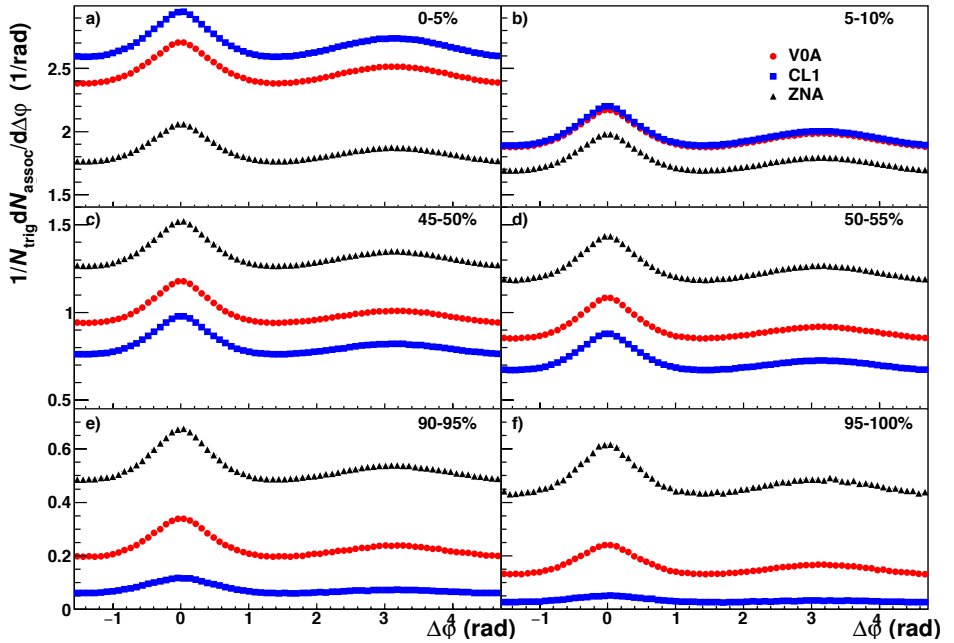


Figure 8.1: Two-particle correlation distributions as a function of azimuthal angular difference $\Delta\varphi$ for $p_{\text{Tcut,trig}} = p_{\text{Tcut,assoc}} = 0.7$ GeV/c. The six panels show six different multiplicity classes: 0–5% (top left), 5–10% (top right), 45–50% (middle left), 50–55% (middle right), 90–95% (bottom left), 95–100% (bottom right). Three different estimators are shown: V0A (red circles), CL1 (blue squares) and ZNA (black triangles).

8.2 $\Delta\varphi$ per-trigger long-range-subtracted yields

By looking at the highest multiplicity class, i.e. 0–5% in the top left panel, one sees that the $\Delta\varphi$ correlation for particles from events selected with CL1 is the highest, followed by the ones selected with V0A and ultimately the ones selected with ZNA. This is expected, as CL1 in 0–5% selects an average number of particles higher than the other estimator (see Table 5.1 or Fig. 5.4). Moving further towards lower multiplicity classes, the peaks and the uncorrelated background decrease. This is not surprising either, since events with lower multiplicity have less number of tracks by definition.

The lower-multiplicity events are characterized by more activity in the forward region than at midrapidity. Already in the 45–50% class (middle left), the highest $\Delta\varphi$ correlation is for tracks from ZNA-selected events, while the lowest is for tracks from CL1-selected events.

8.2 $\Delta\varphi$ per-trigger long-range-subtracted yields

The long-range subtraction allows isolating the minijet contribution. The six panels of Fig. 8.2 represent the jet-like two-particle correlations for six multiplicity classes and the three colors are assigned to different multiplicity estimators. As an effect of the subtraction, the baseline is removed, i.e. all the combinatorial pairs are discarded. The differences between the amplitude of the correlation functions for the three estimators are not striking, apart from the lowest-multiplicity classes. As regards the shape, the away-side peak in the highest multiplicity class appears to be flattened. This has been understood as an effect of the symmetrization of the odd components of the Fourier decomposition. More details are given in Sec.7.6.

8.3 Near-side and away-side per-trigger yields

The near-side and away-side per-trigger yields are studied as a function of multiplicity. Since the definition of multiplicity class in percentile bins is estimator-dependent, in order to compare the yields among different multiplicity estimators, they are reported as a function of $\langle N_{\text{ch}} \rangle^{\text{estimator}}(|\eta| <$

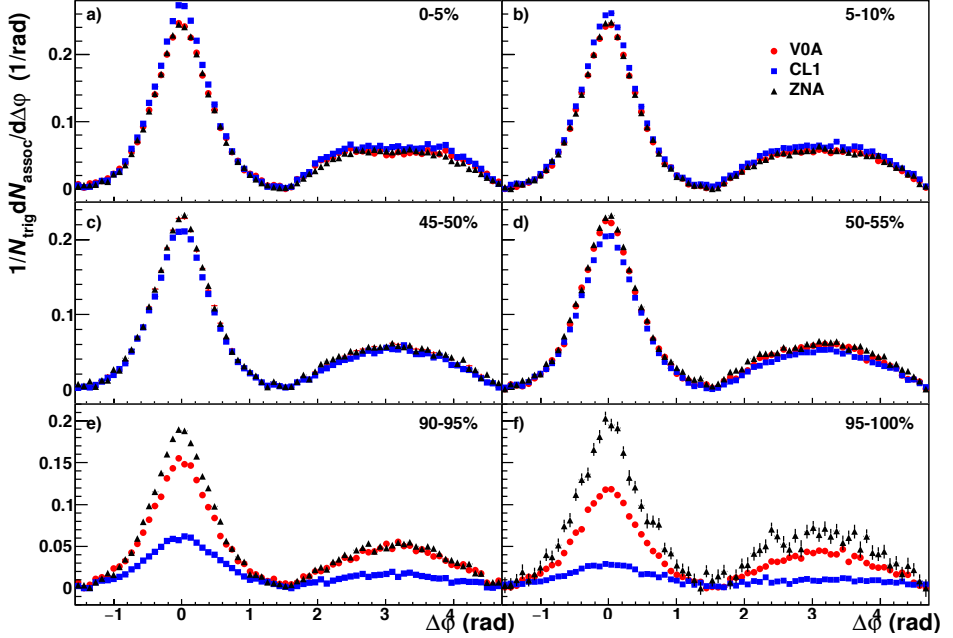


Figure 8.2: Two-particle correlation distributions after subtraction of the long-range component as a function of azimuthal angular difference $\Delta\phi$ for $p_{T,\text{cut},\text{trig}} = p_{T,\text{cut},\text{assoc}} = 0.7$ GeV/ c . The six panels show six different multiplicity classes: 0–5% (top left), 5–10% (top right), 45–50% (middle left), 50–55% (middle right), 90–95% (bottom left), 95–100% (bottom right). The different estimators are shown in different colors: red circles for V0A, blue squares for CL1 and black triangles for ZNA.

0.9, $p_T > 0.2$ GeV/ c). The correspondent values for each estimator are listed in Table 5.1. Shown in Fig. 8.3 are the near-side (left panel) and away-side (right panel) per-trigger yields for V0A, CL1 and ZNA. The multiplicity range covered by these estimators depends on their pseudorapidity acceptance (see Sec. 5.5): the ZNA covers the smallest range, followed by

8.3 Near-side and away-side per-trigger yields

the V0A and then by the CL1, which covers the largest one. The yields for CL1 exhibit the steepest slope. This behaviour is expected from the event-selection bias imposed by the overlapping η -region of event selection and tracking: on the near side the value increases from about 0.04 to 0.27 and on the away side from about 0.02 to 0.15. In the overlapping region between V0A and ZNA, the respective yields assume similar values: a mild increase from about 0.2 to about 0.25 in the near side and from about 0.1 to 0.13 in the away side is observed between 10 and 45 charged particles. Above 45 charged particles, the yields for V0A remain roughly constant, while below 10 charged particles, they decrease to 0.14 in the near side and 0.08 in the away side for V0A.

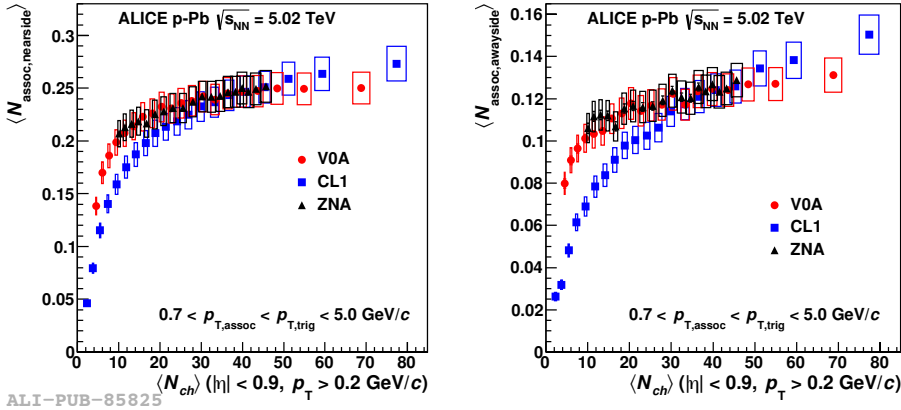


Figure 8.3: Near-side (left) and away-side (right) per-trigger yield as a function of average charged particle multiplicity at midrapidity $|\eta| < 0.9$ with $p_T > 0.2$ GeV/c for three multiplicity estimators: V0A (red circles), CL1 (blue squares), ZNA (black triangles). Statistical (lines) and systematic uncertainties (boxes) are shown, even though the statistical ones are mostly smaller than the symbol size.

The yields are further compared for different $p_{T\text{cut,trig}}$ and $p_{T\text{cut,assoc}}$. Figure 8.4 shows the near-side (left panel) and the away-side (right panel)

per-trigger yield as a function of V0A multiplicity class for different minimum p_T of the trigger and associated particles: the standard requirement $p_{T\text{cut,assoc}} = p_{T\text{cut,trig}} = 0.7 \text{ GeV}/c$ in red, $p_{T\text{cut,assoc}} = 0.7 \text{ GeV}/c$ and $p_{T\text{cut,trig}} = 2 \text{ GeV}/c$ in blue, and $p_{T\text{cut,assoc}} = p_{T\text{cut,trig}} = 2 \text{ GeV}/c$ in black. The different absolute value of the yields depend on the different number of possible combinations of trigger and associated particles: less pairs are possible in the case $p_{T\text{cut,assoc}} = 0.7$ and $p_{T\text{cut,trig}} = 2 \text{ GeV}/c$, but the number of triggers is smaller, hence the yield is higher. It is more interesting, though, to look at the trend of the yields with multiplicity in the different cases. For the standard requirement, the yields both in the near and in the away side increase in the lowest multiplicity classes, but remain nearly constant from 60% to the higher class. Same behaviour is observed for $p_{T\text{cut,assoc}} = 0.7 \text{ GeV}/c$ and $p_{T\text{cut,trig}} = 2 \text{ GeV}/c$, where the multiplicity region of plateau extends up to 80%. When both thresholds are increased $2 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$, no deviation from a flat trend is observed. This may suggest that the multiplicity dependence of the yields is mainly driven by the soft component of the particle spectrum.

The near-side and away-side per-trigger yields are defined as the number of associated particles normalized to the number of trigger particles, in a certain multiplicity class of events. To understand the meaning of the trend of the per-trigger yield with multiplicity, one has to focus on the processes in which the trigger and associated particles are produced. The only requirement for trigger particles is to have their p_T above a certain threshold. Hence, all particles both from *semi-hard* and *soft* processes can act as *trigger* particles. On the other hand, a particle with its p_T above a certain threshold plays the role of *associated* particle by definition if it belongs to the minijet pinpointed by the trigger particle, thus it originates from *semi-hard* processes. Therefore, in the region where the associated yields per trigger particle show a plateau, the hard processes and the number of soft particles must exhibit the same evolution with multiplicity. To clarify this concept, one can imagine an example event containing N_{minijets} with N_{assoc} associated particles each and a background of N_{soft} particles with no

8.3 Near-side and away-side per-trigger yields

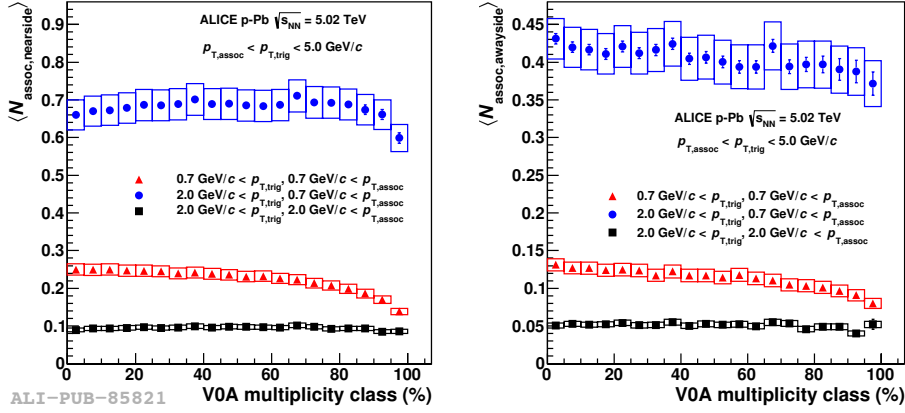


Figure 8.4: Near-side (left) and away-side (right) per-trigger yield as a function of V0A multiplicity class for three different trigger and associated particle selection: $0.7 < p_{T,assoc} < p_{T,trig} < 5 \text{ GeV}/c$ (red), $0.7 < p_{T,assoc} < 5 \text{ GeV}/c$ and $2 < p_{T,trig} < 5 \text{ GeV}/c$ (blue), $2 < p_{T,assoc} < p_{T,trig} < 5 \text{ GeV}/c$ (black). Statistical (lines) and systematic uncertainties (boxes) are shown, even though the statistical ones are mostly smaller than the symbol size.

azimuthal correlation. In this scenario, the per-trigger associated yield is:

$$\frac{\text{associated yield}}{\text{trigger particle}} = \frac{N_{\text{minijets}} \cdot N_{\text{assoc}}(N_{\text{assoc}} - 1)/2}{N_{\text{minijets}} \cdot N_{\text{assoc}} + N_{\text{soft}}} \quad (8.1)$$

So, for different events with different multiplicity, which translates into different number of trigger particles in the denominator of the equation, the ratio is constant if N_{minijets} from hard processes and N_{soft} from soft processes increase by the same factor. The given example can be extended to several events and to a different number of associated particles per minijet.

8.3.1 Comparison with pp results

A comparison with the published pp results for the near-side and away-side per-trigger yields can give some hints on the similarities and differences in

the physical phenomena involved in the different collision systems.

Such a comparison is shown in Fig. 8.5 for the near-side and away-side per-trigger yield as a function of charged-particle multiplicity at midrapidity for pp collisions at $\sqrt{s} = 0.9$ TeV, 2.76 TeV and 7 TeV and for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Since the dependence on the energy is small as one can see in Fig. 8.5, one can expect to have similar trend for pp collisions at $\sqrt{s} = 5.02$ TeV. Therefore, in the comparison between pp and p-Pb one can consider negligible the difference in center-of-mass energy.

The caveat in the shown comparison is that the definition of event multiplicity is different in the two analyses. In pp, events are not subdivided in multiplicity classes: the N_{ch} is measured with detectors at midrapidity and corrected by an unfolding procedure and is shown in the x axis in bins of unitary width. In p-Pb, instead, event multiplicity is defined as the average of the number of charged particles in a certain multiplicity class measured with either forward (V0A, ZNA) or *central* (CL1) detectors. However, one can qualitatively look at the values from CL1 in p-Pb, coherently with the pseudorapidity range of multiplicity measurement in pp.

The CL1 trend looks qualitatively consistent with the pp published results: the yields increase with no apparent stagnation at high multiplicity.

8.4 Jet-like yields and double ridge

So far, the per-trigger yields have revealed interesting properties of the minijet fragmentation and particle production as a function of multiplicity. Additional information can be obtained by comparing the jet-like per-trigger yields, i.e. after long-range correlation subtraction, with the per-trigger yields before subtraction. The latter, in fact, keep the information about the particles that belong to the 'ridge' structure.

The determination of the yields without long-range subtraction is, however, slightly different. The non-subtracted distribution does not have a zero baseline by construction. In this case, the baseline is determined in the long-range correlations region ($1.2 < |\Delta\eta| < 1.8$) between the near-side ridge and the away-side peak at $1.05 < |\Delta\varphi| < 1.22$.

8.4 Jet-like yields and double ridge

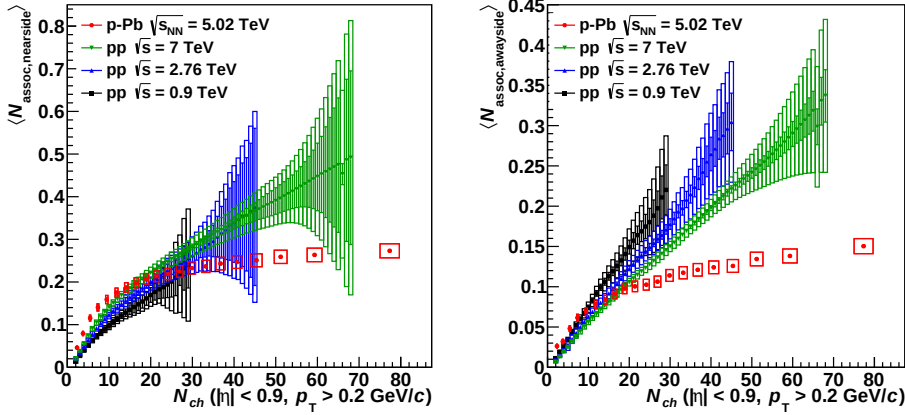


Figure 8.5: Near-side (left) and away-side (right) per-trigger yields with $p_{T,\text{cut,assoc}} = p_{T,\text{cut,trig}} = 0.7 \text{ GeV}/c$ as a function of charged particle multiplicity in p-Pb collisions estimated with CL1 (red circles) and in pp collisions at $\sqrt{s} = 0.9$ (black squares), 2.76 (blue triangles) and 7 TeV (green inverted triangles). Data for pp collisions taken from [122]. The N_{ch} in the x axis is unfolded with central barrel detectors in pp results, while it is averaged in CL1 multiplicity classes in p-Pb results, as the scope of these analyses were different.

Figure 8.6 shows the near-side (left) and away-side (right) per-trigger yield as a function of V0A multiplicity class, extracted for $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$. The yields agree with each other in the multiplicity classes from 50% to 100%. Nonetheless, unlike the case with long-range subtraction, the yields obtained without any subtraction do not show a plateau from intermediate to high multiplicity class: the near-side yield increases up to about 0.34, compared to the 0.25 after subtraction, while for the away side the value goes up to about 0.23, compared to 0.13. Thus, in the highest multiplicity class, the subtraction procedure removes 30–40% of the measured yields. The same observation is valid for other multiplicity estimators.

In conclusion, the scaling of the hard processes and the number of soft particles with the multiplicity of the event is only true for jet-like correlations, i.e. only when long-range correlations are subtracted. This may imply that the physical mechanism at the basis of the 'ridge' structure has a different nature from the physics of the minijet fragmentation. The observation is consistent with a picture where the minijet-associated yields in p-Pb collisions originate from the incoherent fragmentation of multiple parton-parton scatterings, while the long-range correlations appear unrelated to minijet production.

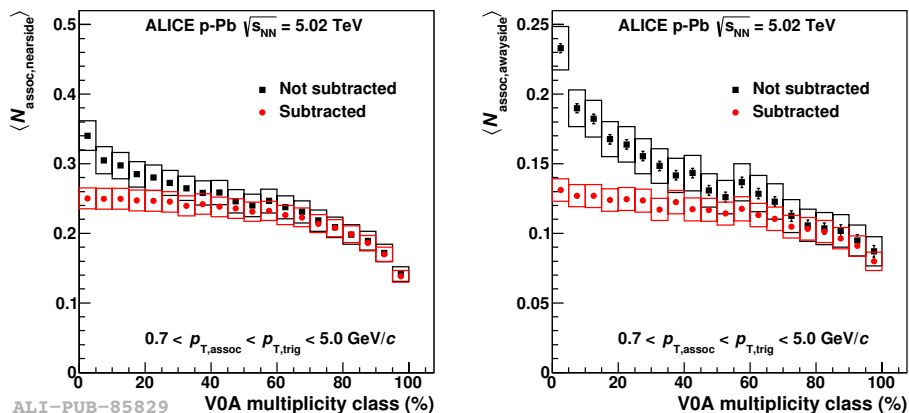


Figure 8.6: Near-side (left) and away-side (right) per-trigger yield with $p_{T\text{cut,assoc}} = p_{T\text{cut,trig}} = 0.7\text{ GeV}/c$ as a function of V0A multiplicity class, with (red circles) and without (black squares) long-range correlation subtraction. Statistical (lines) and systematic uncertainties (boxes) are shown, even though the statistical ones are mostly smaller than the symbol size.

8.5 Uncorrelated sources of minijet production

As demonstrated in Sec. 6.3.2, this analysis is able to give an estimate of the number of uncorrelated sources of minijet production, the so-called

8.5 Uncorrelated sources of minijet production

uncorrelated seeds. These have been found to be proportional to the number of MPI in PYTHIA.

As for their definition, the uncorrelated seeds can only be calculated for symmetric p_T bins for trigger and associated particles. Figure 8.7 shows the uncorrelated seeds estimated in V0A multiplicity classes, as a function of $\langle N_{ch} \rangle^{V0A}(|\eta| < 0.9, p_T > 0.2 \text{ GeV}/c)$. The results for two different p_T cuts are presented. In the range $0.7 < p_{T,assoc} < p_{T,trig} < 5 \text{ GeV}/c$ the seeds increase from 2 to about 20, while in $2 < p_{T,assoc} < p_{T,trig} < 5 \text{ GeV}/c$ they go up from 0 to about 3. The uncorrelated seeds increase linearly with midrapidity charged particle multiplicity, especially at high $\langle N_{ch} \rangle$. To quantify this behavior, a linear fit is performed in the 0–50% multiplicity range (full markers are included, empty markers are excluded). The dashed lines draw the fit for the two sets of data points. The ratio between data and fit, shown in the lower panel with the band of systematic uncertainty, reveal that the linear description of the data is valid for $N_{ch} > 20$ (a slightly wider range of linearity for the smallest p_T range), while deviations at lower multiplicity are observed. Those deviations are not surprising as other observables, e.g. $\langle p_T \rangle$ [125] and R_{pA} [126] show a change in dynamics as a function of multiplicity.

This result, together with the flatness of the yields, suggests that events with high multiplicity are not built up by a higher number of associated particles in the minijets, but rather by a higher number of uncorrelated sources of particle production, i.e. a higher number of independent minijets.

8.5.1 Comparison with pp results

The study in [122] has examined the evolution of the average uncorrelated seeds with multiplicity in pp collisions at several center-of-mass energies: $\sqrt{s} = 0.9 \text{ TeV}, 2.76 \text{ TeV}, 7 \text{ TeV}$. The ratio between the data and a linear fit is shown in Fig. 8.8. The deviation from linearity is important in the regions of low and high multiplicity. In the low-multiplicity region, the behavior is the same as in p–Pb. In the high-multiplicity region, the deviation increases with center-of-mass energy and, unlike in p–Pb, data deviate from linear fit up to 10%. In this scenario, high multiplicities may be obtained with a

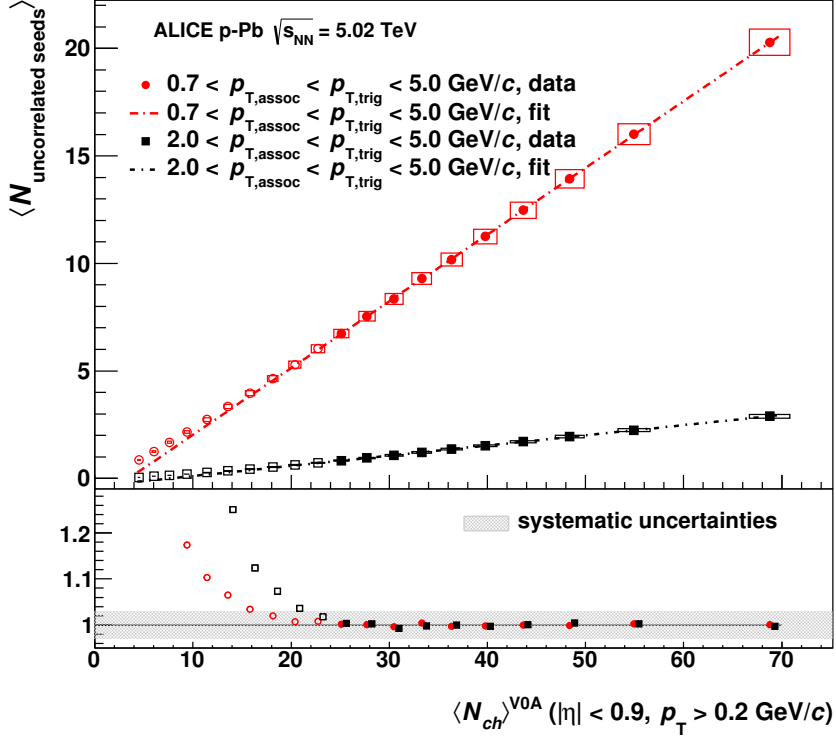


Figure 8.7: Top panel: number of uncorrelated seeds as a function of $\langle N_{\text{ch}} \rangle^{\text{V0A}} (|\eta| < 0.9, p_{\text{T}} > 0.2 \text{ GeV}/c)$. Shown are results for two p_{T} cuts: $0.7 < p_{\text{T,assoc}} < p_{\text{T,trig}} < 5 \text{ GeV}/c$ (red circles) and $2 < p_{\text{T,assoc}} < p_{\text{T,trig}} < 5 \text{ GeV}/c$ (black squares). Each of them is fitted with a linear function in the 0–50% multiplicity classes; open symbols are not included in the fit. Statistical (lines) and systematic uncertainties (boxes) are shown, even though the statistical ones are smaller than the symbol size. Bottom panel: ratio between the number of uncorrelated seeds and the linear fit functions. Black points are displaced slightly for better visibility.

8.5 Uncorrelated sources of minijet production

high number of associated particles in the minijets.

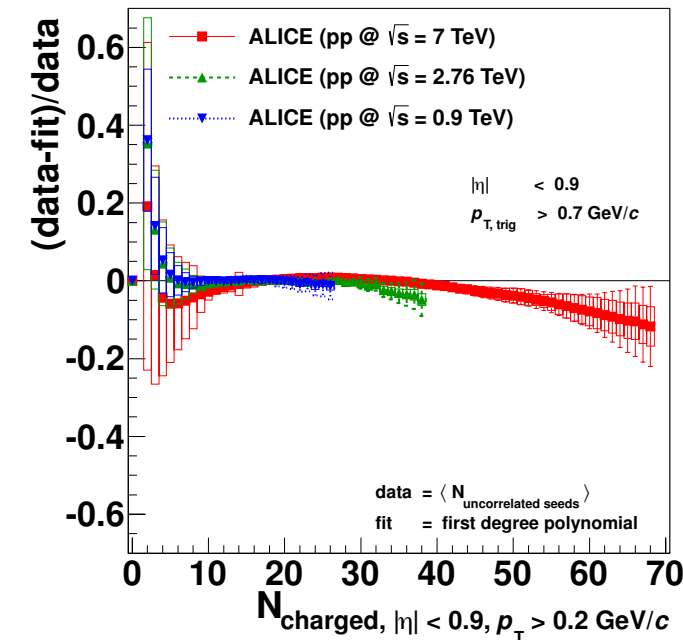


Figure 8.8: Ratio between average number of uncorrelated seeds and linear fit for $\sqrt{s} = 0.9$ TeV (blue triangles), 2.76 TeV (green triangles) and 7 TeV (red squares) as a function of charged particle multiplicity at midrapidity ($|\eta| < 0.9$) with $p_T > 0.2$ GeV/c. Figure taken from [122].

The average number of uncorrelated seeds measured in p-Pb collisions in bins of CL1 multiplicity is shown in Fig. 8.9 together with the pp data. The results for the different collision systems agree within systematic errors, although the uncorrelated seeds obtained in p-Pb collisions appear to be slightly higher than expected for pp collisions at similar energies.

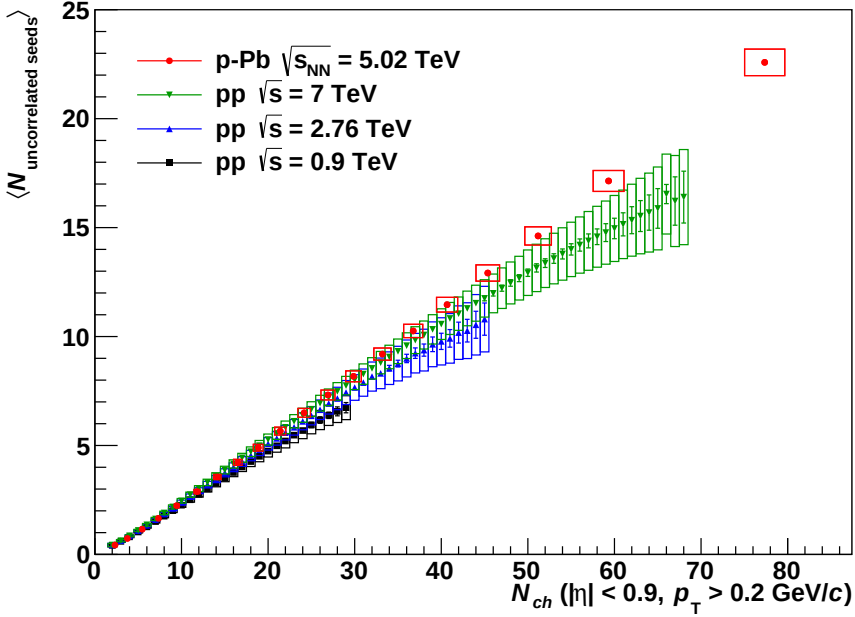


Figure 8.9: Number of uncorrelated seeds with $p_{T\text{cut,assoc}} = p_{T\text{cut,trig}} = 0.7 \text{ GeV}/c$ as a function of charged particle multiplicity in p-Pb collisions estimated with CL1 (red markers) and in pp collisions at $\sqrt{s} = 0.9$ (black), 2.76 (blue) and 7 TeV (green).

8.5.2 From MPI to nucleon–nucleon collisions

The uncorrelated seeds are proportional to the number of MPI, which, as already discussed in Sec. 4.2.2, could be affected by some biases. In fact, the latter is found not to scale with N_{coll} at low and high multiplicity, because of a multiplicity-based selection of event classes. In particular, the average number of uncorrelated seeds is higher than the average N_{coll} at high multiplicity and lower at low multiplicity, meaning that one selects event with respectively higher and lower number of semihard processes.

8.5 Uncorrelated sources of minijet production

Figure 8.10 shows the multiplicity dependence of the ratio between the average number of uncorrelated seeds and the average N_{coll} values extracted from a Glauber Monte Carlo in V0A multiplicity classes. These values are listed in Tab.4.1. Data points are plotted for two sets of p_T cuts: $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ and $2 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$. The latter are scaled so to have the same value as the former in the 50–55% multiplicity bin. A scaling only holds within 3% between 25% and 55% multiplicity classes for both sets. At high multiplicity, the deviation from the scaling amounts up to 25% for the $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ and to 60% for the $2 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$. At low multiplicity, the deviation amounts down to 30% for the former and 25% for the latter p_T cut. The breaking of the scaling means that the number of uncorrelated seeds is not strictly proportional to the number of binary collisions. However, as showed in Fig. 4.4, some of these deviations could be due to a bias induced by the centrality estimator, which reflect in a bias in the number of MPI. In addition, the geometric bias affects the lowest multiplicity classes, where the number of MPI is expected to decrease. Concluding, the bias in the number of uncorrelated seeds reflects the same multiplicity and geometric biases undergone by the number of MPI. This would confirm that this observable is a good estimation of the number of MPI.

The same verification has been made using the ZNA-based selection for event classes. In Fig. 8.11, the same ratio as Fig. 8.10 is shown, but the average number of uncorrelated seeds is computed in ZNA classes and the average number of N_{coll} is calculated from the hybrid model. More details about the hybrid model and a list of the $\langle N_{\text{coll}} \rangle$ values can be found in Sec. 4.3.2 and Tab.4.1. As expected, the biases are very much reduced (only the geometric bias is expected to arise deviations at low multiplicity). For $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ a deviation within 20% appears, in the opposite direction than Fig. 8.10. Moreover, the bias seems to be dampened the higher the p_T cut: for $2 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ the multiplicity dependence of the ratio is definitely small.

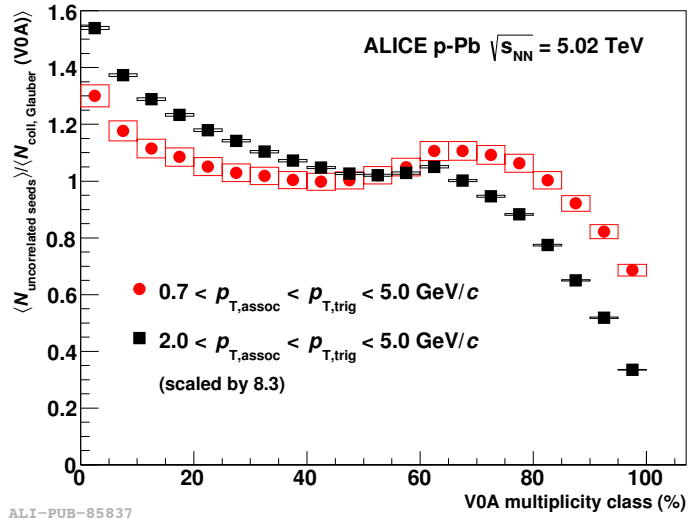


Figure 8.10: Multiplicity dependence of the ratio between average number of uncorrelated seeds and $\langle N_{\text{coll}} \rangle$ estimated with a Glauber Monte Carlo in V0A multiplicity classes. The red circles correspond to the $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ range, while the black squares to the $2 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5 \text{ GeV}/c$ range. The latter has been scaled by a factor of 8.3, in order to have it superimposed to the former in the 50–55% multiplicity bin.

8.5 Uncorrelated sources of minijet production

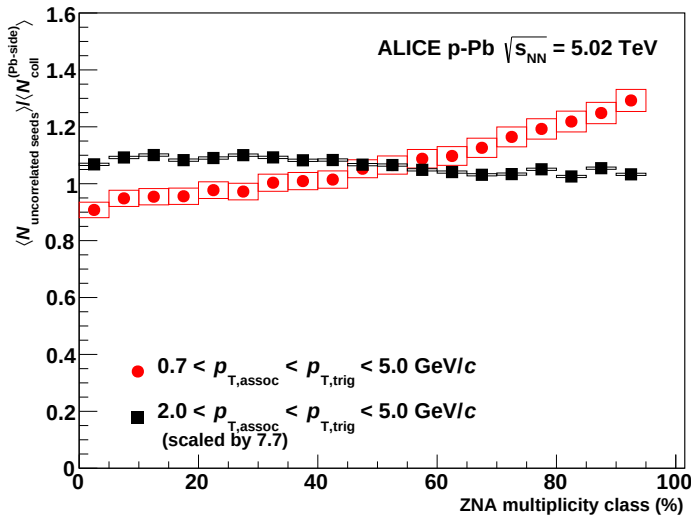


Figure 8.11: Multiplicity dependence of the ratio between the average number of uncorrelated seeds in ZNA multiplicity classes and the average number of N_{coll} estimated via the hybrid model (with inner-ring V0A multiplicity scaling in ZNA multiplicity classes). The red circles correspond to the $0.7 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5$ GeV/ c range, while the black squares to the $2 < p_{T,\text{assoc}} < p_{T,\text{trig}} < 5$ GeV/ c range. The latter has been scaled by a factor of 7.7, in order to have it superimposed to the former in the 50–55% multiplicity bin.

Summary

The Large Hadron Collider has provided collisions at unprecedented energy and a plethora of new and intriguing results have emerged from the analysis of the collision data. In particular, after a pilot run in September 2012 and a full run between January and February 2013, data from collisions of proton and lead-ion beams have been available. Those systems are thought to be sensitive to the cold nuclear matter, i.e. the modification of the matter due to the binding of nucleons in the nuclei. Therefore, they are supposed to be an important reference for the heavy-ion events. However, puzzling results have raised the questions whether some form of collectivity, similar to the effect attributed to the Quark Gluon Plasma in heavy-ion collisions, is present also in smaller colliding systems. The interpretation is still under debate also because some of the results can be explained by models for initial conditions, such as a state of gluon condensate.

Physics with hadronic beams is complicated by the fact that each colliding particle is constituted of a system of elementary partons, which can participate simultaneously in binary parton-parton collisions to a single hadronic interaction. This phenomenon is certainly important as a background for many physics processes, but it is especially crucial in order to fully understand the particle production in hadron collisions. The first evidence of multiple parton interactions (MPI) came from the study of charged-particle multiplicity distribution in the final state. Nowadays, the MPI are included in many of the modern simulating softwares, such as the PYTHIA Monte-Carlo, which has been reviewed in this thesis.

When the MPI are highly energetic, meaning that a high momentum is transferred in the interaction, they produce distinctive topological events, e.g. 4-jets balanced in pairs. On the other hand, when they happen at low energy, they produce low-energy jets (minijets) which are impossible to reconstruct with the standard jet-reconstruction algorithms. Due to its low material budget, ALICE is highly sensitive to low-momentum particles. Thus it is perfectly suited for studying the particle production in the soft regime, where theoretical calculations are limited by the non-applicability of perturbative QCD.

In this thesis, the minijet analysis is performed with the method of two-particle correlations. It consists of pairing particles, according to a selection of transverse momentum, and studying the distribution of their difference in pseudorapidity and azimuthal angle. The analysis strategy combines the information from the two-particle correlations with the event activity, i.e. the charged-particle multiplicity of each event.

In the two-particle correlations, not only the peaks from the minijets appear, but also structures that span a large range in pseudo-rapidity, known as double ridge. In order to examine the jet-like correlations, the ridge-like structures have to be removed. This thesis describes a new subtraction method to disentangle the two contributions.

The number of particles in each minijet are then measured for each event class, characterized by a certain average multiplicity. The first important result of this thesis is the finding that the jet-like yield is invariant with the multiplicity of the event, with a small deviation only at very low multiplicities, which however happens to disappear with a selection of higher-momentum particles. The invariance of the yields is only valid when the multiplicity is measured with a different detector than the one used for detecting the particles to be correlated, otherwise an autocorrelation bias is introduced and the trend becomes slightly increasing.

Summary

This analysis has also been able to compare the multiplicity dependence of the jet-like yields and of the yield in the ridge structures. In fact, without subtracting the latter component, the two-particle correlation yield smoothly increases with the event multiplicity. This is the first hint that different physical phenomena are responsible for the minijet production and for the long-range correlation structures. These findings are consistent with a picture where independent parton-parton scatterings with subsequent incoherent fragmentation produce the measured minijet associated yields.

The measurement of the minijet yields is crucial also for another reason. It has been shown that it is possible to estimate the number of clusters of particle production and this is proportional to the number of MPI in PYTHIA simulations. The so-called uncorrelated seeds, which are defined as the uncorrelated sources of minijet-pair production, have been measured in this thesis work for the first time in p-Pb collisions. Their average number is computed in each event class, resulting in a linear correlation with the average multiplicity of the event, apart from low-multiplicities region, where the probability of MPI to occur is very low. The interpretation of this result is that events with higher multiplicity in the final state, which, as discussed earlier, are not characterized by more populated minijets, but rather contain a larger number of uncorrelated sources of minijets. In other words, the higher multiplicity is due to a higher number of MPI.

Further, it has been checked whether the number of uncorrelated seeds scales with the number of binary hadron-hadron collisions N_{coll} . It has been found that the scaling holds in the intermediate multiplicity region. On the other hand, high (low) multiplicity events are characterized by higher (lower) number of uncorrelated seeds with respect to a linear trend. A similar bias, with the same multiplicity dependence, had been previously found in studying the scaling with N_{coll} of the number of MPI from Monte-Carlo. The confirmation of such bias validates the reliability of the uncorrelated seeds as an estimate of the number of multiple parton interactions.

In summary, this thesis provides an indirect estimate of the behavior

of MPI and the minijet production in the regime of soft probes, where experimental results are crucial for constraining theoretical models. This measurement, in particular, imposes significant constraints on models which aim at describing p–Pb collisions: they must reproduce such an incoherent superposition while also describing observations like the ridge structures and the increase of mean p_T with event multiplicity. This work is therefore expected to help shed light on the current investigation of the intriguing small colliding systems.

Samenvatting

De Large Hadron Collider produceert botsingen van een ongeken- de energie en uit de analyse daarvan is een overvloed aan nieuwe en intrigerende resul- taten voortgekomen. Na een pilot run in september 2012, heeft in januari en februari een volledige run plaatsgevonden met botsingen van een bundel loodkernen met een bundel protonen. De verwachting was dat met deze botsingen de effecten van koude kernmaterie, zoals veranderingen in de nu- cleonstructuur door de binding in de loodkern, onderzocht kon worden. De resultaten zouden dan te gebruiken zijn als referentie voor botsingen tussen loodkernen. Er is echter een aantal onverwachte effecten waargenomen, die het best te verklaren zijn als er collectief gedrag optreedt, dat nor- maal gesproken wordt toegeschreven aan de de vorming van een quark- gluonplasma. Het was niet verwacht dat deze effecten ook zouden optreden in kleine botsingssystemen, zoals proton-loodbotsingen. De interpretatie van de resultaten wordt actief bediscussieert, met name omdat sommige re- sultaten ook verklaard kunnen worden door modellen van de begintoestand waarin de toestand van de kernmaterie wordt beschreven als een gluoncon- densaat.

Protonen en neutronen zijn samengestelde deeltjes die bestaan meer fun- damentele quarks en gluons, gezamenlijk partons genoemd. In een botsing tussen protonen of protonen en loodkernen kunnen tegelijkertijd meerdere partonbotsingen plaatvinden. Dit fenomeen is een belangrijke achtergrond voor andere productieprocessen, maar is ook van cruciaal belang om de deeltjesproductie in hadronbotsingen te begrijpen. De eerste aanwijzingen

dat meerdere partoninteracties (MPI) kunnen plaatsvinden in een botsing zijn gevonden in de studie van de verdeling van het aantal deeltjes (de zgn. multipliciteit) dat in botsingen wordt geproduceerd. De effecten van MPI zijn sindsdien opgenomen in moderne botsingssimulatiesoftware, zoals de PYTHIA Monte-Carlo, die in dit proefschrift wordt beschreven.

Als de MPI gepaard gaan met grote energie- en impulsoverdracht, produceren ze gebeurtenissen met een specifieke topologische structuur, zoals vier jets met impulsbalans in twee jet-paren. Bij lagere impulsoverdracht worden mini-jets geproduceerd, die niet gereconstrueerd kunnen worden met standaard jet-reconstructie-algoritmen. Door de lichte constructie (laag materie-budget), is ALICE bijzonder geschikt voor het detecteren van deeltjes met lage impuls. Daardoor is de detector bij uitstek geschikt voor het bestuderen van het zogenaamde zachte regime voor deeltjesproductie, waar theoretisch inzicht beperkt is, doordat storingsrekeningstechnieken voor QCD niet van toepassing zijn.

In dit proefschrift worden twee-deeltjescorrelaties gebruikt als methode om minejets te analyseren. De analysemethode bestaat eruit deeltjes te selecteren binnen een impulsbereik en de verdelingen van het azimutale hoekverschil en het pseudorapiditeitsverschil van paren van deeltjes te bestuderen. Bij de analyse worden de gebeurtenissen verdeeld in klassen met verschillende activiteit, oftewel het aantal geproduceerde deeltjes (multipliciteit) in de botsing.

In de twee-deeltjesverdeling zijn naast de pieken van mini-jets ook structuren zichtbaar over een groot pseudorapiditeitsverschil, de zogenaamde 'double ridge'. Om de jet-achtige correlaties te onderzoeken, moet de bijdrage van de double ridge eerst worden verwijderd. Dit proefschrift beschrijft een nieuwe methode om de twee bijdragen te onderscheiden.

Met deze methode is het aantal deeltjes per mini-jet gemeten, voor gebeurtenissen met verschillende gemiddelde activiteit. Het eerste belangrijke resultaat van dit proefschrift is de waarneming dat het aantal deelt-

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jes dat per mini-jet wordt geproduceerd niet afhangt van de multiplicititeit van de botsingen. Alleen bij gebeurtenissen met een heel klein aantal geproduceerde deeltjes is een kleine afwijking van dit gedrag te zien, en dan alleen voor deeltjes van lage impuls. De onafhankelijkheid van het aantal deeltjes per mini-jet van het totaal aantal deeltjes in de gebeurtenis is alleen geldig als de totale deeltjesproductie in een ander pseudorapiditeitsgebied wordt gemeten dan waar de mini-jets worden gemeten; anders treedt er een autocorrelatieeffect op, waarbij de multiplicititeit per minijet wel verandert als bijvoorbeeld gebeurtenissen met een groot aantal deeltjes in hetzelfde rapiditeitsgebied worden geselecteerd.

Een tweede belangrijk resultaat van de analyse is het aantal deeltjes in de double-ridge-structuur. De deeltjesproductie in de double-ridge neemt gelijkmatig toe met de totale deeltjesproductie in de gebeurtenissen. Het verschillende gedrag van de deeltjesproductie in de double-ridge en in de mini-jets als functie van de totale deeltjesproductie is een sterke aanwijzing dat het onderliggende productiemechanisme voor beide effecten verschilt. Deze bevindingen komen overeen met de verwachting voor modellen waarin de mini-jets geproduceerd worden in onafhankelijke parton-partonbotsingen, gevolgd door fragmentatie van elke mini-jet, die niet afhangt van bijvoorbeeld hoeveel mini-jets er in de gebeurtenis worden geproduceerd.

De analyse gaat nog een stap verder door een schatting te maken van het aantal mini-jets per gebeurtenis, de zogenaamde ongecorreleerde bronnen. Het is aangetoond dat in het PYTHIA-model, het aantal mini-jets per gebeurtenis evenredig is met het aantal MPI. In dit proefschrift is voor het eerst het aantal mini-jets in p-Pb-botsingen gemeten. Het aantal mini-jets blijkt evenredig te zijn met de totale deeltjesopbrengst in de gebeurtenis, behalve bij zeer lage deeltjesopbrengst, waar ook de kans op meerdere parton-interacties zeer laag is. De interpretatie van dit resultaat is dat gebeurtenissen met groot aantal geproduceerde deeltjes niet worden gekenmerkt door mini-jets die uit meer deeltjes bestaan, maar door een groter aantal mini-jets. Dit wijst erop dat de totale deeltjesopbrengst in p-Pb-botsingen wordt bepaald door het aantal botsingen tussen partons van de inkomende protons

en loodkernen.

Verder is nagegaan of het aantal ongecorreleerde bronnen proportioneel is met het aantal binaire hadron-hadronbotsingen N_{coll} . Deze proportionaliteit blijkt te gelden voor gebeurtenissen over het middenbereik in totale deeltjesproductie. In gebeurtenissen met een erg hoge (lage) multipliciteit is het aantal ongecorreleerde bronnen hoger (lager) dan de lineaire trend. Een soortgelijk effect is gevonden voor de afhankelijkheid van het aantal ongecorreleerde bronnen van het aantal MPI van Monte-Carlo. Dit is een verdere bevestiging van het idee dat het aantal ongecorreleerde bronnen een goede schatting is voor het aantal partonbotsingen in een hadronbotsing.

Al met al is in dit proefschrift een methode toegepast om het gedrag van MPI en mini-jetproductie in zachte productieprocessen, dat wil zeggen met lage impuls, te bestuderen, waar de experimentele resultaten van cruciaal belang zijn om theoretische modellen te valideren. De metingen in dit proefschrift leveren veel informatie op voor het modelleren van p-Pb-botsingen: de modellen moeten in staat zijn de incoherente superpositie van mini-jets te beschrijven, en tegelijkertijd ook de double-ridge-structuren beschrijven en de toename van de gemiddelde p_T met de totale deeltjesproductie in de gebeurtenissen. Daarmee draagt dit werk bij aan het begrijpen van de recent ontdekte intrigerende effecten in botsingen van protonen met elkaar en met loodkernen.

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