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Search for jet quenching effects in high-multiplicity pp collisions with ALICE

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Declaration

I hereby declare I have written this doctoral thesis independently and quoted all the sources of information used in accordance with methodological instructions on ethical principles for writing an academic thesis. Moreover, I state that this thesis has neither been submitted nor accepted for any other degree.

In Prague, April 2024

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Abstrakt

V koncovém stavu malých srážkových systémů, které vykazují velkou multiplicitu produkovaných částic, byly pozorovány kolektivní jevy, které se podobají dobře zavedeným projevům kolektivity vznikajícím ve srážkách těžkých iontů, kde vzniká kvarkové–gluonové plazma (QGP). U těchto malých systémů však nebylo doposud potvrzeno, že by docházelo k efektu zhášení jetu. Určení velikosti energetických ztrát partonů v těchto systémech nebo alespoň stanovení limitů na jejich velikost je tak zásadní otázkou důležitou pro pochopení podmínek vzniku QGP z hlediska objemu vytvářeného média. Tato práce pojednává o výsledcích hledání efektu zhášení jetů v proton–protonových srážkách s velkou multiplicitou produkovaných částic v koncovém stavu při těžištové energii $\sqrt{s} = 13$ TeV na základě dat nabraných detektorem ALICE v letech 2016–2018. Množství nabitých částic produkovaných ve srážkách je charakterizováno pomocí amplitudy signálu z dopředného scintilačního detektoru V0. V nabraných datech zkoumáme semi-inkluzivní rozdělení azimutálního úhlu sevřeného mezi zvoleným hadronem s velkou příčnou hybností ($p_{T, \text{trig}}$) a jety. Fyzikální motivací pro naše měření je hledat, zda ve srážkách s velkou multiplicitou částic dochází k vychylování partonu vzhledem k jeho původnímu směru, které by mohlo být zapříčiněno interakcí partonu s horkým a hustým médiem. Takovéto interakce by měly za následek rozšíření měřého úhlového rozdělení ve srovnání s vakuem. Rekonstrukce jetů se provádí pomocí algoritmu anti- k_T s parametrem rozlišení $R = 0,4$. Analýza provádí odečet příspěvku od nekorelovaných párů hadron–jet, včetně příspěvků z multipartonových interakcí, pomocí metody založené na datech. Naměřená distribuce úhlového rozdělení pro srážky s velkou multiplicitou vykazuje potlačení výtěžku produkovaných jetů a jeví se též širší ve srovnání s rozdělením získaným ve srážkách triggerovaných jako minimum bias. Tyto rysy jsou také vidět v událostech pp simulovaných Monte Carlo generátorem PYTHIA 8, který energetické ztráty partonů vlivem zhášení jetů neuvažuje. Následné studium simulovaných dat z PYTHIA ukazuje, že pozorované potlačení a rozšíření vzniká vlivem triggeru vybírajícího srážky s vysokou multiplicitou částic v detektoru V0, který mimojiné dává pozitivní signál také v případech, kdy detektor V0 zasáhl jet. Učiněný závěr má ovšem širší platnost a lze jej vztáhnout na všechna měření, kde se pátrá po zhášení jetů v malých srážkových systémech s velkou multiplicitou částic v koncovém stavu a kde je výběr takových srážek založen na nějakém podobném principu.

Abstract

Several unforeseen collective phenomena have been observed in high-multiplicity (HM) final states of small collision systems that resemble the well-established signatures of the quark–gluon plasma (QGP) formation in heavy-ion collisions. However, jet quenching has not been observed in small collision systems. Quantification or setting limits on the magnitude of jet quenching in small systems is essential for understanding the limits of the QGP formation. This thesis discusses the outcomes of a search for jet quenching effects in HM proton–proton collisions at $\sqrt{s} = 13$ TeV collected by the ALICE detector. The multiplicity of charged particles in an event is characterized by the signal amplitude from the V0 forward scintillator detector. The analysis measures the semi-inclusive azimuthal acoplanarity distribution of charged-particle jets recoiling from a trigger hadron with high transverse momentum ($p_{T, \text{trig}}$). Propagating through the hot and dense medium, hard scattered parton is deflected relative to its original direction, thus resulting in the broadening of the hadron-jet acoplanarity compared to a vacuum. Jet reconstruction is performed utilizing the anti- k_T algorithm with cone radius $R = 0.4$. The analysis performs a data-driven subtraction of contributions from uncorrelated trigger–recoil jet pairs, including contributions from multiple partonic interactions. The measured acoplanarity distribution for HM events exhibits suppression and broadening compared to minimum-bias events. These peculiar features are also seen in pp events simulated by the PYTHIA 8 Monte Carlo event generator, which does not incorporate jet quenching. Further studies of PYTHIA simulations reveal that the observed suppression and broadening originate from a bias induced by the ALICE high-multiplicity trigger. The identification of this bias has implications for all jet quenching searches in small collision systems where selection is made on the event activity.

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List of Acronyms

- APTS** Analogue pixel test structure. xvi, 92–95, 97
- DPTS** Digital pixel test structure. xvi, 92
- FF** Fragmentation function. 8
- HM** High multiplicity. xii–xv, 2, 29, 30, 33, 40–49, 53–61, 72–77
- IRC** Infrared and collinear safe. 11, 12
- ITS** Inner Tracking System. ix, xv, 2, 3, 35, 36, 80–85
- LHC** Large Hadron Collider. 1, 2, 11, 17–20, 22, 23, 25, 27, 28, 33, 35, 36, 38
- LO** Leading order pQCD. xiv, 8, 72, 73, 75, 77
- MAMI** Mainz Microtron. xvi, 94, 98, 99
- MAPS** Monolithic Active Pixel Sensor. 87, 88
- MB** Minimum bias. xii–xv, 2, 30, 40–51, 53–61, 72–76
- MOSS** Monolithic Stitched Sensor. xvi, xvii, 99–103
- MPI** Multiple parton interaction. 47
- NLO** Next-to-leading order pQCD. xiv, 11, 73
- nPDF** Nuclear parton distribution function. 14, 15
- PDF** Parton distribution function. 8, 14
- pQCD** Perturbative quantum chromodynamics. x, xiv, 4, 6–8, 14, 30, 72, 73
- PS** Proton Synchrotron. xvi, 101, 102
- QC** Quality Control. 83–85
- QCD** Quantum chromodynamics. x, 3–7, 9, 11, 27, 35, 77
- QGP** Quark–gluon plasma. 1–3, 14–16, 23, 28, 30, 33, 77
- RHIC** Relativistic Heavy Ion Collider. 1, 17, 18, 20, 22, 23, 33

SPD Silicon Pixel Detector. 36, 44, 45

SPS Super Proton Synchrotron. xvi, 23, 94, 98, 99

TPC Time Projection Chamber. 35–37, 44

TT Trigger track. 29–32, 47, 77

UE Underlying event. 13, 14

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Chapter 1

Introduction

The hadronic matter at extreme temperatures and energy densities forms the hot and dense, strongly interacting quark–gluon plasma (QGP), a state in which quarks and gluons are deconfined [1]. In the Big Bang model, the Universe was filled with the QGP a few microseconds after its birth [2]. At present, the QGP is created in heavy-ion collisions at the Relativistic Heavy Ion Collider (RHIC) at Brookhaven National Laboratory and the Large Hadron Collider (LHC) at CERN. Production of the QGP in laboratory conditions provides a unique opportunity to study its properties. Comparison of the experimental results with theoretical calculations based on viscous relativistic hydrodynamics suggests that the QGP is opaque and almost perfect liquid with a noticeably small shear viscosity to entropy density ratio [3].

Experimental studies of the QGP have many open questions [4]. One of them concerns the limits of the QGP formation in terms of the size of colliding nuclei. To study this question, measurements have been carried out for *small collision systems*, where the projectile is a proton or a light nucleus and the target is a proton or a nucleus. Measurements in small collision systems exhibit experimental signatures, which are associated with the QGP creation in heavy-ion collisions [5], such as collective flow [6]–[13] and enhanced production of multi-strange hadrons [14]. However, these observations do not provide reliable evidence for the QGP formation in small collision systems. Collective flow in small systems is not unique to hydrodynamic description [15]. Strangeness enhancement can also be qualitatively described by the rope hadronization model [16]. Under these circumstances, it is crucial to search for other prominent signatures of the QGP formation.

Jets are the direct consequence of hard scatterings (high momentum transfer Q^2) of partons from hadronic projectiles. Their production is well understood in elementary collisions, where theoretical calculations based on perturbative quantum chromodynamics (pQCD) accurately describe data [17]–[19]. In heavy-ion collisions, hard-scattered parton is produced before the QGP is formed. Interaction of hard-scattered parton with the

medium results in modification of jet production rates and jet sub-structure, so-called jet quenching [20], [21]. Quenched jets retain information about their interaction with the medium. Thus, they serve as sensitive probes of the QGP.

Jet quenching is a requisite consequence of the QGP formation. In small collision systems, theoretical calculations predict that the jet quenching signal is small [22]–[26]. One of the manifestations of jet quenching is the suppression of inclusive high- p_T hadron or jet yield when compared to the corresponding yield measured in pp collisions, which is scaled by the average number of binary nucleon–nucleon collisions calculated using the Glauber model [20], [21], [27]. Yield suppression of inclusive jets was measured in d–Au collisions at RHIC [28] and in p–Pb collisions at the LHC [29] for different intervals of event activity (EA) characterized based on forward multiplicity and transverse energy E_T , respectively. However, Glauber modelling in these kinds of measurements is subjected to substantial non-geometric bias [30]–[32]. An alternative characterization of EA based on energy from the zero-degree calorimeter (ZDC) was shown to be less biased, though the relation of the ZDC signal to the collision geometry is model dependent [33]. These biases and uncertainties related to Glauber modelling constrain the sensitivity of inclusive observables to jet quenching in small systems. At present, there is no persuasive evidence of jet quenching in small systems from inclusive yield measurements of hadrons/jets within experimental uncertainties [29].

This thesis presents a novel approach to search for jet quenching effects in high-multiplicity (HM) pp collisions at $\sqrt{s} = 13$ TeV. Inclusive observables are not applicable in such measurements since Glauber modelling is undefined for HM pp collisions. Rather, we investigate the semi-inclusive hadron-jet acoplanarity distribution [34]–[36]. It is expected that the acoplanarity distribution gets additional broadening when compared to its vacuum shape due to in-medium multiple scatterings [37], [38]. The acoplanarity measurements are based on the Δ_{recoil} observable [34], which provides a purely data-driven way for suppression of uncorrelated background yield, including yield from multi-parton interactions that would otherwise obscure the acoplanarity broadening from jet quenching. The acoplanarity distribution measured in minimum-bias (MB) events is utilized as a reference for the search for medium-induced acoplanarity in HM events.

The goals addressed in this Doctoral thesis are as follows:

1. Perform the hadron–jet acoplanarity measurements in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV recorded by the ALICE experiment;
2. Contribute to the development of the *Quality Control* software utilized for the quality assessment of data collected by the upgraded Inner Tracking System (ITS2) of the ALICE experiment;

3. Analyze data from test beam campaigns carried out with silicon sensor prototypes for the future ITS3 upgrade.

The thesis is organized as follows: Section 2 provides a concise introduction to quantum chromodynamics (QCD) theory and jet quenching. Further, it discusses the QGP-like signatures measured in small collision systems; Section 4 briefly presents the parts of the ALICE detector, which are of particular importance for this analysis; Section 5 provides detailed description of the analysis steps, including event selection, jet reconstruction, data unfolding, and estimation of systematic uncertainties; Section 6 presents and discusses the final results; Sections 8 and 9 are related to the upgrade projects of the Inner Tracking System of the ALICE detector, where my contribution is discussed. The thesis is concluded with a summary of the achieved results.

Statement of author's contribution

My contribution to the material presented in this thesis can be summarized as follows:

1. I carried out the analysis of hadron–jet acoplanarity in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV studying the semi-inclusive Δ_{recoil} observable. I performed all stages of the analysis, including processing of raw data, data correction via unfolding, assessment of systematic uncertainties, and interpretation of the final results. Also, I performed standalone PYTHIA 8 simulations shown in Chap. 6. I am a principal author of the corresponding paper released by the ALICE Collaboration [39];
2. I developed several ITS Quality Control tasks (asynchronous Cluster and Track tasks, post-processing Cluster and Front-End electronics tasks);
3. I analysed test beam data from Mainz Microtron and CERN Proton Synchrotron to characterize sensor prototypes for the ITS upgrade beyond LHC Run 3.

Besides that, I was also involved in radiation hardness tests of sensor prototypes for the ITS3. The tests were carried out at the cyclotron of the Nuclear Physics Institute of the Czech Academy of Sciences.

The list of publications and the list of conferences are presented in Appendix E and F.

Chapter 2

State of the art

2.1 Quantum chromodynamics

Quantum chromodynamics (QCD) [40] is a gauge field theory which describes interactions between strongly interacting particles, quarks with a half-integer spin and massless gluons having a spin equal to one. The dynamics of the strong interaction is governed by the QCD Lagrangian,

$$\mathcal{L}_{\text{QCD}}(x) = -\frac{1}{4}\vec{F}_{\mu\nu}^a(x)\vec{F}_a^{\mu\nu}(x) + \sum_{f=1}^{n_f} \bar{\psi}_i^f(x) [i\gamma^\mu \partial_\mu \delta^{ij} + g_s \gamma^\mu A_\mu^a t_a^{ij} - m_f \delta^{ij}] \psi_j^f(x), \quad (2.1)$$

which is invariant under local transformations of the colour degrees of freedom. The gauge symmetry group of QCD is the non-Abelian Lie group $\text{SU}(3)$. The indices $i, j = (1, 2, 3)$ and $a = (1, \dots, 8)$ in (2.1) denote color and adjoint color indices, respectively. Lorentz indices are represented by Greek letters and the Einstein summation convention is adopted. The summation goes through all quark flavours. The γ_μ and δ^{ij} are the Dirac matrices and the Kronecker symbol, respectively. The $\psi_j^f(x)$ and $\bar{\psi}_i^f(x)$ represent a quark and antiquark spinor field operator for quark flavor f and mass m_f , with Dirac adjoint $\bar{\psi} = \psi^\dagger \gamma^0$. The coupling constant g_s is the only one free parameter of the Lagrangian (2.1), if one assumes that quark masses are fixed, and it plays a fundamental role in formulating the perturbative QCD (pQCD) framework. Finally, the gluon field tensor $F_{\mu\nu}^a$ is given by

$$F_{\mu\nu}^a(x) = \partial_\mu A_\nu^a(x) - \partial_\nu A_\mu^a(x) + g_s f_{abc} A_\mu^b(x) A_\nu^c(x), \quad [t_a, t_b] = i f_{abc} t^c, \quad (2.2)$$

where ∂_μ is a covariant derivative, A_μ^a is a gluon field, f_{abc} and t_a are the structure constants and the generators of the $\text{SU}(3)$ group, respectively. The bilinear term in (2.2) reflects the non-Abelian character of the theory, thus QCD, in contrast to quantum electrodynamics

(QED), contains a self-interaction of three and four gluons that is depicted in terms of the Feynman diagrams in Fig. 2.1.



Figure 2.1: Feynman diagrams for gluon self-interaction in QCD: a) 3-gluon vertex; b) 4-gluon vertex. Taken from Ref. [41].

2.1.1 Running coupling constant

In order to calculate the cross section of elementary QCD processes, one needs to utilize the perturbative theory, which is applicable only under the condition $\alpha_s = g_s^2/4\pi < 1$ [42]. In QCD, however, α_s changes its value according to the selected renormalization scale μ_R [43], i.e. the QCD coupling constant is running. The expression for α_s can be obtained from the renormalization group equation [43], setting the physical scale Q^2 equal to the renormalization scale μ_R^2 ,

$$\alpha_s = \frac{4\pi}{\beta_0 \ln(Q^2/\Lambda^2)}, \quad \beta_0 = 11 - \frac{2}{3}n_f, \quad (2.3)$$

where β_0 is a 1-loop β -function coefficient, which depends on the number of quark flavors n_f accessible at the scale Q^2 . QCD introduces a scale parameter Λ , which is renormalization-scheme dependent and which sets a limit $Q^2 \rightarrow \Lambda^2$, when the perturbation theory no longer applies.

The effect of running coupling was verified experimentally for a variety of processes. A summary plot in Fig. 2.2 compares theory and experimental data. Good agreement between theory predictions and experimental data suggests that when Q^2 tends to infinity in a hard scattering process, the coupling constant $\alpha_s \rightarrow 0$. Thus, the interaction between colored particles is negligible at short distances. This peculiar property of non-Abelian field theories is known as *asymptotic freedom*. On the other hand, as the momentum Q^2 decreases, the strong interaction becomes indeed strong, preventing quarks and gluons from being in a free state. Therefore, solely colorless bound states of hadrons are physically observable. This remarkable property of QCD is known as *color confinement*. Since color confinement involves non-perturbative processes, the state-of-the-art QCD theory cannot describe it from the first principles.

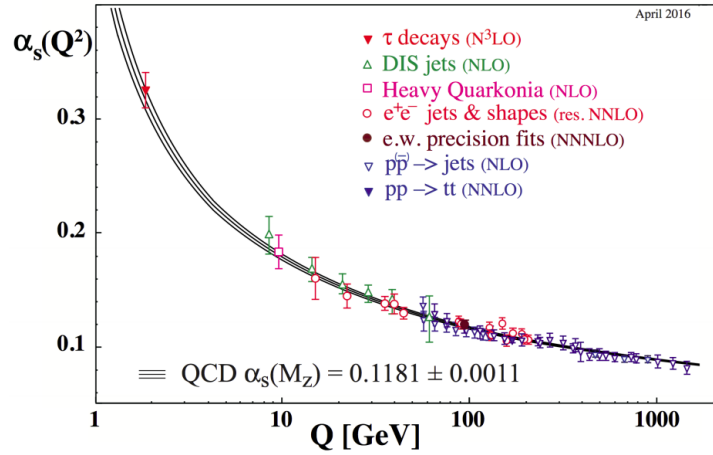


Figure 2.2: Summary of measurements of the strong running coupling constant α_s as a function of transferred momentum Q . Experimental data are represented by a bunch of markers corresponding to different processes where pQCD is involved. The solid line and its associated uncertainty band correspond to the prediction of pQCD theory. Taken from Ref. [44].

2.1.2 Quark–gluon plasma

The confinement property is inherent to QCD at normal temperatures and nuclear matter densities. Collisions of ultrarelativistic heavy ions can, however, create such high energy densities and temperatures that the QCD matter can undergo a phase transition into a new state, where quarks and gluons are deconfined, the so-called quark–gluon plasma (QGP) [45]. The Hot Big Bang model suggests that the QGP existed in the early Universe about $\sim 10^{-6}$ s after the Big Bang [46]. Therefore, heavy-ion collisions provide a unique opportunity to study the physics of the early Universe in laboratory conditions.

Figure 2.3 shows a sketch of the QCD-matter phase diagram expressed in terms of temperature T and baryon chemical potential μ_B . The origin of the coordinate system in the phase diagram at $(T = 0, \mu_B = 0)$ corresponds to the QCD vacuum state. The ordinary nuclear matter with density $\rho_0 \approx 0.15 \text{ fm}^{-3}$ is located on the phase diagram at $T \approx 0 \text{ MeV}$ and $\mu_B \approx 938 \text{ MeV}$. When the nuclear matter is heated up, nucleons evaporate from nuclei and form the hadron–gas phase. Further heating of the hadron gas leads to the production of other hadron species. When the energy density reaches about 1 GeV/fm^3 , a phase transition to the QGP state occurs [1]. The lattice calculations at $\mu_B \approx 0 \text{ MeV}$ for realistic masses of u, d , and s quarks suggest that a phase transition between the QGP and the hadron gas proceeds as a rapid crossover [48]. Note that due to the fermionic-sign problem, the lattice QCD can provide calculations only for the region near the temperature axis where μ_B is close to 0 [49]. At finite values of μ_B , effective field models based on the chiral perturbation theory [50] are used. These models predict that the crossover should turn into a first-order phase transition at some critical temperature

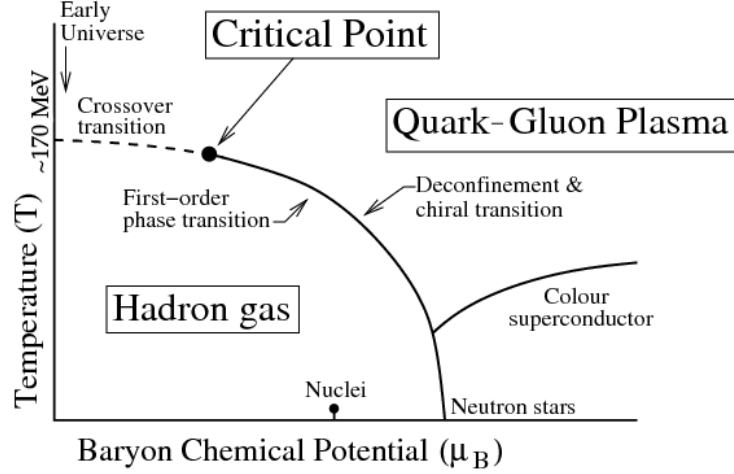


Figure 2.3: Schematic phase diagram of QCD matter as a function of temperature T and baryon chemical potential μ_B . The arrow “Early Universe” depicts the evolution path of the early Universe. Further details can be found in the text. Taken from Ref. [47].

T_c . The point on the phase diagram, at which this transition is expected to occur, is called the critical point. It is important to emphasize that model calculations only predict the presence of the critical point, however, none has succeeded in restricting its location.

The phase diagram in Fig. 2.3 also shows other expected phases of the QCD matter, which are formed when the nuclear matter is enormously compressed at $T \approx 0$. In this regime, the hadrons start overlapping, thus quarks and gluons become weakly coupled. This state of matter is called cold quark matter. Presumably, such extreme conditions can be reached at the cores of superdense neutron stars [51].

2.1.3 QCD factorization theorem

The cross section of elementary partonic processes, like $q + q, g + g$, etc., can be calculated within the pQCD framework. When calculating the scattering of a more complex hadron–hadron system, one solely relies on the QCD factorization theorem [52]. The theorem assumes that the cross section of a hadron–hadron scattering can be expressed as a convolution of a non-perturbative part, which describes the incoming and outgoing states of a hadron at time scales well separated from the scattering process, with a perturbative hard-scattering partonic cross section describing an elementary QCD process. Schematically, it can be written as

$$\sigma^{p+p \rightarrow h+X} \propto \sum_{a,b,c} \hat{\sigma}^{a+b \rightarrow c+X} \otimes f_{a/p}(x_a, Q^2) \otimes f_{b/p}(x_b, Q^2) \otimes D_{h/c}(z, Q^2), \quad (2.4)$$

where $\hat{\sigma}^{a+b \rightarrow c+X}$ is a parton-level cross section for scattering of two partons a and b leading to inclusive production of a parton c and a state X permitted by the conservation laws.

The non-perturbative contribution is encoded in the parton distribution function (PDF) $f_{a/p}(x_a, Q^2)$ [53] and in the fragmentation function (FF) $D_{h/c}(z, Q^2)$ [54]. At leading order (LO), the PDF gives a probability to observe a parton a with a momentum fraction x_a of the parent hadron p at the scale Q^2 , whereas the probability to find a hadron h which emerges from the parton c and carries a fraction z of parton's momentum at the scale Q^2 is governed by the FF.

A crucial consequence of the factorization theorem is the universality of the PDFs and FFs, i.e. the non-perturbative part of the cross section (2.4) is process independent. Therefore, the PDFs and FFs can be measured in different processes and at different scales Q^2 [55]–[58].

2.1.4 Parton shower

High-energy hadron–hadron collisions inevitably involve the formation and evolution of parton showers. Simplistically, a parton shower can be viewed as a cascade of successive splittings of “parent” partons into two “daughter” partons. Historically, the showers are divided into initial- and final-state radiation, depending on the time scale of their formation compared to the occurrence of the hard-scattering process [59].

Insofar as gluons possess a color charge, a parton can branch into four configurations, specifically $g \rightarrow q\bar{q}$, $q \rightarrow gq$, $q \rightarrow qg$, and $g \rightarrow gg$. Each splitting channel is characterized by the splitting kernel $P_{a \rightarrow bc}(z)$ [60], which gives the probability that a “parent” parton a splits into two “daughter” partons b and c . The variable z represents the energy sharing, with the “daughter” b taking away a fraction z and the “daughter” c carrying the remaining fraction $1 - z$ of the “parent” energy. At LO of pQCD for massless quarks, only the first two splitting channels do not contain a singularity in their kernels $P_{g \rightarrow q\bar{q}}(z)$ and $P_{q \rightarrow gq}(z)$, whereas the $P_{q \rightarrow qg}(z)$ and $P_{g \rightarrow gg}(z)$ kernels diverge in the limit of soft gluon emission for $z \rightarrow 0$.

Simulations of a parton shower in popular Monte Carlo event generators are based on the use of the Sudakov form factors [60],

$$\Delta_a(Q_{\max}^2, Q_a^2) = \prod_{i \in \text{splitting channel}} \exp \left[- \int_{Q^2}^{Q_{\max}^2} \frac{dQ^2}{Q^2} \frac{\alpha_s(Q^2)}{2\pi} \int_{z_{\min}}^{z_{\max}} dz P_{a \rightarrow i}(z, Q^2) \right], \quad (2.5)$$

that represent the unconditional probability for a parton a not to split between the scales $Q_{\max}^2$ and Q^2 , where the former is the maximum possible virtuality of a parton a . Choice of $Q_{\max}^2$ is ambiguous. It depends on model assumptions or is considered as a parameter to fit hadron/jet spectra in proton–proton collisions [61]. Further details about implementing parton showers in Monte Carlo generators can be found in Ref. [60].

In literature, variables Q^2 and z are called evolution and splitting variables, respectively. They can have different definitions depending on a parton shower model implemented in a given event generator. For instance, the “simple” parton-shower model, implemented in the PYTHIA event generator [62], [63], is a p_T -ordered shower algorithm, i.e. parton transverse momentum p_T is set as the evolution variable. The splitting variable has in this case a meaning of energy fraction. For most $2 \rightarrow 2$ hard QCD processes in the PYTHIA, the maximum $p_T^{\max}$, which replaces $Q_{\max}^2$ in this case, is set to be $p_T^{\max} = 4p_T^{\text{hard}}$, where p_T^{hard} is a transverse momentum of hard scattering [59].

The PYTHIA 8 simulations with the Monash tune [64] that were carried out in this thesis employ the “simple” parton-shower model. As of version 8.3, the PYTHIA is equipped with three different complete parton-shower models, see the recent PYTHIA manual for further details [62].

2.1.5 Lund fragmentation model

The transition of colored partons to colorless hadron states is hidden from our understanding by the nonperturbative regime of QCD. The impossibility of applying perturbative methods has led to the emergence of a variety of phenomenological models that describe the process of hadronization based on various physical assumptions. For example, the event generators Sherpa [65] and HERWIG [66] consider hadronization as a process of formation of clusters from color-connected quarks with subsequent decay of the clusters into hadrons. At the same time, in the PYTHIA, the hadronization process is based on the Lund string model [67], which is considered below.

The Lund fragmentation model is based on the idea that the color field between a $q\bar{q}$ pair can be represented as a color string connecting the quarks. The corresponding potential energy of two quarks separated by a distance r can be described by the so-called Cornell potential [68],

$$V_{\text{QCD}}(r) \approx -\frac{4\alpha_s}{3r} + \sigma r, \quad (2.6)$$

where the first part represents the Coulomb-like potential. The σ coefficient is called the string tension. As the quarks move away from each other, the potential energy grows linearly with the distance, which gives a qualitative explanation for the confinement effect. The presence of the linear term was confirmed by lattice QCD calculations [69] and emergence of the Regge trajectories [70], which constrained the value of σ to be about 1 GeV/fm. At large distances, the linear term dominates, and only this term is significant for the Lund model.

The string model describes a meson state with $p_T = 0$ as a massless $q\bar{q}$ -pair. When the $q\bar{q}$ -pair is created, the quarks move in opposite directions with the speed of light.

While moving apart from each other, the quarks stretch a color string, thereby increasing its potential energy. If the $q\bar{q}$ -pair does not have enough energy to break the string, the quarks will move towards each other. This periodic movement, called a “yo-yo” motion, will be repeated with period $\tau = 4r_0/c$, where $2r_0$ is the maximum separation distance. The invariant mass of the system is equal to the maximum potential energy of the stretched string, namely $m = 2\sigma r_0$.

The string tension does not grow indefinitely. At a certain moment, it is energetically beneficial for the system to break the string, while creating a new $q\bar{q}$ -pair at the broken ends to compensate for the color charge. In general, the initial color string can break up into several mesons. Each newly created meson carries a fraction z of the light-cone momentum, which is given by the Lund symmetric fragmentation function [67],

$$f(z) \propto \frac{(1-z)^a}{z} \exp\left(-b\frac{m^2}{z}\right), \quad (2.7)$$

where a and b are parameters of the model. The sketch of the Lund string breakup is shown in Fig. 2.4. A space-time point, where the string breaks, is called vertex and it is marked as a black dot. Due to local flavor conservation, a $q\bar{q}$ -pair should be created in the same vertex. Closed rectangular regions depict the “yo-yo” motion of $q\bar{q}$ -pairs. The area of each of the regions is proportional to a mass of the corresponding meson state.

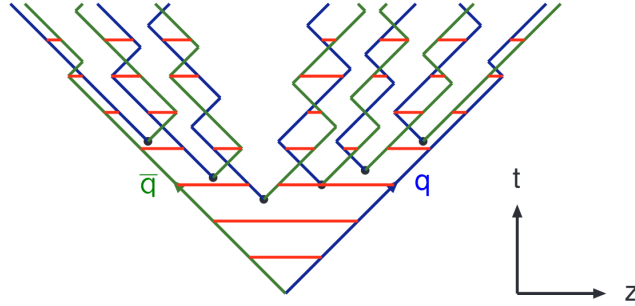


Figure 2.4: Breakup of the Lund string between a $q\bar{q}$ -pair in space-time. Taken from Ref. [71].

The discussed fragmentation model becomes more complicated when considering massive quarks with a non-zero transverse momentum. In this case, quarks are produced via the Schwinger mechanism [72]. They appear on the mass-shell after tunneling through the distance $(m + p_T)/\sigma$. The consequence of the tunneling process is the appearance of the Gaussian factor in the probability function,

$$\mathcal{P}_{q\bar{q}} \propto \exp\left(-\pi \frac{m_q^2 + p_T^2}{\sigma}\right), \quad (2.8)$$

which “determines” a flavour of the created $q\bar{q}$ -pair. The relative probability of produc-

ing a $c\bar{c}$ -pair during the string breaking is 10^{-11} when compared to light quarks [73]. Therefore, c and b quarks are produced only in a hard-scattering process.

The addition of gluons and baryons to the Lund model further substantially complicates the theory, therefore it is not considered here. An elaborate discussion of the Lund fragmentation model is available in Refs. [62], [67].

2.2 Jets

A hard scattered parton forms a collimated parton shower. The evolution of the shower continues until the parton virtuality reaches the non-perturbative scale Λ , after which the partons become confined inside colorless hadrons. The resulting collimated shower of energetic hadrons is called a jet [74]. Jets play a crucial role in understanding the hard-scattering QCD processes. The large production rate of jets at the LHC energies provides high-precision jet data in a wide range of transverse momenta and rapidities, which are used in the global-fit analyses of parton distribution functions [75] and for extraction of the QCD running coupling constant [76].

2.2.1 Jet reconstruction algorithm

The above-mentioned description of a jet provides only a qualitative physics picture without specifying any rules for identifying and handling this object at the partonic and hadronic levels. Therefore, the definition of a jet based on a jet clustering algorithm is introduced. This definition establishes a set of rules for the jet reconstruction from a list of input particles. A properly designed jet algorithm must ensure finite jet production cross sections at any order of perturbation theory. Starting from the next-to-leading order (NLO), the emission of soft and collinear partons occurs. In order to establish a well-defined connection between the definitions of jets at partonic and hadronic levels, a jet reconstruction algorithm must be infrared and collinear (IRC) safe [77]. In other words, the IRC safe jet reconstruction algorithm ensures that the number of reconstructed hard jets is unaffected by infrared emission and collinear splitting processes. An example of the IRC safe algorithm is the generalized k_T algorithm, which belongs to the family of sequential recombination algorithms [74]. The algorithm commences the jet reconstruction from a list of pre-clusters, which can be formed from partons, detected particles or calorimeter cells. Further, the algorithm evaluates which particles are the closest ones to each other based on a specified definition of a “proximity measure”:

- For each pair of particles (i, j) one defines two sets of distances:

$$d_{ij} = \min(p_{T_i}^{2p}, p_{T_j}^{2p}) \frac{\Delta R_{ij}^2}{R^2} \quad \text{and} \quad d_{iB} = p_{T_i}^{2p}, \quad (2.9)$$

where the notation “ $\min(\dots)$ ” means that from the two input values one chooses the lower one, R is a jet resolution parameter, sometimes called also as a jet radius, and ΔR_{ij}^2 is defined as

$$\Delta R_{ij}^2 = (y_i - y_j)^2 + (\varphi_i - \varphi_j)^2, \quad (2.10)$$

where y and φ are particle rapidity and azimuthal angle, respectively;

- If $d_{ij} < d_{iB}$ then particles (i, j) are merged into a new particle, which is afterwards returned back to the list of pre-clusters. Otherwise, the particle “ i ” is declared as a final jet and removed from the list.

The parameter p is equal to -1 in the case of the anti- k_T algorithm [78]. There are two other IRC-safe algorithms with a different choice of p : $p = 0$ corresponds to the Cambridge/Aachen algorithm [79], and $p = 1$ corresponds to the k_T algorithm [80]. It can be seen that the anti- k_T algorithm initializes jet reconstruction by merging nearby high- p_T particles. The merging of soft particles with hard ones happens before the clusterization of soft particles among themselves. Thus, the reconstructed anti- k_T jets are resilient to a soft background. Furthermore, the anti- k_T jets have a perfect conical shape if there is no hard particle within a distance of $2R$ [78]. On the contrary, the k_T algorithm first recombines low- p_T particles that are close to each other in the $(y - \varphi)$ plane. This results in the algorithm being sensitive to a soft background, like the underlying or pile-up events. The background presence causes the back reaction [81] of the algorithm, which leads to the energy and shape alteration of the reconstructed jets. This feature of the k_T algorithm is nevertheless utilized in the background estimations in hadron–hadron and heavy-ion collisions.

The jet algorithm needs to be supplemented by a specification of a recombination scheme that determines the procedure of combining the 4-momenta of merged particles. The simplest recombination scheme is the E -scheme, which sums particles’ 4-momenta. In addition, the boost-invariant p_T recombination scheme is broadly applicable for pp collisions. Further information about the existing recombination schemes can be found in Ref. [82].

2.2.2 Background subtraction

Jets are a direct consequence of rare scattering processes with high- Q^2 momentum transfer. In heavy-ion collisions, hard-scattering processes are accompanied by an intense soft background, which originates from the underlying event (UE) (multiple parton interactions, initial- and final-state radiation, etc.). Under these conditions, the identification of jets originating from hard-scattering processes becomes a non-trivial task insofar as jet reconstruction algorithms cluster also background particles into jets. Moreover, the properties of reconstructed jets are largely affected by contamination from a soft background. The yield of a soft background must be suppressed to determine the kinematic properties of jets as precisely as possible.

An experiment-independent approach for background subtraction was proposed in Ref. [83]. This approach is based on the idea that the UE particles are distributed fairly uniformly in the $(y - \varphi)$ plane, at least after averaging over many events. Therefore, contamination of hard jets due to soft background should be, on average, proportional to jet area. However, the definition of jet area is rather complicated in itself. A widely used definition of jet area, called the “active area”, is based on the concept of ghost particles, which are massless particles with a negligible transverse momentum ($p_T \sim 10^{-100}$ GeV/ c) [81]. Ghosts with some density per unit area in the $(y - \varphi)$ plane are added into an event, after which a jet reconstruction algorithm is launched. It is important to stress that the “active area” definition can be used solely for infrared-safe algorithms, which are insensitive to the presence of soft particles like ghosts. Ghosts can either be clustered with a hard jet or clustered among themselves. Finally, the jet area is evaluated based on the number of ghost particles within a given jet, see Ref. [81] for further details. The “active area” approach is implemented in the FastJet package, where one can find other methods for jet area measurements [82]. Having a definition of jet area, one can move towards background subtraction from reconstructed jets. If the UE has a transverse momentum density per unit area ρ , then the reconstructed jet p_T is corrected using the following formula:

$$p_{T, \text{jet}}^{\text{corr}} = p_{T, \text{jet}}^{\text{raw}} - \rho \times A_{\text{jet}}, \quad (2.11)$$

where $p_{T, \text{jet}}^{\text{raw}}$ and A_{jet} are jet transverse momentum before correction and jet area, respectively. Formula (2.11) is applied to all jets reconstructed in a given event. The background density ρ is calculated on the event-by-event basis as the median value of the $p_{T, \text{jet}}^{\text{raw}}/A_{\text{jet}}$ distribution, $\rho = \text{median}(p_{T, \text{jet}}^{\text{raw}}/A_{\text{jet}})$. To estimate ρ , additional jet reconstruction is performed separately from the main analysis utilizing either the k_T algorithm or the Cambridge/Aachen algorithm since they are more susceptible to soft background than the anti- k_T algorithm [84]. Formula (2.11) does not take into account the effect of

local background fluctuations, which smear the jet p_T spectrum. The smearing due to local background fluctuations needs to be corrected later on by means of unfolding.

The average transverse momentum of particles created in a hard process is greater than that of background particles. This fact is often utilised to reduce the population of purely combinatorial jets by introducing a minimum p_T cutoff for a leading jet constituent. However, this approach leads to a bias in jet fragmentation. The impact of UE particles on the jet energy scale can also be reduced by considering smaller jet radii or higher energy jets since the contribution of background particles is proportional to jet area. Hard-scattering processes and the UE are independent, therefore, the relative contribution of the latter to jet energy decreases with jet energy. Interaction of a jet parton with the medium leads to the transfer of parton energy outside the jet cone [85]. Therefore, studies of narrow high-energy jets alone cannot provide a complete picture of the parton energy loss in the QGP. Studies of low-energy jets with large radii are also of great importance. Note that theoretical calculations of jets with a large cone radius are better defined than those of jets with a small radius, see, for instance, Refs. [86], [87].

There are several approaches to background subtraction. A description of the most popular approaches used by the largest collaborations in the heavy-ion community can be found in Refs. [35], [88]–[90]. This thesis will follow the approach to background subtraction customarily used by the ALICE Collaboration [91]. The method is based on the “jet area” approach described above. The jets are reconstructed in two stages. In the first stage, all accepted tracks in an event are clustered into jets using the k_T algorithm. The two leading k_T jets with the highest p_T are excluded from the analysis. The remaining k_T jets are used to calculate the ρ value. In the second stage, jets are reconstructed using the anti- k_T algorithm, where p_T of jets is corrected utilizing Eq. (2.11).

2.3 Jet quenching

Jet production is extensively studied in pp collisions, where pQCD computations show good agreement with experimental measurements [92], [93]. It has been proposed to utilize jets to probe the hot and dense medium produced in collisions of ultra-relativistic heavy ions [94]. Jets originate from hard scattering processes, which occur prior to the QGP formation. Therefore, they experience the complete dynamical evolution of the medium carrying crucial information about its properties.

The physical interpretation of jet measurements in relativistic heavy-ion collisions requires the separation of two distinct nuclear matter effects. First, the PDF $f_{a/A}(x_a, Q^2)$, measured in the nucleus (nPDF), differs from the PDF $f_{a/p}(x_a, Q^2)$ corresponding to a free proton. This phenomenon is referred to as the initial state cold nuclear matter effects [95].

The nPDFs are studied typically through the measurements of the nuclear modification ratio,

$$R_a^A(x_a, Q^2) = f_{a/A}(x_a, Q^2)/f_{a/p}(x_a, Q^2), \quad (2.12)$$

where a parton species is denoted by the subscript a . The ratio can be extracted by employing a global-fit method to the deep inelastic scattering data of proton/deuteron–nucleus collisions [96]–[98]. Figure 2.5 shows the averaged nuclear modification ratios for lead nucleus obtained by three independent analyses at the scale $Q^2 = 10 \text{ GeV}^2$, see Ref. [97] and references therein for further details. The left and middle plots correspond to the valence- and sea-quark modifications, respectively, while the right plot shows the data for gluon nuclear modification. The shaded areas and dashed lines represent uncertainty bands. Experimental data suggest that the nuclear modification ratio contains several particular effects, which are general for nPDFs:

- suppression of R_a^A below 1 at low $x < 0.1$ refers to the *shadowing* effect, which is followed by the *antishadowing* ($R_a^A > 1$) in the region $x \sim 0.1$;
- in the valence region corresponding to $x \sim 0.4$, the so-called *EMC effect* is observed.

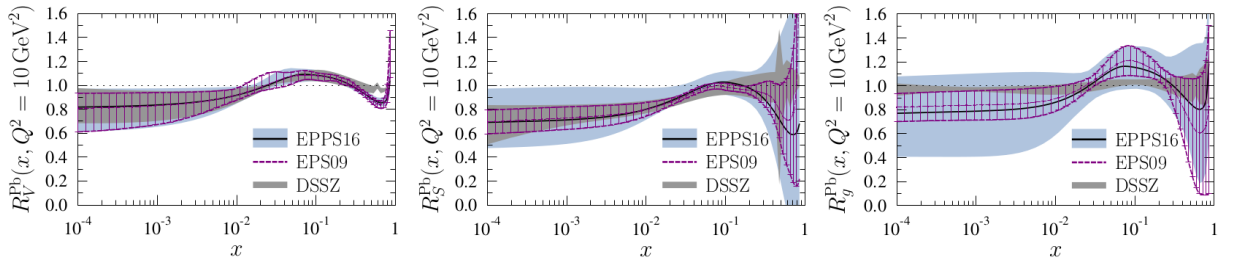


Figure 2.5: The averaged nuclear modification factors obtained for Pb nucleus as a function of x at scale $Q^2 = 10 \text{ GeV}^2$ for valence (left) and sea (middle) quarks, and also for gluon (right). Data are shown for the EPPS16, EPS09, and DSSZ analyses. The shaded areas and dashed lines represent the uncertainty bands. Taken from Ref. [97].

The second nuclear effect is associated with producing the hot and dense QGP medium. While propagating through the opaque medium, a jet parton interacts with the QGP constituents. Multiple elastic collisions and processes of medium-induced gluon radiation [99] result in parton energy loss, which has a strong dependence both on parton characteristics (e.g. energy and mass), and on the QGP properties (e.g. thickness, opacity, temperature, etc.). The energy redistribution within the parton shower leads to the yield suppression of high- p_T hadrons and jets, the so-called *jet quenching effect*.

Radiative energy loss is the dominant energy loss mechanism for light flavor partons at high energies. In the case of heavy quarks, the situation is somewhat modified. At low and intermediate energies, collisional energy loss exceeds the radiative one owing

to the large mass of heavy quarks that tends to reduce the phase space of medium-induced radiation [99]. When moving towards the high energy regime, the radiative loss commences dominating. However, the complete description of the parton energy loss is complicated by various non-trivial interaction mechanisms. Therefore, it is customary to resort to utilizing diverse formalisms for energy loss description. Despite the difference in approaches to the formulation of a particular formalism, three generally made assumptions can be distinguished [99], [100]:

- *Eikonal approximation*: energy E of a parton and energy ω of a radiated gluon are much greater than the transverse momentum $p_\perp$ exchanged with the QGP constituents, i.e. $E, \omega \gg p_\perp$;
- *Collinear approximation*: energy of the radiated gluon ω is much larger than its transverse momentum $k_\perp$, i.e. $\omega \gg k_\perp$;
- *Discrete scattering centers*: the mean free path of a parton propagating through the medium is greater than the Debye screening length.

Several formalisms have been developed to describe radiation energy loss, for instance, Baier-Dokshitzer-Mueller-Peigne-Schiff-Zakharov (BDMPS-Z) [101], [102], Arnold-Moore-Yaffe (AMY) [103], and higher twist (HT) [104]. They can also be extended by the addition of collisional energy loss. More details and discussions on the parton energy loss mechanisms are given in Refs. [99], [100].

2.3.1 Nuclear modification factor

The medium-induced modification of hadron or jet yield can be conveniently expressed in terms of nuclear modification factor R_{AA} [61],

$$R_{AA}(p_T) = \frac{1}{\langle T_{AA} \rangle} \frac{dN_{AA}/dp_T}{d\sigma_{pp}^{\text{INEL}}/dp_T}, \quad (2.13)$$

where N_{AA} is the normalized per event yield of hadrons or jets produced in AA collisions, and $\sigma_{pp}^{\text{INEL}}$ is the inelastic cross section of pp collisions. The nuclear overlap function is determined from Glauber modelling and it is defined as $\langle T_{AA} \rangle = \langle N_{\text{coll}} \rangle / \sigma_{NN}^{\text{INEL}}$, where $\langle N_{\text{coll}} \rangle$ is the average number of binary collisions determined from the Glauber modelling [27], and $\sigma_{NN}^{\text{INEL}}$ is inelastic nucleon–nucleon collision cross section. In the absence of nuclear effects, the R_{AA} is 1, i.e. particle production in AA collisions can be described by the corresponding number of independent pp collisions.

Figure 2.6 presents the measurements of nuclear modification factor in central Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. One can see that the production of inclusive charged high-

p_T hadrons is suppressed in central Pb–Pb collisions, but at the same time, there is no visible suppression in the production of heavy weak bosons and direct photons. The latter measurements suggest that the Glauber model correctly describes the collision geometry. The heavy bosons emerge in the initial hard-scattering processes. However, their leptonic decay products are not affected by the produced medium, therefore, their $R_{\text{PbPb}} \approx 1$.

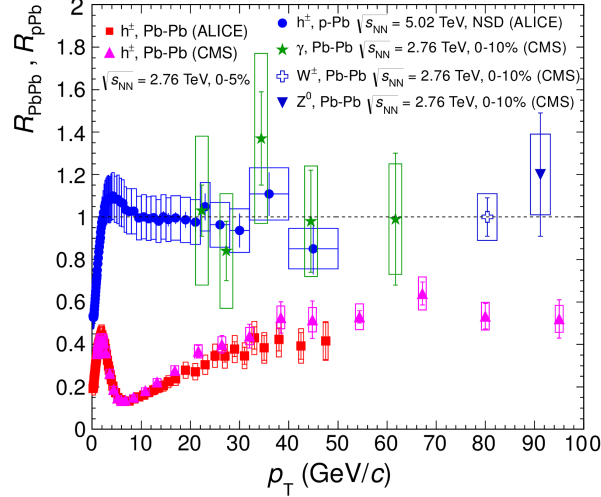


Figure 2.6: R_{AA} of charged hadrons as measured in central Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV and minimum bias p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV as a function of p_T . In addition, R_{AA} measurements of direct photons and weak bosons are shown for Pb–Pb collisions. Taken from Ref. [105].

Figure 2.6 shows also the nuclear modification factor of charged hadrons for minimum bias p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. In this case, the production of high- p_T hadrons does not seem to be affected by jet quenching and R_{pPb} is compatible with unity above 10 GeV/c.

By carrying out systematic studies of diverse experimental manifestations of jet quenching and comparing them with predictions of phenomenological models, one can determine the parameters of these models, known as jet transport coefficients [106], [107]. Particular attention is paid to the coefficient $\hat{q} = d\langle p_\perp^2 \rangle / dL$ that corresponds to the mean square of transverse momentum transferred between a jet parton and the medium per unit length. The JET Collaboration has made an immense contribution to the phenomenological study of jet quenching by extracting $\hat{q}$ and related quantities through the comparison of different model calculations with inclusive hadron R_{AA} measurements at RHIC and the LHC energies [108]. The JETSCAPE Collaboration supplemented and extended the results of the JET Collaboration [109] by employing the state-of-art statistical approaches, which are based on the Bayesian inference methods [110]. Figure 2.7 shows the quark jet transport coefficient $\hat{q}$, scaled by the density of the scattering centres $\sim T^3$, as a function of the medium temperature T for jets with $p = 100$ GeV/c (left) and as a function of jet

momentum for a fixed T at 300 MeV (right). Bayesian extraction of the coefficient $\hat{q}$

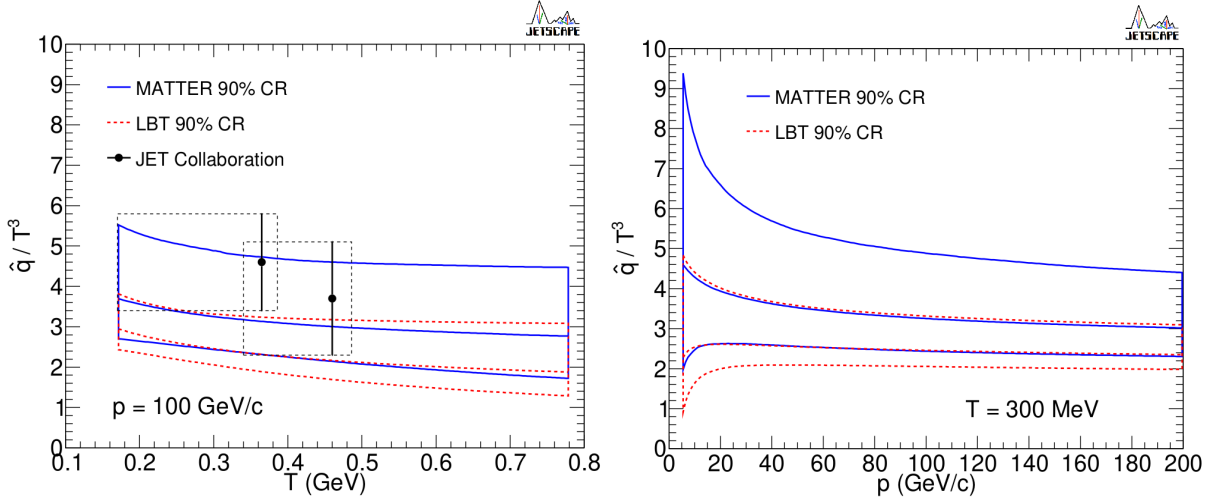


Figure 2.7: Quark jet transport parameter $\hat{q}/T^3$ as obtained from Bayesian parameter extraction utilizing the MATTER and Linear Boltzmann Transport (LBT) models of a parton energy loss. Left: as a function of the medium temperature. Right: as a function of jet momentum. Taken from Ref. [109].

was performed by applying two jet energy loss models, MATTER [111] and LBT [112], to the experimental data on the R_{AA} suppression of inclusive charged hadrons produced at two centralities of Au-Au collisions at $\sqrt{s_{NN}} = 200$ GeV (RHIC) and Pb-Pb collisions at the $\sqrt{s_{NN}} = 2.76$ and 5.02 TeV (LHC). One of the novel features of the JETSCAPE analysis is the extraction of the $\hat{q}$ parameter for the combined data set from both colliders, which enables to obtain a smooth distribution of the $\hat{q}$ dependence on the jet momentum and medium temperature. The JET Collaboration analysis performed calculations of $\hat{q}$ individually for RHIC and the LHC data due to a limitation of the least-squares minimization technique. The corresponding results are shown by the black markers on the left plot of Fig. 2.7. The JETSCAPE results are shown as a solid (dashed) line, corresponding to MATTER (LBT) energy loss models. The lines above and below the median value represent the 90% credible region (CR). The extracted $\hat{q}/T^3$ value exhibits a mild dependence on medium temperature and jet momentum. The JETSCAPE analysis went even further and considered for the first time a multi-stage jet evolution by incorporating the MATTER and LBT models simultaneously. The transition between the MATTER and LBT evolution was carried out using the average switching scale Q_0 , which should be treated as a matching scale between the high and low virtuality phases of parton evolution. Discussion of this issue goes beyond the scope of this thesis, inspect Ref. [109] for further details.

2.3.2 Dijet momentum imbalance

Since the scope of this thesis is a search for jet quenching via measurements of hadron–jet acoplanarity, it is also illustrative to consider how jet quenching affects dijets. The correlation of two probes, such as γ -jet, Z^0 -jet or a dijet, serves as the direct observable of parton energy loss. The ATLAS Collaboration studied the distribution of the dijet energy imbalance [113] to search for jet quenching in Pb–Pb collisions collected during the Run 1 data-taking period at the LHC. A year later, the CMS Collaboration released a study on the transverse momentum imbalance in Pb–Pb at $\sqrt{s_{\text{NN}}} = 2.76$ TeV, which brought additional insights into the transport of the emitted partonic energy outside the jet cone [85]. The latter measurement is discussed in more detail below.

The dijet transverse momentum imbalance is characterized by the asymmetry ratio

$$A_J = \frac{p_{T,1} - p_{T,2}}{p_{T,1} + p_{T,2}}, \quad (2.14)$$

where the subscripts “1” and “2” refer to jets with the highest and the second highest p_T in an event. These jets are called *leading* and *subleading*, respectively. Jet reconstruction was performed utilizing the modified iterative cone algorithm [114] with $R = 0.5$. The analysis required for the *leading* jet to have $p_{T,1} > 120$ GeV/ c , while a candidate for the *subleading* jet was constrained to have $p_{T,2} > 50$ GeV/ c . The azimuthal separation between the *leading* and *subleading* jets $\Delta\varphi_{1,2}$ was required to be greater than $2\pi/3$. According to the conservation law of 4-momentum, at LO and in the absence of the medium, the *leading* and *subleading* jets must have the same p_T and be very close to back-to-back in azimuth. Therefore, the ratio (2.14) should contain a peak at $A_J = 0$. However, at higher orders of perturbation theory, the distribution of the momentum imbalance A_J becomes wider even in pp collisions due to a parton splitting that leads to the creation of 3 or more jets in the final state. Figure 2.8 shows the evolution of the dijet momentum imbalance A_J as a function of the collision centrality. The solid line corresponds to the PYTHIA simulation embedded into the data. It can be observed that the peripheral Pb–Pb events in Fig. 2.8 (a) exhibit a peak at $A_J = 0$ and are in good agreement with PYTHIA + Data simulation, therefore supporting the idea of no partonic energy loss. However, when moving towards central Pb–Pb collisions (Fig. 2.8 (b) and (c)), a substantial suppression of events with the balanced *leading* and *subleading* jets is observed. Additionally, in most central 0–10% events, the mean value of the distribution is shifted. This observation is consistent with the assumption about partonic energy loss due to interaction with the medium created in central Pb–Pb collisions.

The dijet momentum imbalance can be restored by considering particles emitted outside the *subleading* jet cone. To investigate this aspect, a new quantity $\langle p_T^\parallel \rangle$ was introduced

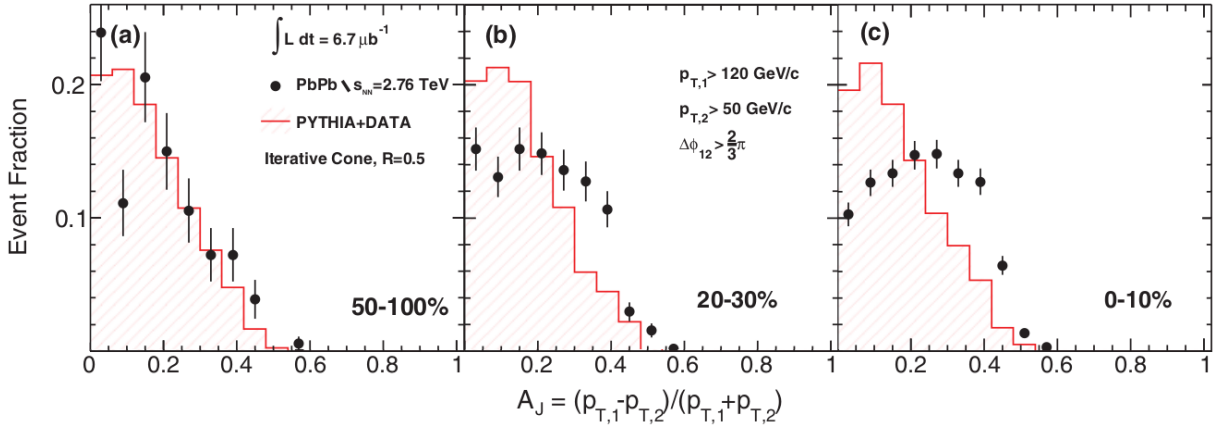


Figure 2.8: Momentum asymmetry ratio A_J (2.14) in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV in several centrality bins: (a) 50–100%, (b) 20–30%, and (c) 0–10%. The data are shown by the black markers. The solid line represents PYTHIA events embedded into Pb–Pb data. Taken from Ref. [85].

that is defined as the average sum of the track momenta projected onto the *leading* jet axis:

$$\langle p_T^{\parallel} \rangle = - \left\langle \sum_i p_T^i \times \cos(\varphi^i - \varphi^{\text{leading jet}}) \right\rangle_{\text{over events}}, \quad (2.15)$$

where p_T^i and φ^i are the momentum and azimuthal angle of a track “ i ”, respectively; $\varphi^{\text{leading jet}}$ is the azimuthal angle of the *leading* jet. Manifestly, the negative values of $\langle p_T^{\parallel} \rangle$ correspond to the tracks that have direction likewise as the *leading* jet, whereas the positive ones attribute to the tracks pointing in the direction of the *subleading* jet. Figure 2.9 shows the distribution of $\langle p_T^{\parallel} \rangle$ as a function of A_J for tracks with $p_T > 0.5$ GeV/ c and $|\eta| < 2.4$ in PYTHIA + HYDJET simulation (top) and central 0–30% Pb–Pb collisions (bottom). For this study, the azimuthal angle $\Delta\varphi_{1,2}$ was set to be larger than $5\pi/6$. The left and right panels illustrate the data for tracks inside and outside the cones with $R = 0.8$ placed around the *leading* and *subleading* jets, respectively. The size of the shaded areas reflects the amount of p_T carried by the tracks with p_T in the listed intervals to $\langle p_T^{\parallel} \rangle$. The in-cone imbalance of $\langle p_T^{\parallel} \rangle$, shown by the circle markers, both for the Monte Carlo (MC) and data, is compensated by $\langle p_T^{\parallel} \rangle$ from the out-of-cone region. More importantly, in the data, $\langle p_T^{\parallel} \rangle$ value in the out-of-cone region is composed mainly of soft tracks with $p_T \in (0.5, 4)$ GeV/ c , while in the MC simulation, the primary contribution comes from tracks with $p_T > 4$ GeV/ c . One can conclude that the interaction with the medium affects the jet fragmentation pattern, making it softer, and results in the transport of emitted parton energy at large angles with respect to the *subleading* jet axis.

It should be noted that dijet measurements at the LHC were preceded by similar measurements of dihadron correlations at RHIC. These measurements observed a substantial suppression of the hadron yield on the away-side in Au–Au collisions when compared to

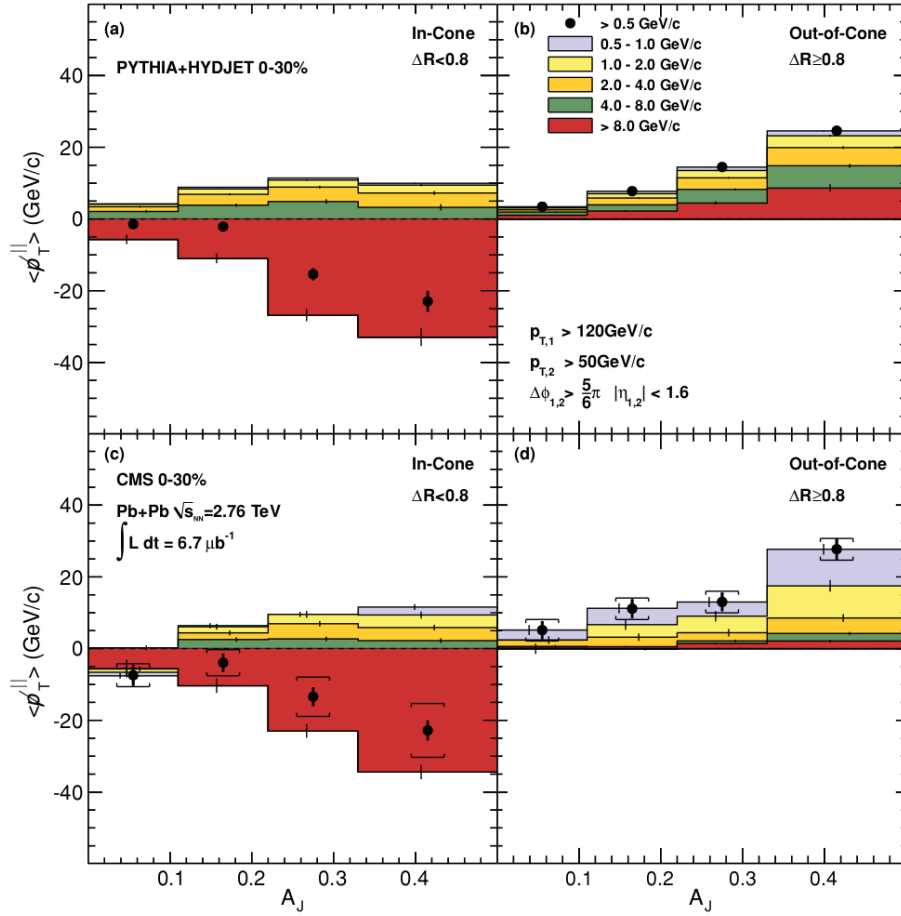


Figure 2.9: Average momentum $\langle p_T^{\parallel} \rangle$ (2.15) for tracks with $|\eta| < 2.4$ and $p_T > 0.5$ GeV/c, which are inside $\Delta R < 0.8$ (left column) and outside $\Delta R \geq 0.8$ (right column) cones with $R = 0.8$ placed around the *leading* and *subleading* jets. The distributions are shown as a function of the asymmetry ratio A_J for PYTHIA + HYDJET simulation (top row) and measured 0–30% central Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV (bottom row). The circle markers represent the sum of the negative and positive $\langle p_T^{\parallel} \rangle$ parts. The color boxes show the contribution to $\langle p_T^{\parallel} \rangle$ from the listed intervals of track p_T . Statistical uncertainties are illustrated as vertical bars. Systematic uncertainties are shown by empty brackets. Taken from Ref. [85].

pp and d–Au references [115], [116].

When interpreting results from dijet measurements, one needs to keep in mind that both jets can lose energy in the medium. Thus, it becomes impossible to obtain information about their initial properties before the interaction with the medium. Triggering on two probes that can experience energy loss results in a scattering bias towards the surface edge [117]. Instead of dijet production, one can consider γ/Z^0 -jet production, which serves as a “pure probe” [118] since the gauge bosons do not interact with the medium. Therefore, it becomes possible to deduce the energy of the recoil jet from the energy of the triggered gauge boson. Moreover, it is expected that this observable does not involve surface bias. Measurements of γ/Z^0 -jet production are performed both at

the LHC and RHIC [119]–[122]. The obtained results complement and reinforce previous measurements on dijet production in heavy-ion collisions. However, this observable has its drawbacks. For instance, it suffers from a much lower yield when compared to the production of dijets.

Chapter 3

Droplet of quark–gluon plasma in small collision systems?

The QGP formed in high-energy collisions of heavy nuclei has an extremely rapid dynamical evolution with a duration of about 10^{-22} s. Its formation leaves an “imprint” on the properties of the produced final-state particles. The global program for studying the properties of nuclear matter under extreme conditions started at the CERN SPS in 1986 [123]. A major breakthrough in this area was marked by discoveries that were made on RHIC in the early 2000s [124]–[127]. The measured collective flow of particles and jet quenching established the paradigm that the QGP is a strongly interacting, nearly perfect and almost opaque liquid. It is important to note that the QGP also has other formation signatures. Among them, one can single out the enhanced production of strange hadrons [128] and melting of the heavy-flavor quarkonia [129].

The first results from the LHC confirmed and extended many of the key RHIC results at much higher collision energies [130]–[134]. However, the LHC data also demonstrated that flow-like patterns emerge in high-multiplicity pp and p–Pb collisions [6], [9], [135], [136]. Furthermore, these systems exhibit enhanced production of strange hadrons [14]. These observations will be discussed in detail in the next section since they raise a critical question about the system size limits for the QGP formation. This question is also relevant for the QGP studies in heavy-ion collisions, where small collision systems are utilized as a reference measurement to distinguish the initial-state effects from the final-state in-medium interactions. Namely, this comparison assumes that the QGP cannot be produced in small systems.

To discard or prove the hypothesis for the QGP creation in small systems, one needs to carry out further searches for other QGP signatures. One of the prominent signatures is jet quenching, which is a direct consequence of opaque medium formation. The results of a search for jet quenching in small collision systems are discussed in more detail in

Sect. 3.3.

3.1 Evidence of collective flow

A non-central collision of two heavy nuclei results in an almond-shaped transverse profile of the interaction volume. The initial spatial anisotropy with respect to the reaction plane evolves by means of pressure gradients into a particle-momentum anisotropy. The azimuthal anisotropy of the final state particles can be parameterised in terms of Fourier coefficients [137],

$$v_n = \langle \langle \cos[n(\varphi - \Psi_{\text{RP}})] \rangle \rangle, \quad (3.1)$$

where $\langle \langle \dots \rangle \rangle$ denotes averaging firstly over all particles in an event and then over all events, φ and Ψ_{RP} are the azimuthal angle of a particle and the reaction plane angle, respectively. The coefficient v_2 for the second-order harmonic n is called the elliptic flow. The flow coefficients are customarily obtained from the angular correlations between the reconstructed particles. For heavy-ion collisions, values of the flow coefficients are reasonably well predicted by the contemporary hydrodynamic models [138].

3.1.1 Ridge-like structure

Studies of the two-particle correlations in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [139] showed that a pair of particles with $p_{\text{T}} \in (1, 3)$ GeV/ c , separated by a large interval of pseudorapidity $\Delta\eta = |\eta_1 - \eta_2|$, is correlated in azimuth. The measured distribution of the $\Delta\varphi = \varphi_1 - \varphi_2$ angle exhibits a long-range ridge structures at $\Delta\varphi \approx 0$, see Fig. 3.1 (a). Since both particles are separated by a large pseudorapidity distance, the ridge structure has to emerge early in the collision process [140]. The correlation function in

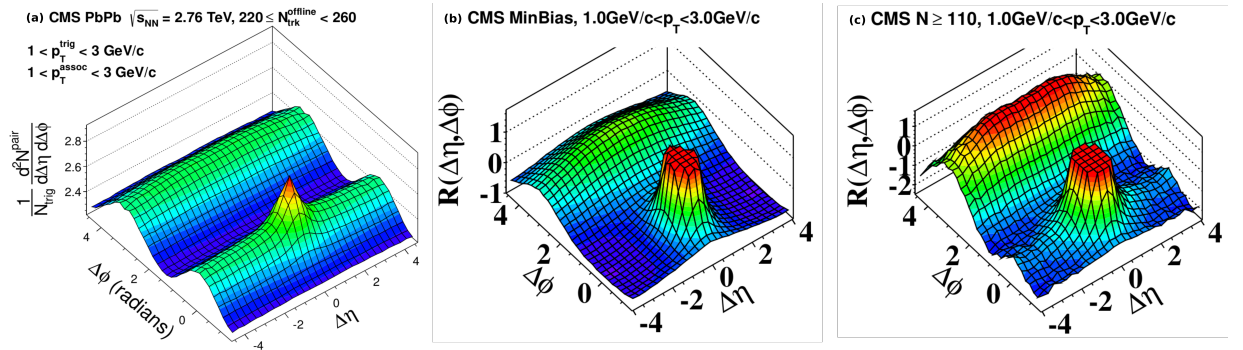


Figure 3.1: Two-particle correlation function for particles with $1 < p_{\text{T}} < 3$ GeV/ c in: (a) Pb–Pb events at $\sqrt{s_{\text{NN}}} = 2.76$ TeV with charged-track multiplicity $220 \leq N_{\text{trk}}^{\text{offline}} < 260$ [139]; (b) minimum bias pp events at $\sqrt{s} = 7$ TeV [6]; (c) high-multiplicity pp events at $\sqrt{s} = 7$ TeV with $N_{\text{trk}}^{\text{offline}} \geq 110$ [6].

Fig. 3.1 (a) has two additional notable features. The pronounced peak at $\Delta\eta, \Delta\varphi \approx 0$ is populated by particles from the near-side jet fragmentation, resonance decays, and pairs of identical bosons, which are correlated due to the Bose–Einstein correlations [141]. The second structure is a broad extended ridge at $\Delta\varphi \approx \pi$ formed by particles resulting from the fragmentation of back-to-back jets.

The CMS Collaboration performed measurements of the two-particle correlation function in pp collisions at $\sqrt{s} = 7$ TeV [6], see Fig. 3.1 (b), (c). The analysis was carried out utilizing minimum bias and high-multiplicity events. High-multiplicity events contained at least $N_{\text{trk}}^{\text{offline}} \geq 110$ particles after the offline selection. While the correlation function of particles with $p_{\text{T}} \in (1, 3)$ GeV/ c in minimum-bias events exhibits a clear plateau at $\Delta\varphi \approx 0$ for large $\Delta\eta$ values, the correlation function measured in high-multiplicity events unexpectedly exhibits a prominent and substantial near-side long-range correlation, similar to that in heavy-ion collisions. The absence of the ridge structure in minimum-bias events may be associated with a contribution of a non-flow component that strongly dominates in low-multiplicity collisions.

It is worth noting a recent measurement of two-particle angular correlations for charged particles produced in e^+e^- collisions at 183–209 GeV [142]. Events with high multiplicity featuring more than 50 charged particles demonstrate a remarkably similar ridge structure as measured in high-multiplicity pp collisions at different $\sqrt{s}$ [136].

3.1.2 Azimuthal correlations between several particles

Observation of the ridge in high-multiplicity p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [7]–[9] and central d–Au collisions at $\sqrt{s_{\text{NN}}} = 200$ GeV [143], [144] had stimulated further investigation of the collective-like behaviour in small collision systems utilizing the multi-particle cumulant technique [145], which has been successfully employed in nucleus–nucleus collisions.

In a collectively expanding bulk, the azimuthal correlations are shared among many particles. By construction, the multi-particle correlations substantially reduce the non-flow contribution. Figure 3.2 shows the elliptic flow harmonics $v_2\{n\}$ measured in pp, p–Pb, and Pb–Pb collisions at the LHC utilizing multi-particle cumulant analysis [136]. The notation $\{n\}$ denotes the four-, six-, eight-, and all-particle (based on the Lee–Yang zeros (LYZ) method [146]) azimuthal correlations. The harmonic $v_2^{\text{sub}}\{2, |\Delta\eta| > 2\}$ was obtained through the correlation of two particles separated by a large pseudorapidity interval $|\Delta\eta| > 2$. The results were averaged across particles with $p_{\text{T}} \in (0.3, 3)$ GeV/ c and expressed as a function of event multiplicity. In high-multiplicity pp collisions, the higher-order cumulant $v_2\{4\}$ and $v_2\{6\}$ values are comparable with each other within statistical and systematic uncertainties. Similar behaviour is observed in p–Pb and, more

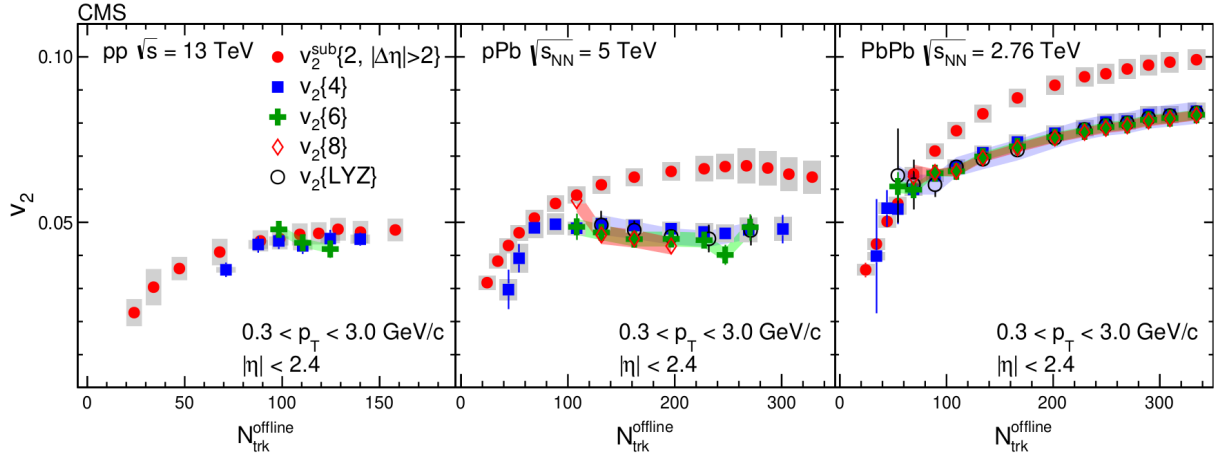


Figure 3.2: Values of v_2 extracted from multi-particle cumulants as a function of charged-particle multiplicity $N_{\text{trk}}^{\text{offline}}$ in pp collisions at $\sqrt{s} = 13$ TeV (left), p-Pb collisions at $\sqrt{s_{\text{NN}}} = 5$ TeV (middle) and Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV (right). The results were averaged over all particles with $0.3 < p_T < 3$ GeV/c and $|\eta| < 2.4$. For pp collisions, $v_2\{8\}$ and $v_2\{\text{LYZ}\}$ values were not calculated due to limited statistics. The shaded areas represent the systematic uncertainties. Taken from Ref. [136].

importantly, in Pb-Pb collisions. This observation supports the hypothesis of a collective origin of the long-range correlations measured in high-multiplicity pp and p-Pb collisions. An explicit relation $v_2\{2\} > v_2\{4\}$ is seen in p-Pb and Pb-Pb collisions. Within a hydrodynamic description, this difference comes from flow fluctuations on an event-by-event basis [147]. However, no difference $v_2\{2\} \approx v_2\{4\} \approx v_2\{6\}$ is observed in pp events within experimental uncertainties. Currently, there is no consensus on what causes the difference in pp data when compared to Pb-Pb.

In nucleus-nucleus collisions, it is assumed that elliptic anisotropies at low p_T arise from the hydrodynamic flow. However, the non-zero value of v_2 coefficient also emerges for high- p_T particles. This azimuthal anisotropy is expected to arise from the path-length dependence of the parton energy loss [148]. Figure 3.3 shows that v_2 of high- p_T particles is non-zero also in p-Pb events at $\sqrt{s_{\text{NN}}} = 8.16$ TeV. The data correspond to the minimum-bias triggered (MBT) events and events with a jet, which has $p_T > 100$ GeV/c [149]. Both distributions were scaled by a factor of 1.5 to enable visual comparison with the measurement of v_2 for 20–30% central Pb-Pb events at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. Remarkably, the overall trends of the v_2 values are similar for both collision systems.

Despite the similarity of experimental results for collective flow measurements in small systems with those from heavy-ion collisions, the question of the physical origin of multi-particle correlations remains open. In these circumstances, the only reliable and correct way is to compare the experimental results with theoretical predictions. Remarkably, hydrodynamic models applied to small systems provide a reasonable agreement with the

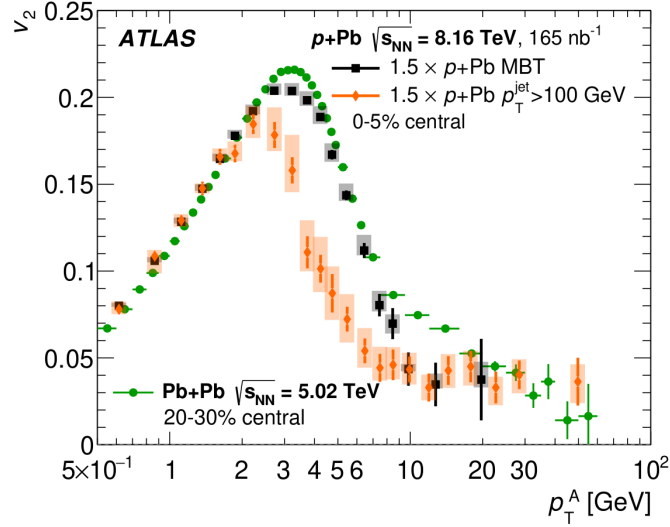


Figure 3.3: Elliptic flow coefficients v_2 as a function of particle p_T in p–Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV. The green markers show the data from Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV with 20–30% centrality. Color boxes denote the systematic uncertainties. Taken from Ref. [149].

experimental results [11], [12], [150]–[152], notwithstanding the system size being much smaller than in heavy-ion collisions. Note that the hydrodynamic description of small systems strongly depends on the applied initial state model. A detailed discussion about this topic can be found in Ref. [153]. However, sometimes hydrodynamic models provide a severe misinterpretation of the experimental results. For instance, the paper [154] demonstrates that the hydrodynamic calculations of $v_2\{4\}$ for pp collisions at $\sqrt{s} = 13$ TeV are always imaginary. This prediction contradicts the measured value of $v_2\{4\}$, see the left plot of Fig. 3.2. Besides the hydrodynamic description, other alternative models exist that provide a qualitative description of the observed flow effects in small collision systems. For instance, the Color Glass Condensate (CGC) model [155] predicts that the ridge structure is produced by the glasma graphs [140]. Theoretical work [156] demonstrated that the intrinsic QCD interference effects can lead to non-zero $v_2\{2\}$ and $v_2\{4\}$ coefficients, which are similar in shape and amplitude to the experimental data for pp collisions at the LHC. Additional material and reflections on this topic can be found in Refs. [147], [157]–[159].

3.2 Strangeness enhancement

Measurements of the strangeness-suppression factor (Wroblewski factor) [160] in e^-e^+ , $p\bar{p}$, and pp collisions reveal that the yield of $s\bar{s}$ pairs is about 20% of the light quark yields. However, this factor almost doubles in collisions of heavy ions. It indicates that other processes contribute to $s\bar{s}$ pair formation. Production of strange and multi-strange

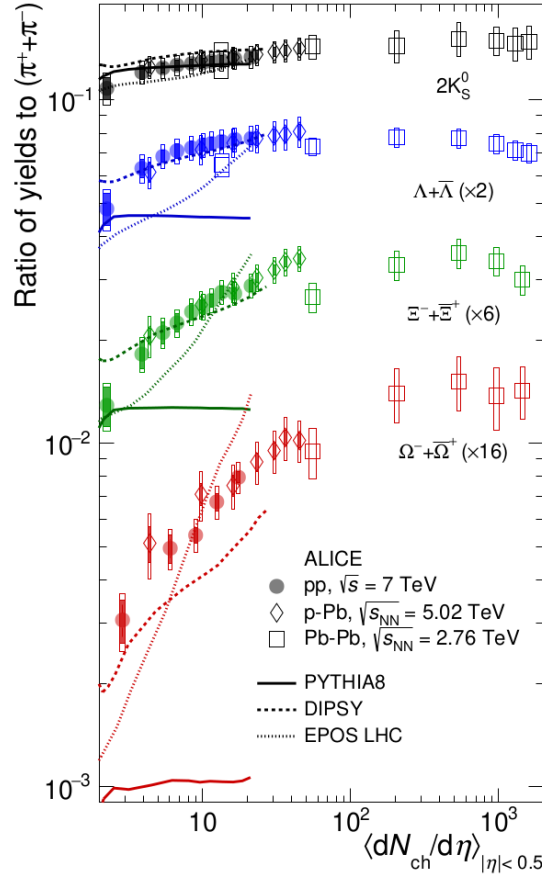


Figure 3.4: Relative yield of strange hadrons with respect to pions measured as a function of charged-particle multiplicity $\langle dN_{\text{ch}}/d\eta \rangle_{|\eta| < 0.5}$ at mid-rapidity in pp collisions at $\sqrt{s} = 7$ TeV and p-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The data are compared to Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV and model predictions. Taken from Ref. [14].

hadrons in heavy-ion collisions can be enhanced due to the thermal production of the $s\bar{s}$ pairs via gluon fusion $gg \rightarrow s\bar{s}$ within the QGP medium and chiral symmetry restoration [128]. Detailed theoretical calculations are presented in Ref. [161].

Strangeness enhancement has been observed many times in heavy-ion collisions in a wide centre-of-mass energy interval ranging from the SPS to the LHC; see Ref. [162] and references therein. The ALICE Collaboration has measured the ratio of the yields of K_S^0 , Λ , Ξ , and Ω hadrons to the yield of pions as a function of charged-particle multiplicity density at mid-rapidity in pp, p-Pb, and Pb-Pb collisions [14], see Fig. 3.4. The data suggest that the enhancement of strangeness production grows with the charged-particle density. The amount of enhancement is similar across different collision systems. Note, the ratio trend grows particularly faster for multi-strange Ξ and Ω hyperons.

The measured production of strange hadrons was compared to Monte Carlo event generators, see Fig. 3.4. The pure PYTHIA 8 simulation shows no signs of strangeness enhancement with multiplicity, while the DIPSY simulation [163] with the implemented rope hadronization model [16] qualitatively reproduces the data trends.

3.3 Jet quenching measurements in small collision systems

Phenomenological models predict that the partonic energy loss increases with the path length L traversed within the medium [101], [164]. Therefore, jet quenching is expected to be a fairly weak effect in pp and p–Pb collisions due to the smaller system size when compared to heavy-ion collisions [95], [165]. On the other hand, some early model calculations suggested a substantial jet quenching in high-multiplicity pp collisions [166]. However, the inclusive measurements of the nuclear modification factor for high- p_T hadrons and jets (Sect. 2.3.1) in minimum bias p–Pb collisions are consistent with unity, thus supporting the hypothesis of small parton energy loss [167]–[171].

3.3.1 Semi-inclusive jet measurements

Search for jet quenching in high-multiplicity (HM) small collision systems such as pp or p–Pb collisions is a challenging task due to multiplicity selection biases [172]–[174]. In the case of HM pp collisions, the objective becomes more complicated since even the standard measurements of the nuclear modification factor are not feasible due to the inability to determine the number of binary collisions $\langle N_{\text{coll}} \rangle$ for HM pp collisions. The semi-inclusive measurements of hadron–jet correlations [34] can overcome this issue without using the Glauber modelling.

The semi-inclusive correlation technique analyses charged-particle jets that recoil from a high- p_T trigger track (TT) [34]. The trigger track can be any charged particle with p_T in a specified range $\{X, Y\}$ GeV/ c . The presence of the TT ensures that the selected events contain at least one hard scattering process. In the next step, one considers all reconstructed jets that fall into the opposite azimuthal hemisphere with respect to the TT. It is worth noting that the transverse momentum scale of combinatorial jets in a high-multiplicity environment can be comparable to the p_T scale of true physical jets. Thus, combinatorial jets contaminate the measured jet spectrum. In inclusive measurements, their contribution is commonly suppressed by requiring the presence of a high- p_T track within each reconstructed jet [175]. However, this approach imposes a fragmentation bias on the jet population. Another way to suppress the yield of combinatorial jets is based on measurements of jets, which have p_T above a threshold [176]. In this case, the applied p_T cut can discard quenched jets with energy compatible with background jets due to substantial energy loss. Thus, the unbiased jet population is most suitable for studies of quenched jets at low p_T . The semi-inclusive measurement applies a purely data-driven method for removing the background jet component that does not bias the reconstructed spectrum of recoil jets. The method is based on the observation that combinatorial

background jets do not correlate with the hard scattering process. This aspect will be discussed later in this section.

The primary input to the hadron–jet correlation analysis is the per-trigger normalized yield of recoil jets associated with the $TT\{X, Y\}$, see the left-hand side of Eq. (3.2). This observable can be approximated by the ratio of two inclusive cross sections calculable in pQCD theory,

$$\frac{1}{N_{\text{trig}}^{\text{AB}}} \frac{d^3 N_{\text{jet}}}{dp_{\text{T, jet}}^{\text{ch, reco}} d|\Delta\varphi| d\eta_{\text{jet}}} \bigg|_{TT \in \{X, Y\}} \approx \frac{1}{\sigma^{\text{AB} \rightarrow h + X}} \frac{d^3 \sigma^{\text{AB} \rightarrow TT + \text{jet} + X}}{dp_{\text{T, jet}}^{\text{ch, reco}} d|\Delta\varphi| d\eta_{\text{jet}}} \bigg|_{TT \in \{X, Y\}}, \quad (3.2)$$

where $\sigma^{\text{AB} \rightarrow h + X}$ is the cross section for the inclusive TT production and $d^3 \sigma^{\text{AB} \rightarrow h + \text{jet} + X}$ is the differential cross section for coincidence production of the TT and recoil jet. Notation “AB” corresponds to a collision of some nuclei A and B. The jet yield is expressed as a function of jet transverse momentum $p_{\text{T, jet}}^{\text{ch, reco}}$, corrected for the presence of the underlying events using Eq. (2.11), and azimuthal opening angle $\Delta\varphi = \varphi_{\text{TT}} - \varphi_{\text{jet}}$ between the direction of the TT and recoil jet. Note that the $\Delta\varphi$ angle is measured in the range $(-\pi, \pi)$, therefore, due to the symmetry arguments, $|\Delta\varphi| \in (0, \pi)$. The approximation in (3.2) becomes more precise when the TT p_{T} range is narrow.

Jet quenching measurements require a comparison with a reference distribution customarily obtained from minimum bias (MB) pp collisions at the same centre-of-mass energy $\sqrt{s_{\text{AB}}}$. It is assumed that the QGP is not formed in MB pp collisions. In case of no nuclear matter effects, the cross sections in Eq. (3.2) can be expressed through the corresponding cross sections for MB pp collisions, taking into account the A–B collision geometry with the nuclear overlap function $\langle T_{\text{AB}} \rangle$:

$$\begin{aligned} \sigma^{\text{AB} \rightarrow h + X} &= \sigma^{\text{pp} \rightarrow h + X} \times \langle T_{\text{AB}} \rangle, \\ d^3 \sigma^{\text{AB} \rightarrow h + \text{jet} + X} &= d^3 \sigma^{\text{pp} \rightarrow h + \text{jet} + X} \times \langle T_{\text{AB}} \rangle. \end{aligned} \quad (3.3)$$

Substituting (3.3) into (3.2), one can notice that the terms $\langle T_{\text{AB}} \rangle$ in the numerator and denominator cancel each other identically. This means the semi-inclusive approach does not require the Glauber modelling when the recoil jet yield in A–B collisions is compared to the reference. Since no requirements are imposed on the species of A and B nuclei, one can also apply these arguments to HM p–Pb and pp collisions.

The workflow of the hadron–jet correlation analysis will be illustrated on the measurement performed by the ALICE Collaboration for p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [36]. The analysis was carried out in various intervals of event activity. Event activity was measured by the forward detectors located in the Pb-going direction: the scintillator counter V0A and the zero-degree calorimeter ZNA. The self-normalization property of

the semi-inclusive observable enabled the comparison of recoil jet p_T distributions obtained in events with a different activity. The left plot of Fig. 3.5 shows the per-trigger normalized p_T distributions of charged-particle recoil jets obtained for the 20% of p-Pb events that correspond to the largest signal in the ZNA detector. The analysis introduced

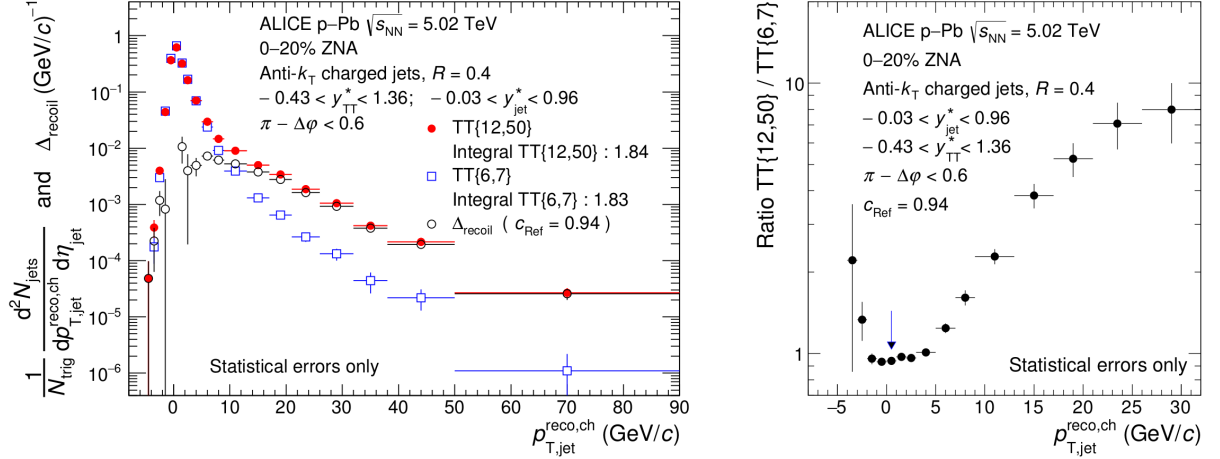


Figure 3.5: Left: the semi-inclusive distributions of charged-particle recoil jets in p-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The data correspond to 0–20% percentile interval of the ZNA signal. The Δ_{recoil} distribution (see Eq. (3.4)) is shown by the open circles. Right: ratio of per-trigger normalized recoil jet yields $\text{TT}\{12, 50\}/\text{TT}\{6, 7\}$. Taken from Ref. [36].

two exclusive $p_{T, \text{trig}}$ classes of the trigger track: $\text{TT}\{6, 7\}$ and $\text{TT}\{12, 50\}$. The data sample was split into two exclusive sets in proportion 10:90, where 10% of all events were used to search for $\text{TT}\{6, 7\}$, and the remaining 90% of events were used to search for $\text{TT}\{12, 50\}$. Selection of recoil jets applied a constraint $\pi - \Delta\varphi < 0.6$ rad on the opening azimuthal angle between the TT and recoil jets. There are two distinct p_T regions where the per-trigger normalized recoil jet yields exhibit different dependence on $p_{T, \text{trig}}$. At large $p_{T, \text{jet}}^{\text{reco, ch}}$, the recoil jet distribution associated with the $\text{TT}\{12, 50\}$ exceeds that for the $\text{TT}\{6, 7\}$. The observed dependence on $p_{T, \text{trig}}$ can be understood using a straightforward argument. The creation of the TT with a larger p_T biases towards hard scattering processes with a larger momentum transfer Q^2 , thereby the associated spectrum of recoil jets will also be harder.

The region of $p_{T, \text{jet}}^{\text{reco, ch}}$ around 0 GeV/c is populated primarily by combinatorial jets, which have the energy comparable to the energy of underlying events $\rho \times A_{\text{jet}}$ (Eq. 2.11). Notably, in this region, both distributions at the event-ensemble level have a similar shape and lie virtually on top of each other. This is illustrated in the right panel of Fig. 3.5, which shows a ratio of both spectra. To a first approximation, one can conclude that the recoil jet yield around 0 GeV/c does not emerge from a hard-scattering process producing the TT. Therefore, it has no dependence on $p_{T, \text{trig}}$. In Pb–Pb collisions [34], this

combinatorial background component is exceedingly larger than in p–Pb or pp, thereby, the corresponding distributions are substantially broader.

The uncorrelated jet component can be eliminated on a statistical basis via the subtraction of two per-trigger normalized yields associated to exclusive TT classes,

$$\Delta_{\text{recoil}} = \frac{1}{N_{\text{trig}}} \frac{dN_{\text{jets}}}{dp_{\text{T, jet}}^{\text{reco, ch}}} \bigg|_{\text{TT}\{12,50\}} - c_{\text{Ref}} \times \frac{1}{N_{\text{trig}}} \frac{dN_{\text{jets}}}{dp_{\text{T, jet}}^{\text{reco, ch}}} \bigg|_{\text{TT}\{6,7\}}. \quad (3.4)$$

Jet reconstruction fills the entire recoil region with jet candidates. Hence, both per-trigger normalized recoil jet distributions have close values of their integrals, see the legend in the left plot of Fig. 3.5. It means that, on average, the number of recoil jets per TT{6, 7} and TT{12, 50} is equal. However, TT{12, 50} is associated with a larger number of correlated physical jets than TT{6, 7}. Since the acceptance for recoil jets is finite, the TT{12, 50} can be associated with a smaller number of combinatorial jets simply because no acceptance is available for them. The correction factor c_{Ref} was introduced to account for this geometric effect. The value of c_{Ref} was extracted from the data utilizing the ratio between the recoil jet distributions in the bin $p_{\text{T, jet}}^{\text{reco, ch}} \in (0, 1)$ GeV/ c (see the blue arrow in the right plot in Fig 3.5). A typical span of the c_{Ref} values is 0.92–1.0. The open circle markers in the left plot of Fig. 3.5 show the resulting raw Δ_{recoil} distribution, which was further corrected for instrumental effects and momentum smearing due to local background fluctuations.

The same procedure was carried out for events corresponding to the 50–100% percentile interval of the ZNA signal. Figure 3.6 compares the Δ_{recoil} distributions obtained in events with high (0–20% ZNA) and low (50–100% ZNA) activity. The ratio is compatible with unity suggesting that possible jet quenching signals are below the resolution of the measurement. Nevertheless, the data were used to set an upper limit on the out-of-cone energy loss that would be transferred out of $R = 0.4$ jet cone in the central events relative to peripheral ones. The out-of-cone energy loss was found to be less than 0.4 GeV/ c at 90% confidence level. Note that the semi-inclusive analysis of recoil jets with $R = 0.5$ in 0–10% Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [34] showed a substantial medium-induced transport of partonic energy outside the jet cone at a level of 8 ± 2 (stat) GeV/ c when compared to pp data at $\sqrt{s} = 2.76$ TeV simulated by the PYTHIA 6.425 event generator [177].

3.3.2 Medium-induced acoplanarity

Jet quenching is a combination of complex phenomena that reflect various aspects of the parton interaction with the medium constituents. The parton energy loss is inextricably

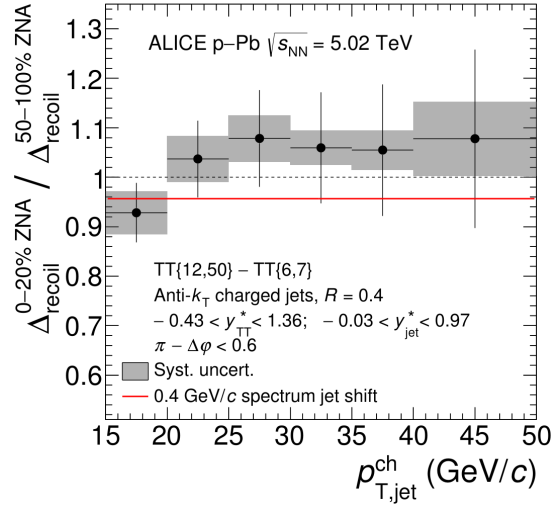


Figure 3.6: Ratio of Δ_{recoil} distributions measured in p-Pb events at $\sqrt{s_{\text{NN}}} = 5.02$ TeV with high (0–20 % ZNA) and low (50–100 % ZNA) event activity. The gray boxes represent the systematic uncertainties. The red line shows the expected value of the ratio provided that jet quenching in the central events would cause a shift of the spectrum measured in the peripheral events by 0.4 GeV/c. Taken from Ref. [36].

associated with the parton transverse momentum broadening $p_{\perp}$ [178]. The medium-induced $p_{\perp}$ broadening can be quantified by the measurements of hadron-jet azimuthal acoplanarity. However, hadron-jet correlations exhibit finite acoplanarity even in the absence of the deconfined medium. Hadron-jet acoplanarity in a vacuum is a direct consequence of the soft gluon emission, associated with Sudakov logarithms [179], and the intrinsic transverse momentum of the initial state partons. Thus, in-medium interaction should lead to an additional broadening of the acoplanarity distribution. Theoretical works [38], [180] have shown that Sudakov radiation has the dominant contribution to hadron-jet acoplanarity for high- p_{T} jets produced at the LHC energies. Thereby, the medium-induced hadron-jet acoplanarity should be substantially pronounced at RHIC energies, while it is necessary to consider low p_{T} jets at the LHC [38].

Multiple small-angle scatterings of a parton within the medium lead to a smearing of the hadron-jet acoplanarity distribution around π . At the same time, the occurrence of a single hard Moliere scattering (like Rutherford scattering [181]) should be visible in the acoplanarity distribution at azimuthal angles far away from π . Hypothetically, at short distances, the strongly interacting QGP could contain quasiparticles, which are weakly bound states of quarks or/and gluons [182]. Thus, a search for an enhancement in the large angle deflection of the recoil jet can provide valuable information about the microscopic structure of the QGP. However, this measurement is challenging due to the requirement of significant event statistics.

In this thesis, a search for jet quenching effects in HM pp collisions at $\sqrt{s} = 13$ TeV is

carried out by measuring the medium-induced hadron–jet acoplanarity. The acoplanarity measurements are based on the semi-inclusive approach discussed in Sect. 3.3.1. The analysis and main results will be discussed in Chap. 5.

Chapter 4

ALICE detector

A Large Ion Collider Experiment (ALICE) is one of the four large experiments at the LHC. The ALICE experiment has an extensive physics program, primarily aimed at studying properties of the QCD matter produced in collisions of ultrarelativistic lead nuclei [183], [184]. Additionally, the ALICE experiment also investigates pp and p–Pb collisions, which serve not only as a baseline for Pb–Pb data but which are of great interest in their own right. The ALICE apparatus is designed to provide reliable reconstruction and identification of particles in a high-multiplicity environment. The apparatus incorporates many detector systems. Further, only the Inner Tracking System (ITS), the Time Projection Chamber (TPC), and the V0 detectors will be considered since they are particularly important for this analysis. It is worth stressing that some of the presented technical details were only relevant for LHC Run 1 and Run 2 since the discussed detectors underwent a significant upgrade before LHC Run 3 [185]. A detailed description of the other systems can be found in Ref. [186].

4.1 Inner Tracking System

The ITS is a silicon tracker detector, which is designed to resolve the following objectives [187]:

- localization of the primary vertex in Pb–Pb collisions with a resolution of about 100 μm in the transverse plane ($r\varphi$) and about 250 μm along the beam axis (z) for tracks with $p_{\text{T}} = 500 \text{ MeV}/c$;
- reconstruction of the secondary vertices from weak decays of strange, charm, and beauty particles;
- tracking and identification of particles with $p_{\text{T}} < 200 \text{ MeV}/c$;

- improving position, momentum, and angular resolution for tracks reconstructed by the TPC;
- reconstruction of the high- p_T tracks, which pass the dead sectors of the TPC.

During the LHC Run 1 and Run 2, the ITS consisted of 6 coaxial cylindrical layers of position-sensitive silicon detectors. It covered the mid-rapidity range $|\eta| < 0.9$ for all primary vertices reconstructed within $z = \pm 5.3$ cm (i.e. $\pm 1\sigma$ of the luminous region) with respect to the nominal interaction point. The layers were arranged to surround the beam pipe at the radial positions ranging from 39 to 430 mm. The first two layers of the ITS were equipped with Silicon Pixel Detectors (SPD), which provided precise position measurements for track hits necessary for secondary vertex reconstruction. The SPD sensor had a matrix containing 256×160 pixels with a pitch size of $50 \mu\text{m} \times 425 \mu\text{m}$. The intrinsic spatial resolution of the SPD sensor was about $12 \mu\text{m}$ in the plane transverse to the beam axis and about $100 \mu\text{m}$ along the beam axis. The two intermediate layers consisted of Silicon Drift Detectors (SDD) with good multitasking capabilities. The two outer layers were made of Silicon Strip Detectors (SSD). They were particularly important for matching of tracks from the TPC to the ITS. The SDD and SSD layers had an analogue readout. Therefore, they could perform measurements of dE/dx to identify particles with low p_T . Detailed technical specifications concerning the ITS are presented in Refs. [186], [188].

4.2 Time Projection Chamber

The TPC is the main 3-dimensional tracking detector of the ALICE central barrel. It covers a pseudorapidity acceptance of $|\eta| < 0.9$ in full azimuth. The detector construction is based on a cylindrical field cage. The active volume of 90 m^3 is filled with the mixture of Ne/CO₂/N₂ gases at atmospheric pressure. The TPC has an inner and outer radius of 85 cm and 250 cm, respectively. The total length of the detector is about 500 cm. The detector's active volume is divided into two symmetric drift regions by the central electrode, which is kept at -100 kV potential with respect to the forward/backward TPC readout plates. The central electrode and two opposite axial voltage degraders form a highly uniform axial electric field of 400 V/cm . When a charged particle traverses the TPC volume, it ionizes the gas. Under the influence of the electric field, the ionization electrons drift towards the readout multi-wire proportional chambers mounted on the two end-plates of the detector. The drift time of the electron for the full drift path of 250 cm is about $92 \mu\text{s}$. This time is large compared to interaction rates of 100 kHz in pp collisions [189]. Thus, a relatively slow charge collection leads to pileup events in the

detector volume.

The TPC can provide up to 159 space–time points for a track reconstruction [190]. In addition, the TPC provides particle identification by measuring the charge released during ionization. Further information on the TPC construction and operation can be found in Refs. [191], [192].

4.3 V0 detector

The V0 consisted of two arrays of scintillator counters, denoted as V0A and V0C. The detectors covered the pseudorapidity intervals V0A: $2.8 < \eta < 5.1$ and V0C: $-3.7 < \eta < -1.7$ over the entire azimuth. Each detector was divided into 32 individual scintillator counters: 4 rings segmented into 8 equal sectors of 45° . The detectors were installed on both sides of the nominal collision vertex, see Fig. 4.1. The V0A detector was located at a distance of 3.4 m from the vertex, while the V0C detector was placed in front of the hadronic absorber, 0.9 m away from the vertex. The asymmetry in location was solely conditioned by the lack of free space.

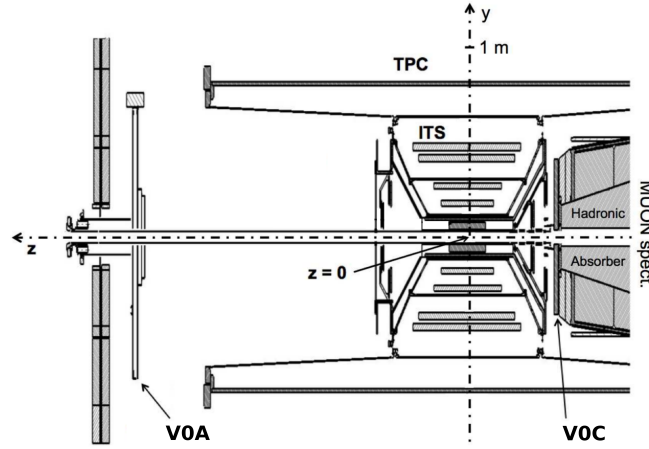


Figure 4.1: Schematic location of the V0A and V0C detectors inside the ALICE apparatus. Adopted from Ref. [193].

The detectors' components, such as scintillation material, photomultiplier tubes, and optical fibres, are susceptible to accumulated radiation loads, leading to detector performance degradation over time. The highest radiation doses were absorbed by the first ring of the V0C detector [194]. The radiation hardness of the scintillator material of the first ring was tested with a 60 MeV proton beam with the maximum accumulated radiation dose up to 300 krad corresponding to 10 years of operation of the ALICE experiment [195]. It was shown that during 10 years of operation, one should expect signal attenuation of the order of 30%. However, by the end of 2012, i.e. after 3 years of operation, a signal amplitude loss of 10 to 40% was observed [193]. Further studies have shown that the dom-

inant contribution to signal attenuation comes from the ageing of V0 photomultipliers. To compensate for the decrease in amplitude, one can adjust the gain to its previous level by increasing the voltage on the photomultipliers. As a consequence, the signal amplitude measured by the V0 detector varies for different periods of data taking.

Within the ALICE experiment, the V0 detector has several key functions. It provides a minimum-bias interaction trigger based on the coincidence of signals from the V0A and V0C detectors. The signal of the V0 detector is utilized to characterize the centrality of collisions since the amplitude of the V0 signal is proportional to the multiplicity of charged particles traversing the V0 scintillators. This multiplicity is expected to correlate with the collision impact parameter. The V0 detector can trigger events with the required centrality in Pb–Pb collisions by imposing a threshold on the measured signal. Note that the trigger efficiency and the charged-particle multiplicity distribution are strongly affected by the background from secondary particles. Any material surrounding the V0A and V0C detectors serves as a source of secondary particles. The transport of particles inside the ALICE was studied in Ref. [194] using simulations of a detailed detector model implemented in GEANT3 [196]. The efficiency of the minimum-bias trigger, which required at least one particle in the acceptance of each of the V0 detectors, was estimated to be 77% and 83% in pp collisions generated by PYTHIA standalone and PYTHIA + GEANT, respectively. The PYTHIA + GEANT simulation was also utilized to study the multiplicity distributions measured by the 4 rings of the V0A and V0C detectors in pp collisions. Figure 4.2 shows the multiplicity distributions for the first ring only. The distributions are close to each other when there are no secondary-generated particles. A similar trend is observed also for the other rings. However, the presence of secondary particles significantly distorts the shape of the distributions, especially in the case of the V0C detector.

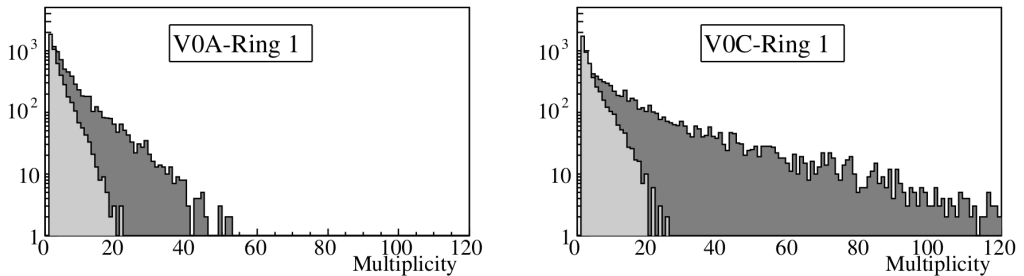


Figure 4.2: Charged-particle multiplicity measured by the first ring of the V0A (left) and V0C (right) detectors in pp collisions simulated by the PYTHIA event generator. The light and dark gray data correspond to events, in which particles are transported in a vacuum and the ALICE environment simulated by a detailed GEANT3 model of the apparatus, respectively. Taken from Ref. [194].

Collisions of the accelerated LHC beams with the residual gas in the beam pipe gener-

ate events that constitute the background for beam–beam collisions. The influence of the beam–gas background is more pronounced in pp collisions than in Pb–Pb [197]. Events from the beam–beam and beam–gas interactions can be discriminated by the arrival time of the signal in the two V0 detectors. A collision with the residual gas produces a signal that arrives earlier in time than a signal coming from the collision vertex. Figure 4.3 shows correlations between the sum and difference of signal times recorded by the V0A and V0C detectors. Three types of events can be distinguished. The largest and brightest spot corresponds to the beam–beam collision. The remaining two faint peaks originate from the p–gas collisions. The latter events are then rejected from data reconstruction.

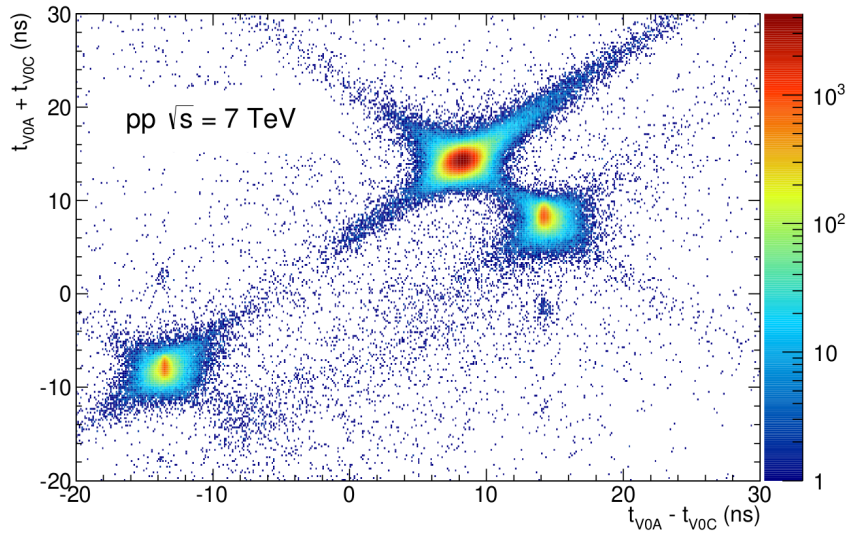


Figure 4.3: The correlation between the sum and the difference of signal times arrival in the V0A and V0C detectors. The main peak corresponds to events from pp collisions at $\sqrt{s} = 7$ TeV. The two other smaller peaks represent events originating from p–gas collisions. Taken from Ref. [197].

In addition to the main functions discussed above, the V0 detector is also used in luminosity and event plane measurements. Further details on these topics are provided in Ref. [197]. Comprehensive information about the technical features of the detector can be found Ref. [194].

Chapter 5

Measurements of hadron–jet acoplanarity in pp collisions

5.1 Data sample

The analysis is based on pp collision data at a centre-of-mass energy of $\sqrt{s} = 13$ TeV collected in 2016, 2017, and 2018. The events were recorded using minimum bias (MB) and high multiplicity (HM) online triggers. Both triggers used the light signal measured by the forward scintillator detectors V0A and V0C. The light yield is proportional to the total multiplicity of charged particles passed through the V0 arrays in an event. The MB trigger required a time coincidence of the signals from the V0A and V0C, while the HM trigger measured the sum amplitude of the signals, denoted as $V0M = V0A + V0C$, which was restricted to be greater than $5 \times \langle V0M \rangle$, where $\langle V0M \rangle$ corresponds to the mean signal amplitude in MB events. The HM trigger selects 0.1% of MB events with the largest value of the V0M signal. Accepted events were required to have the position of the primary vertex within $|z_{\text{vtx}}| < 10$ cm (see Fig. 5.1), where z_{vtx} is the location along the beam axis relative to the nominal interaction point of the ALICE detector. After the offline event selection, the MB and HM data sets comprised 1.4×10^9 and 9.1×10^8 events, respectively.

This analysis also incorporates MC simulations of pp collisions at $\sqrt{s} = 13$ TeV, which are based on the PYTHIA 8 event generator with the Monash tune (hereafter denoted as PYTHIA 8 Monash) [64], and a sophisticated GEANT3 [196] model of the ALICE detector, which emulated the transport of PYTHIA-generated particles through the detector material including production of secondary particles. Events generated by PYTHIA without detector effects are denoted as “particle-level”, while such events passed through GEANT3 simulation are denoted as “detector-level”. The particle-level simulation contained all generated final state charged particles, except for weak decay daughters [198]. The coincidence requirement of the MB trigger was modelled by requiring particle-

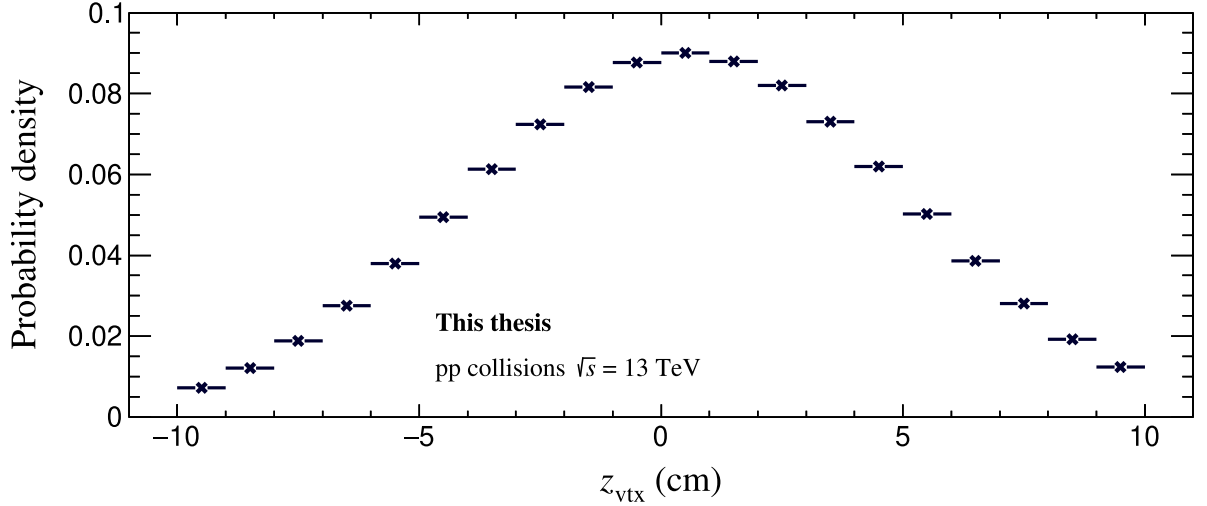


Figure 5.1: Distribution of z coordinate of the primary vertex reconstructed in MB pp collisions at $\sqrt{s} = 13$ TeV.

level events to have at least one charged final state particle in both V0A and V0C, while detector-level events were required to have simulated hits in both V0A and V0C. PYTHIA events were generated in both minimum-bias mode (all soft QCD processes) and triggered mode using bins in $\hat{p}_T$ [199].

5.2 Measurement of event activity

Figures 5.2 (a) and (b) show the measured V0M multiplicity distributions in MB and HM events collected in each year of the considered data sample. During different data-taking periods, the gain of the V0 photomultipliers may differ, resulting in a different amplitude of the measured V0M signal. Moreover, the ageing of the scintillator material also affects the signal amplitude. To eliminate the impact of the detector configuration on the V0M amplitude, this analysis considers the event multiplicity scaled by the mean value $\langle \text{V0M} \rangle$ from MB events, see Fig. 5.2 (c) and (d). After the scaling, the distributions almost match up to $\text{V0M}/\langle \text{V0M} \rangle \approx 9$, while they have a different slope at $\text{V0M}/\langle \text{V0M} \rangle > 9$. This suggests that there is another source of HM events in this region, which starts dominating. Presumably, this source is a residual pileup. To verify this assumption, a toy Monte Carlo study was carried out, where the number of pp interactions per bunch-crossing was simulated. Note that in the analysis, the contribution of the pileup events was suppressed offline by discarding events with multiple reconstructed primary event vertices, which were separated by more than 0.8 cm along the beam axis. Input parameters for the toy MC were taken from the real data. The mean number of interactions per bunch crossing was $\mu = 0.011$. The flow of the toy MC was as follows:

1. sample the number k of pp interactions per bunch-crossing from a Poisson distribu-

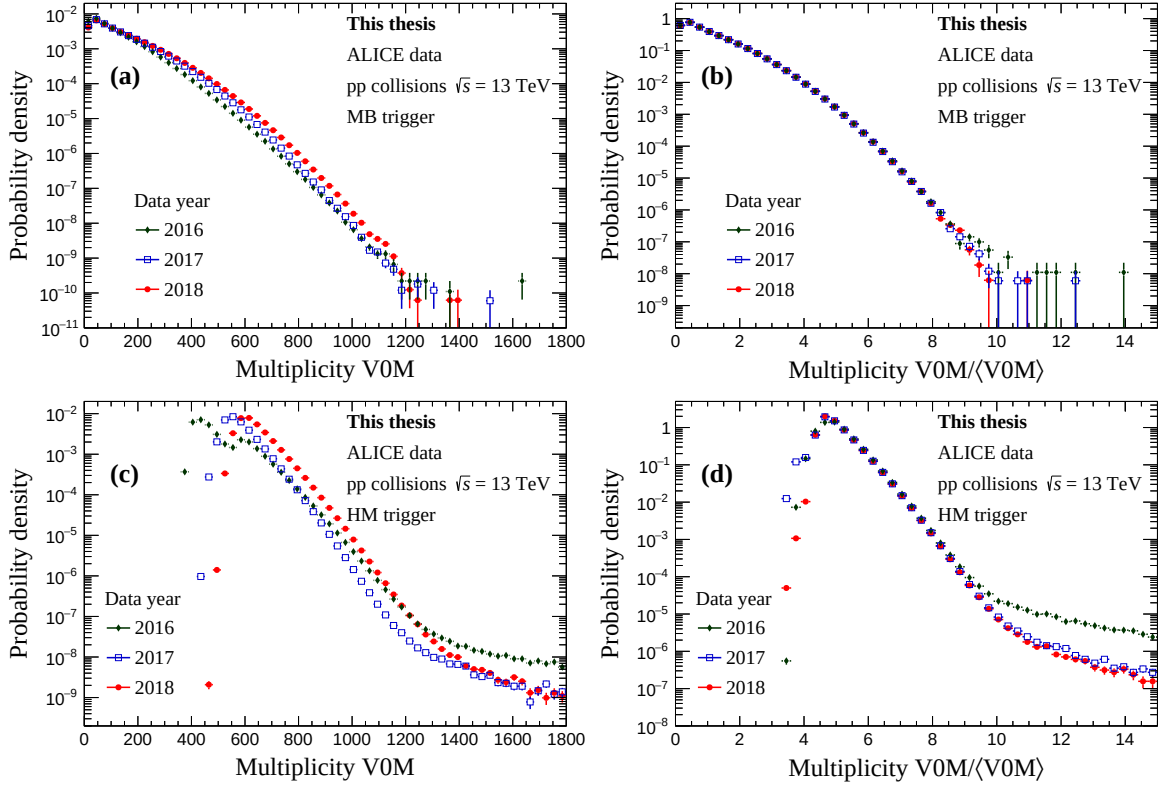


Figure 5.2: V0M multiplicity of events collected by the MB and HM triggers in 2016, 2017, and 2018: (a), (c) absolute signal; (b), (d) scaled by the mean multiplicity $\langle V0M \rangle$ of MB events.

tion [200] with the mean value of $\mu = 0.011$,

$$P(k) = \frac{\mu^k}{k!} e^{-\mu}; \quad (5.1)$$

2. sample a z vertex coordinate from the distribution in Fig. 5.1 for each simulated pp interaction;
3. sample $V0M/\langle V0M \rangle$ value from the total multiplicity distribution summed over 3 years of data taking (Fig. 5.2 (b)). The total event multiplicity is then obtained by summing up the sampled $V0M/\langle V0M \rangle$ values from all simulated pp interactions;
4. the case, when $k \geq 2$ and the distance between *any pair* of primary vertices is *more than* 0.8 cm, is considered as a pileup event that can be rejected (condition for offline pileup suppression). Otherwise, if $k = 1$ or the distance between *all pairs* of primary vertices is *less than* 0.8 cm, then the simulated event is considered as MB-triggered event.

Figure 5.3 shows the results of MC simulations of the $V0M/\langle V0M \rangle$ distribution for all MB events, including those with residual pileup (a reflection of what is measured in the

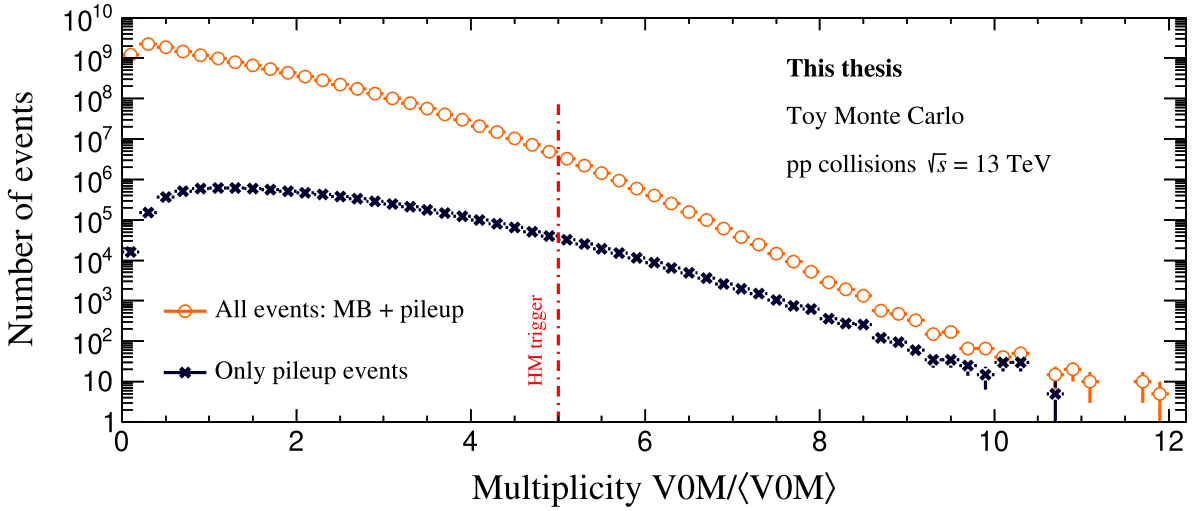


Figure 5.3: $V0M/\langle V0M \rangle$ distributions obtained in the toy Monte Carlo study of residual pileup fraction. See text for further details.

experiment) and for purely pileup events. One can see that the majority of pileup events have low event activity (small $V0M/\langle V0M \rangle$). The contribution of pileup events to MB and HM data samples is equal to the ratio of the integrals of the two curves. In the case of HM events, the integral was calculated for $V0M/\langle V0M \rangle > 5$. The residual pileup constitutes about 0.06% of MB events, while its fraction in the case of HM events is 1.3%. It is important to note that the contribution of pileup events to the HM sample increases with increasing $V0M/\langle V0M \rangle$, and both distributions almost coincide at $V0M/\langle V0M \rangle > 9$. Therefore, it was decided to exclude all events where $V0M/\langle V0M \rangle > 9$ from the analysis.

It should be noted that the multiplicity scaling also enables comparison of the measured data with MC simulation, where the multiplicity has an alternative definition. Figure 5.4 shows the $V0M/\langle V0M \rangle$ distribution for MB pp collisions at $\sqrt{s} = 13$ TeV. The red dashed-dotted line indicates the lower limit for HM event selection, $V0M/\langle V0M \rangle = 5$. The figure also shows the particle-level and detector-level $V0M/\langle V0M \rangle$ distributions for PYTHIA pp events at $\sqrt{s} = 13$ TeV generated in the minimum-bias mode. The particle-level simulations defined the $V0M$ multiplicity as the total number of charged-final-state particles in the acceptances covered by the V0 detectors. In this case, the $\langle V0M \rangle$ roughly equals to 31 particles. The detector-level simulations cannot describe the whole shape of the measured multiplicity distribution. However, the trend is reproduced fairly close. A vast difference in the amplitudes of the particle- and detector-level distributions at high multiplicities arises from the contribution of secondary particles [194] generated within the ALICE environment. Thus, unlike the particle-level spectrum, the measured and detector-level distributions have an extended tail.

In order to compare the particle-level HM-selected distributions to data, it was assumed that the yield of secondary particles generated in the detector material is, on

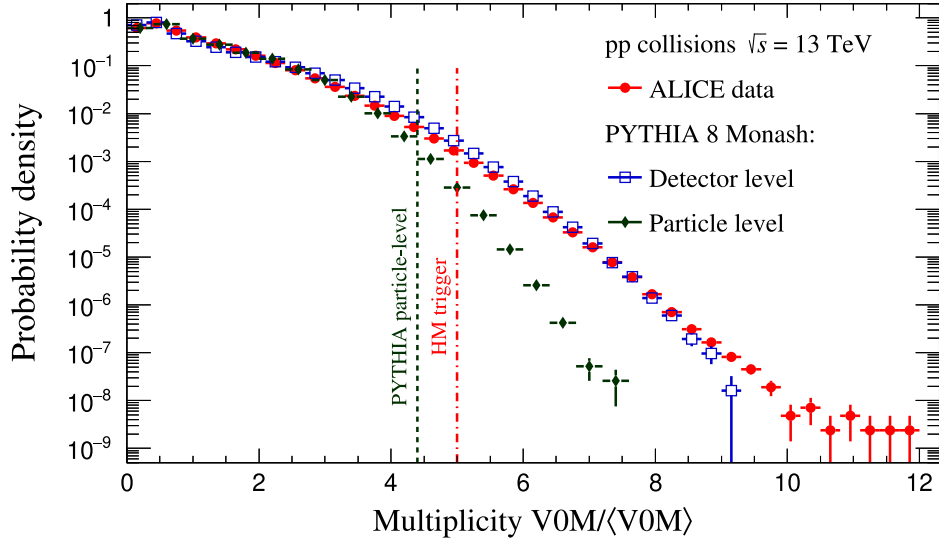


Figure 5.4: $V0M/\langle V0M \rangle$ distribution measured in MB pp collisions at $\sqrt{s} = 13$ TeV. The ALICE data are compared to the particle- and detector-level MC simulation. The vertical dashed lines indicate the lower bound for HM selection for data and particle-level simulations. Accessible in Ref. [39].

average, proportional to the primary multiplicity in the acceptances of the V0A and V0C. Therefore, the HM cut for the particle-level data was chosen to select the same fraction of the MB cross section (0.1%) as the HM selection $V0M/\langle V0M \rangle > 5.0$ used for data. This cut corresponds to $V0M/\langle V0M \rangle > 4.4$, as indicated by the black dashed line in Fig. 5.4. Note that the selection of the same cross section fraction (0.1%) at the particle level only selects the same fraction of the ALICE multiplicity distribution. It cannot select the same population event-by-event due to significant fluctuations in the correlation between particle- and detector-level events. For the detector model used in this analysis, about 35% of events passing the HM selection at the detector level also do so at the particle level. This is a generic issue, with similar features expected for any event selection in small systems based on the measurement of event activity in the forward rapidity.

5.3 Jet reconstruction

The analysis considered global tracks which had at least one hit in the SPD layers with more than 70 TPC spatial points and more than 80% of the geometrically findable TPC spatial points [19]. Reconstruction of the primary event vertex was performed offline based on reconstructed global tracks, see Ref. [201] for further details. During the data taking, the SPD detector had regions of local inefficiency, which resulted in a non-uniform azimuthal distribution of global tracks. Such detector performance would negatively affect jet reconstruction and measurements of φ -dependent observables, for instance, the

hadron-jet acoplanarity. To restore uniformity of the azimuthal distribution, global tracks were supplemented with complementary tracks in the regions, where the SPD detector was inactive. The combined set of global and complementary tracks is called hybrid tracks [19]. Complementary tracks were required to pass the same quality cuts as global tracks, except for a hit in the SPD layers. To ensure a similar momentum resolution as for global tracks, complementary tracks were re-fitted by requiring the primary vertex as the origin of the tracks. Finally, the hybrid tracks with $|\eta| < 0.9$, $\varphi \in (0, 2\pi)$, and $p_T > 0.15$ GeV/ c were accepted for the analysis.

The analysed data have the corresponding anchored MC simulations, which reflect the data-taking conditions and detector performance. Figure 5.5 shows a comparison of the azimuthal angle and pseudorapidity distributions of hybrid tracks reconstructed in MB and HM pp events at $\sqrt{s} = 13$ TeV and in the detector-level MC simulations. The detector-level simulations describe both MB and HM data samples well. Figure

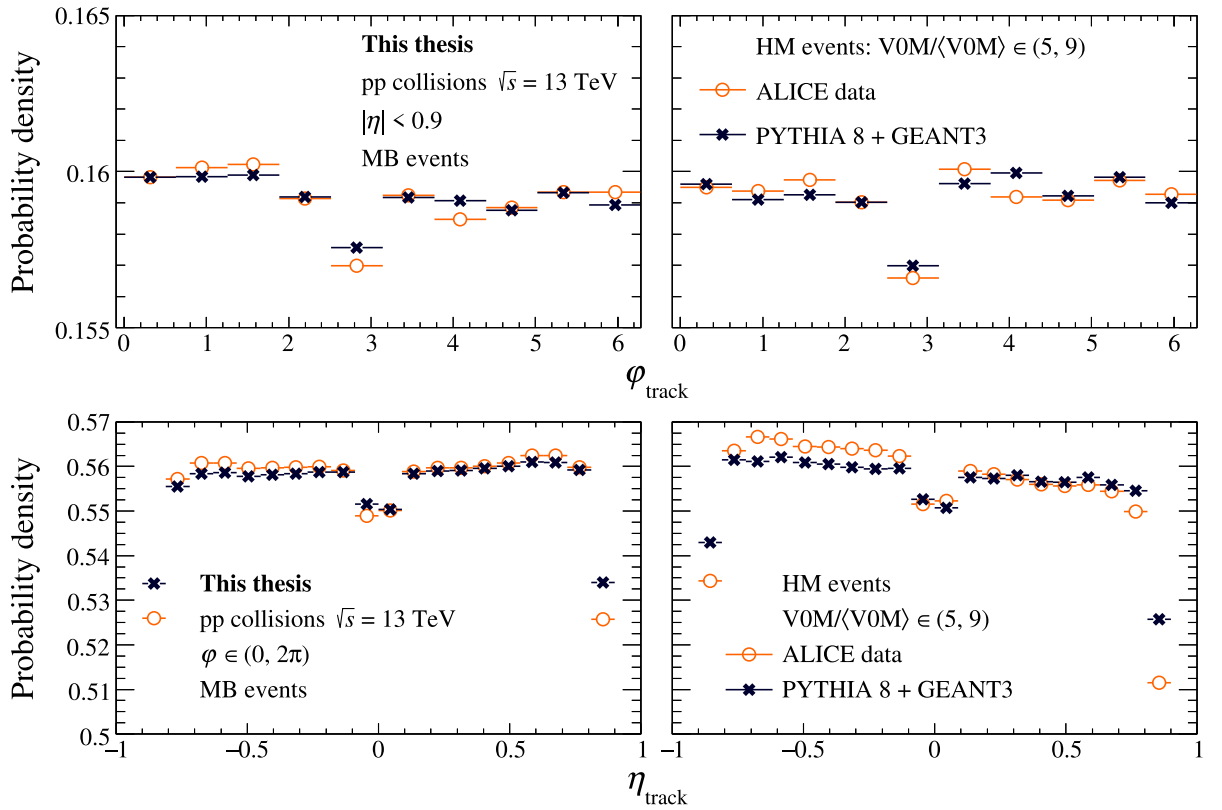


Figure 5.5: Azimuthal angle and pseudorapidity distributions of hybrid tracks reconstructed in raw data and detector-level PYTHIA 8 + GEANT3 events. Left: MB data. Right: HM data.

5.6 illustrates the track reconstruction efficiency for MB and HM events. The tracking efficiencies for MB and HM events are similar. The tracking efficiency for the hybrid tracks with $p_T = 1$ GeV/ c is about 85%, whilst it grows up to 82% at $p_T = 10$ GeV/ c and remains above 68% at higher p_T . The transverse-momentum resolution was obtained from

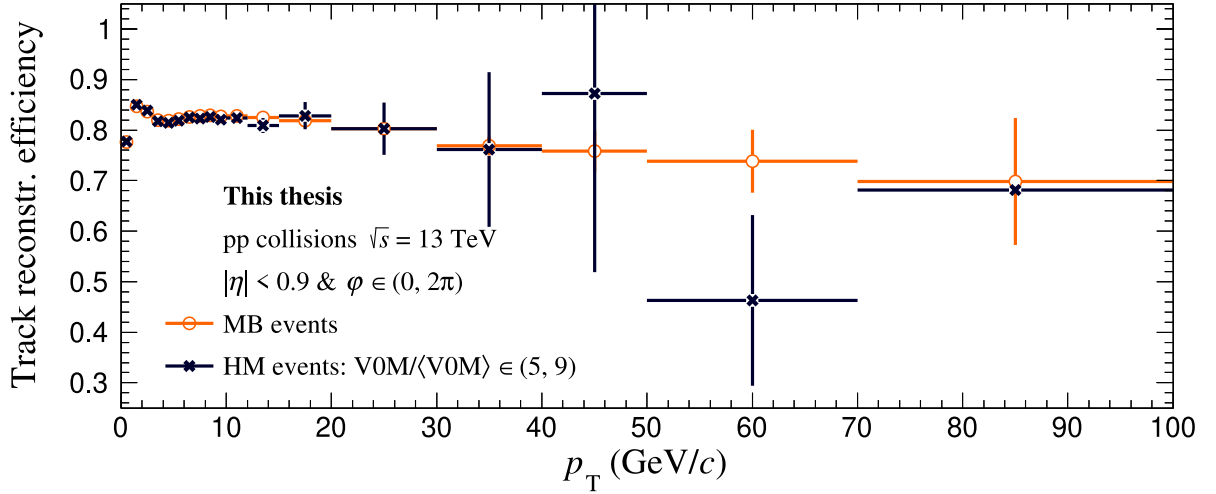


Figure 5.6: Track reconstruction efficiency in MB and HM pp collisions at $\sqrt{s} = 13$ TeV.

the track covariance matrix [189]. The mean relative transverse-momentum resolution of hybrid tracks is about 0.7% at $p_T = 1$ GeV/c and 1.3% at $p_T = 10$ GeV/c.

Jets were reconstructed from accepted hybrid tracks utilizing the anti- k_T algorithm with resolution parameter $R = 0.4$ and the boost-invariant p_T -recombination scheme [82]. The jets containing at least one constituent with $p_T > 100$ GeV/c were excluded from the analysis due to a poor momentum resolution for such tracks. The jet centroid was constrained within the fiducial volume $|\eta_{\text{jet}}| < 0.9 - R$ to assure the full jet reconstruction inside the ALICE central-barrel acceptance. Jet p_T was corrected for the presence of the underlying event following the approach discussed in Sect. 2.2.2. Jet area was calculated utilizing the “active area” approach implemented in the FastJet [82]. The mean value of ρ is 0.09 GeV/c for MB and 1.24 GeV/c for HM-selected events. For events containing a charged particle in $|\eta| < 0.9$ with $20 < p_T < 30$ GeV/c, the mean value of ρ is 0.39 GeV/c for MB and 1.62 GeV/c for HM events.

Jet reconstruction was also performed for the particle-level and detector-level MC simulations. The particle-level and detector-level jets were clustered from the corresponding samples of MC particles and hybrid tracks. The detector-level tracks had $|\eta| < 0.9$, $\varphi \in (0, 2\pi)$, and $p_T > 0.15$ GeV/c, whereas the particle-level tracks have the same constraints besides the cut on p_T .

5.4 Raw data

The goal of the analysis is the search for medium-induced jet scattering in HM compared to MB pp collisions at $\sqrt{s} = 13$ TeV, via broadening of the hadron-jet acoplanarity distribution. To this end, the $|\Delta\varphi|$ observable is introduced, which is the absolute azimuthal-angle separation between the trigger hadron and recoil jet. A significant source of background

in this search arises from multiple parton interactions (MPIs), in which two uncorrelated high- Q^2 scatterings occur within the same pp collision [202], with one producing a high- p_T trigger hadron and the other producing a recoil jet in the experimental acceptance. The acoplanarity distribution of such uncorrelated pairs is uniform, thereby restricting the sensitivity of the measurement to the broadening of the $|\Delta\varphi|$ distribution due to medium-induced scattering of recoil jets correlated with the trigger hadron.

Background yield from MPIs must be subtracted in a purely data-driven way for precise acoplanarity measurements. The Δ_{recoil} observable, introduced in Sect. 3.3.1, is particularly appropriate for this purpose. This analysis defines the Δ_{recoil} observable as a function of jet p_T and $|\Delta\varphi|$ azimuthal opening angle,

$$\Delta_{\text{recoil}} = \frac{1}{N_{\text{trig}}} \frac{d^2 N_{\text{jets}}}{dp_{T,\text{jet}}^{\text{ch, corr}} d|\Delta\varphi|} \Big|_{\text{TT}\{20,30\}} - c_{\text{Ref}} \times \frac{1}{N_{\text{trig}}} \frac{d^2 N_{\text{jets}}}{dp_{T,\text{jet}}^{\text{ch, corr}} d|\Delta\varphi|} \Big|_{\text{TT}\{6,7\}}, \quad (5.2)$$

where two classes of TT are introduced, which have $p_{T,\text{trig}}$ within the exclusive intervals $p_{T,\text{trig}} \in (6, 7)$ GeV/ c and $p_{T,\text{trig}} \in (20, 30)$ GeV/ c . These p_T intervals were chosen to optimize the opposing requirements of obtaining high-statistical precision and limiting the kinematic range for a more precise comparison of the same observables with different event activity selections. Hereafter, shorthand notations TT{20, 30} and TT{6, 7} will be adopted throughout this chapter. The azimuthal-angle separation between the TT and recoil jet was defined in the region $(0, \pi)$. The scaling factor c_{Ref} accounts for the different phase space available to observe the uncorrelated yield in the distributions associated with TT{20, 30} and TT{6, 7} [34], [35]. The c_{Ref} was extracted from the ratio of the first and second terms on the right-hand side of Eq. (3.4) in the bin $p_{T,\text{jet}}^{\text{ch, corr}} \in (0, 1)$ GeV/ c , see Fig. 5.7. In this region, jets have p_T compatible with the energy density ρ of the underlying event. The c_{Ref} values obtained for MB and HM events were 0.95 and 0.94, respectively. Differentiation of Eq. (5.2) with respect to η_{jet} was applied but omitted in the notations since the size of the pseudorapidity window for recoil jets in this analysis is $d\eta_{\text{jet}} = 1$.

After the offline selection, each of the MB and HM datasets was divided into two statistically independent subsets, where 95% of all events were assigned to search for TT{20, 30}, while the remaining 5% of events were utilized to search for TT{6, 7}. Any hybrid track reconstructed in $|\eta| < 0.9$ and $\varphi \in (0, 2\pi)$ was treated as a candidate for the TT if the $p_{T,\text{trig}}$ condition was satisfied. In rare cases, when several tracks satisfied the TT selection criteria, one was randomly chosen. The requirement of the TT presence in each event vastly reduced the analysis statistics. Table 5.1 summarizes the total number of events selected for the hadron-jet correlation measurement.

Figure 5.8 shows the projections of the raw trigger-normalized recoil jet yields and

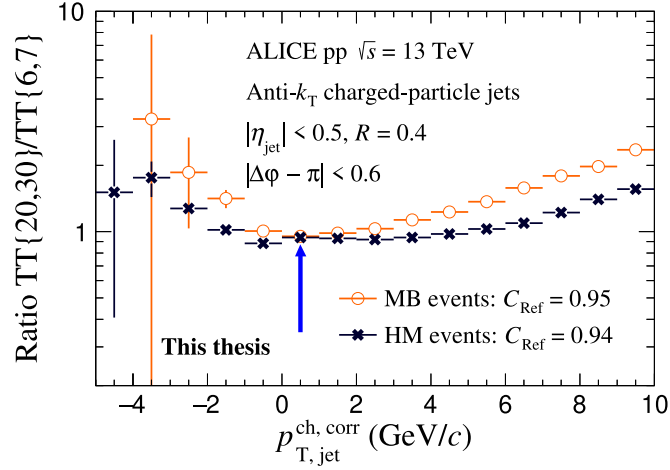


Figure 5.7: Ratio of recoil jet yields for $TT\{20,30\}$ and $TT\{6,7\}$ in MB and HM pp collisions at $\sqrt{s} = 13$ TeV. The arrow indicates the 0–1 GeV/ c bin, content of which is utilized to calculate c_{Ref} .

Data sample	$TT\{6,7\}$	$TT\{20,30\}$
MB	204.3k	62.9k
HM	506.0k	201.2k

Table 5.1: Total number of events remained for the analysis after searching for a trigger track.

the corresponding Δ_{recoil} distributions as a function of $p_{T,\text{jet}}^{\text{ch,corr}}$ and $|\Delta\varphi|$ for MB and HM pp events at $\sqrt{s} = 13$ TeV. The subtraction of uncorrelated background yield has a significant effect only in the region $p_{T,\text{jet}}^{\text{ch,corr}} < 20$ GeV/ c . The measured acoplanarity distributions exhibit a clear peak at $\Delta\varphi \approx \pi$, corresponding to the back-to-back topology of dijet events. Comparing the amplitudes of the Δ_{recoil} distributions for the MB and HM events, one can notice that the HM data are substantially suppressed. Furthermore, the HM acoplanarity spectra exhibit a broadening when compared to the corresponding MB distributions. Suppression of the recoil jet yield and hadron-jet medium-induced acoplanarity are signatures of jet quenching [38], [116], [182]. However, any interpretation of the measured acoplanarity distributions at this analysis stage is premature and inconclusive since the data are subject to the detector response. Thus, the raw data must be corrected for the detector effects via the unfolding procedure.

5.5 Data correction for detector effects

The physical properties of reconstructed jets are distorted by a collection of instrumental effects related to the detector response. The jet energy scale and resolution are strongly af-

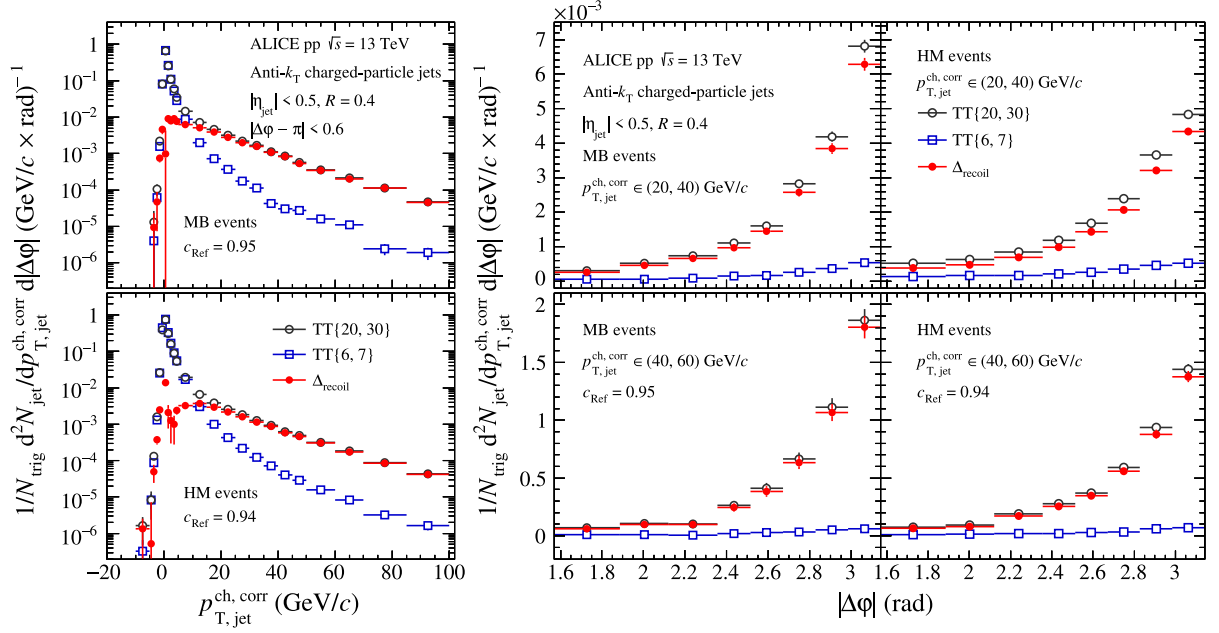


Figure 5.8: Raw trigger-normalized recoil jet distributions and the resulting Δ_{recoil} distribution for MB and HM pp collisions at $\sqrt{s} = 13$ TeV. Left panels: projection onto $p_{T,\text{jet}}^{\text{ch,corr}}$ for $|\Delta\varphi - \pi| < 0.6$. Right panels: projection onto $|\Delta\varphi|$ in two $p_{T,\text{jet}}^{\text{ch,corr}}$ intervals. Accessible in Ref. [39].

ected by limited track reconstruction efficiency and finite momentum resolution of tracks. The instrumental effects are encoded in the response matrix, R_{instr} , which maps the true distribution of Δ_{recoil} onto the measured distribution,

$$\Delta_{\text{recoil}}^{\text{Meas}}(|\Delta\varphi|^{\text{det}}, p_{T,\text{jet}}^{\text{det}}) = R_{\text{instr}}(|\Delta\varphi|^{\text{det}}, p_{T,\text{jet}}^{\text{det}}; |\Delta\varphi|^{\text{part}}, p_{T,\text{jet}}^{\text{part}}) \otimes \Delta_{\text{recoil}}^{\text{True}}(|\Delta\varphi|^{\text{part}}, p_{T,\text{jet}}^{\text{part}}), \quad (5.3)$$

where the notations *part* and *det* denote the particle-level and detector-level jets. Data correction for the instrumental effects is performed by solving Eq. (5.3) to acquire the true $\Delta_{\text{recoil}}^{\text{True}}$ spectrum for a given response matrix and the measured $\Delta_{\text{recoil}}^{\text{Meas}}$ spectrum. However, an endeavour to solve this equation for a singular R_{instr} via matrix inversion often leads to an incorrect true distribution, which can suffer from large fluctuations in central values. A physically appropriate solution of Eq. (5.3) is obtained via the unfolding procedure with imposed regularization. This analysis utilized the iterative 2D Bayesian unfolding technique [203] implemented in the RooUnfold package [204]. Besides the specification of the response matrix and measured distribution, the unfolding procedure also requires a prior spectrum, which serves as an initial presumption about the shape of the true $\Delta_{\text{recoil}}^{\text{True}}$ distribution. The subsequent sections will describe all requisite inputs for the unfolding in more detail. Throughout this chapter, particle-level and detector-level distributions,

as well as raw data, will be presented using the following rebinning scheme:

$$\begin{aligned}
 p_{\text{T, jet}}^{\text{ch, corr}} &= \begin{cases} (10, 15, 20, 25, 30, 35, 40, 50, 60, 100) \text{ GeV}/c & \text{detector-level/raw data} \\ (0, 10, 20, 30, 40, 50, 60, 80, 100, 130) \text{ GeV}/c & \text{particle-level data} \end{cases}, \\
 |\Delta\varphi| &= \left(\frac{20}{40}\pi, \frac{24}{40}\pi, \frac{28}{40}\pi, \frac{32}{40}\pi, \frac{34}{40}\pi, \frac{36}{40}\pi, \frac{38}{40}\pi, \pi \right) \text{ rad} \quad \text{particle-/detector-level data.}
 \end{aligned} \tag{5.4}$$

5.5.1 Response matrix

The instrumental response matrix is formed by pairs of the inclusive particle-level and detector-level jets reconstructed in simulated MB events, which have centroids separated by the minimum phase-space angular distance,

$$\min \left\{ \sqrt{(\eta^{\text{part}} - \eta^{\text{det}})^2 + (\varphi^{\text{part}} - \varphi^{\text{det}})^2} \right\} \leq \Delta R_{\text{min}}, \tag{5.5}$$

where ΔR_{min} is the maximum allowed angular distance between a pair of the matched jets. The value of ΔR_{min} was set to 0.3 for this analysis. The matching with ghost jets (see Sect. 2.2.2) was prohibited. No additional criterion was imposed on the matching between jets' constituents. The successful matching assures the injective mapping of detector-level jets to particle-level jets. Unmatched particle-level jets were considered as “missed” jets, and their yield was utilized to estimate jet reconstruction efficiency. Figure 5.9 shows the jet reconstruction efficiency, which gives the probability of reconstructing the particle-level jet after it passes the detector response simulations. The efficiency of jet reconstruction is 98% at $p_{\text{T, jet}} = 10 \text{ GeV}/c$ and almost 100% at $p_{\text{T, jet}} \in (15, 100) \text{ GeV}/c$.

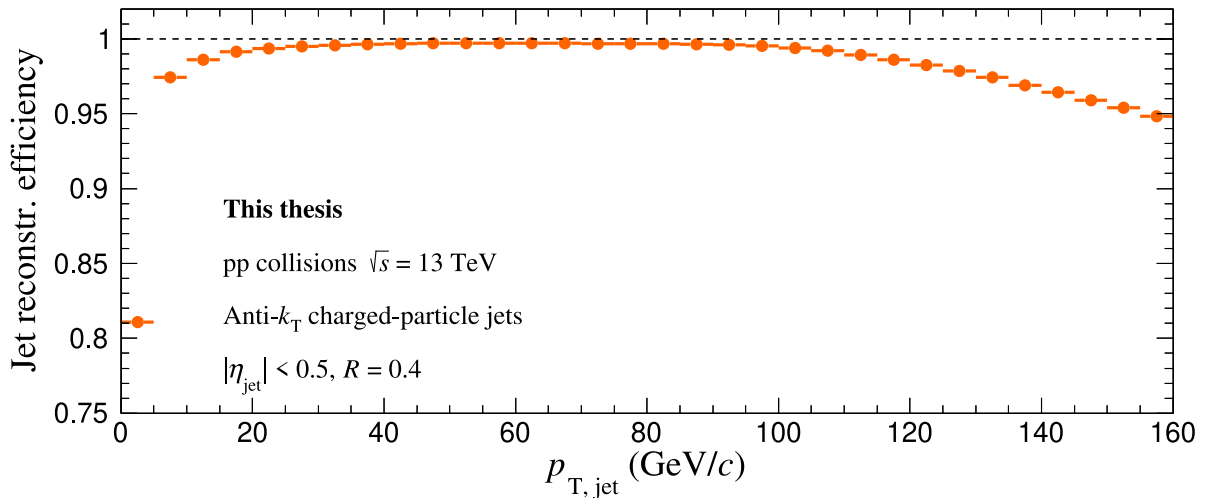


Figure 5.9: Reconstruction efficiency of the particle-level jets after applying the instrumental effects in pp collisions at $\sqrt{s} = 13 \text{ TeV}$.

A drop in the efficiency above 100 GeV/ c is due to the rejection of detector-level jets, which had at least one constituent with $p_T > 100$ GeV/ c .

The choice of the inclusive jet population instead of recoil jets for response matrix determination is motivated by the larger statistics of this jet sample. It was verified that the azimuthal angle and p_T smearing of recoil jets are close to that of inclusive jets. Figure 5.10 shows the smearing of $p_{T, \text{jet}}$ due to instrumental effects for inclusive jets and recoil jets associated with $\text{TT}\{20, 30\}$ in MB events. The momentum smearing of

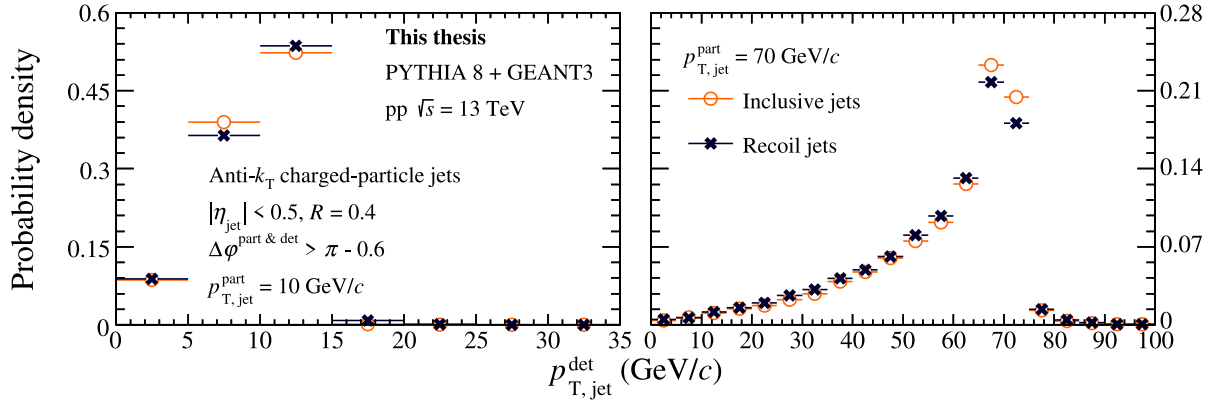


Figure 5.10: Momentum smearing of inclusive jets and recoil jets associated with $\text{TT}\{20, 30\}$ due to detector response. The distributions have a maximum at the p_T values corresponding to the selected $p_{T, \text{jet}}^{\text{part}}$. See text for further details.

inclusive jets has $\text{RMS} \approx 2.43$ GeV/ c at $p_T = 10$ GeV/ c and $\text{RMS} \approx 15.68$ GeV/ c at $p_T = 70$ GeV/ c , whereas the momentum smearing of recoil jets has $\text{RMS} \approx 2.64$ GeV/ c and $\text{RMS} \approx 16.02$ GeV/ c for the same p_T , respectively. The RMS of the azimuthal angular smearing of inclusive jets in MB events decreases from 2.6×10^{-2} rad at 10 GeV/ c to 10^{-2} rad at 70 GeV/ c . Similar RMS values are observed for recoil jets, namely 3.4×10^{-2} rad and 1.3×10^{-2} rad for the same p_T , respectively. The difference is minor, thereby, with sufficient accuracy for this analysis, one can consider utilizing inclusive jets to construct the response matrix. Note that the closure test (same as in Sect. 5.5.3) was carried out using both matrices. It was found that the difference between the results was within statistical uncertainties.

The correction of the acoplanarity $\Delta_{\text{recoil}}(|\Delta\varphi|)$ distribution for instrumental effects has to take into account not only the $|\Delta\varphi|$ smearing but also the jet p_T smearing, which has a larger impact on the measured spectrum since the p_T spectrum of recoil jets is steeply falling, see Fig. 5.8. Thus, the unfolding procedure requires the instrumental response matrix, which has four axes ($|\Delta\varphi|^{\text{det}}, p_{T, \text{jet}}^{\text{det}}, |\Delta\varphi|^{\text{part}}, p_{T, \text{jet}}^{\text{part}}$), where the particle-level and detector-level $p_{T, \text{jet}}$ are corrected for underlying event density using Eq. (2.11). The smearing of $|\Delta\varphi|$ is approximated by the azimuthal angular smearing of inclusive jets based on the arguments discussed in the paragraph above. However, the $|\Delta\varphi|$ is defined

in the range $(0, \pi)$, while the azimuthal angle of inclusive jets φ is distributed uniformly in the range $(0, 2\pi)$. To avoid loss of statistics due to different ranges of the azimuthal angle of recoil and inclusive jets, the particle- and detector-level values of jet φ angles to be filled to the response matrix are obtained as follows:

- Azimuthal angle smearing due to detector effects is obtained as $\delta\varphi = \varphi_{\text{jet}}^{\text{det}} - \varphi_{\text{jet}}^{\text{part}}$, where $\varphi_{\text{jet}}^{\text{det}}$ and $\varphi_{\text{jet}}^{\text{part}}$ denote the azimuthal angles of the matched inclusive jets;
- Particle-level $|\Delta\varphi|^{\text{part}}$ angle is sampled uniformly in the range $(0, \pi)$ rad;
- Detector-level $|\Delta\varphi|^{\text{det}}$ is then obtained by the smearing of the $|\Delta\varphi|^{\text{part}}$ angle with $\delta\varphi$, $|\Delta\varphi|^{\text{det}} = |\Delta\varphi|^{\text{part}} + \delta\varphi$. The $|\Delta\varphi|^{\text{det}}$ angle is constrained in the range $(0, \pi)$ rad with no underflow/overflow bins.

Figures 5.11 (a) and (b) show the projections of the generated instrumental response matrix onto $|\Delta\varphi|^{\text{det}}$ vs. $|\Delta\varphi|^{\text{part}}$ plane when $p_{\text{T, jet}}^{\text{det}}$ and $p_{\text{T, jet}}^{\text{part}}$ are constrained in the range 20–40 GeV/c and 40–60 GeV/c, respectively. Figure 5.11 (c) shows the projection onto

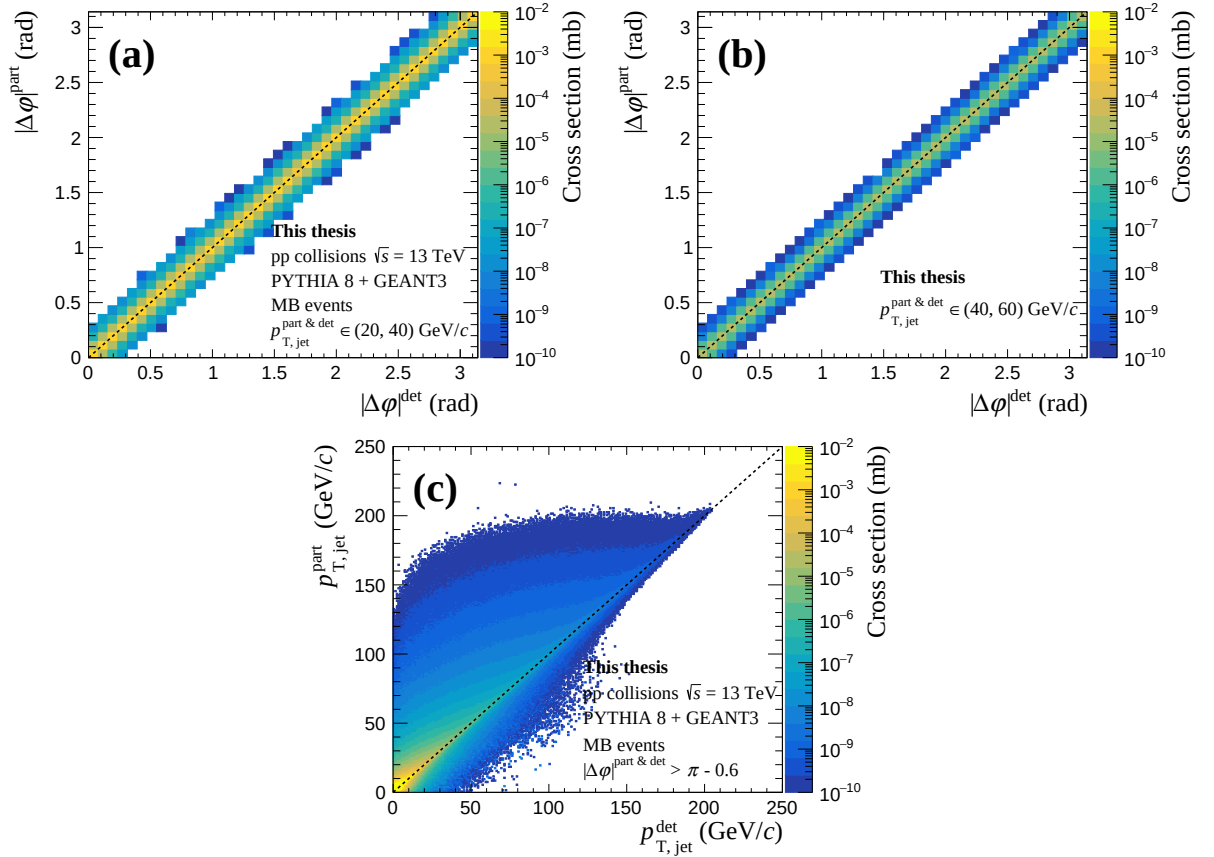


Figure 5.11: Projections of the 4D instrumental response matrix. Top panels: projection onto $|\Delta\varphi|^{\text{det}}$ vs. $|\Delta\varphi|^{\text{part}}$ plane in two intervals of $p_{\text{T, jet}}^{\text{det}}$ and $p_{\text{T, jet}}^{\text{part}}$. Bottom panel: projection onto $p_{\text{T, jet}}^{\text{det}}$ vs. $p_{\text{T, jet}}^{\text{part}}$ plane with 1 GeV/c binning for $|\Delta\varphi|^{\text{part}} \& \text{det} > \pi - 0.6$. The dashed line represents the main diagonal.

$p_{T,\text{jet}}^{\text{det}}$ vs. $p_{T,\text{jet}}^{\text{part}}$ plane for $|\Delta\varphi| > \pi - 0.6$ at the particle- and detector-level. The presented projections have a close-to-diagonal form. One can see that the smearing of the azimuthal angle is relatively tiny compared to the jet p_T smearing. Therefore, the main correction of the $\Delta_{\text{recoil}}(|\Delta\varphi|)$ distribution is associated with jet p_T smearing.

5.5.2 Prior spectrum

A prior spectrum has to be chosen close enough to the expected shape of the true distribution. Moreover, a prior spectrum should be smooth to suppress the impact of random fluctuations on the unfolded distribution. This analysis utilizes the Δ_{recoil} prior spectra obtained for MB and HM pp events at $\sqrt{s} = 13$ TeV simulated by the PYTHIA 8 event generator with the Monash tune [64]. The statistics of PYTHIA HM events were increased by relaxing the event activity selection criterion to $V0M/\langle V0M \rangle \in (4, 9)$. Figure 5.12 shows the projections of the prior Δ_{recoil} distributions for MB and HM events onto jet- p_T and $|\Delta\varphi|$ axes.

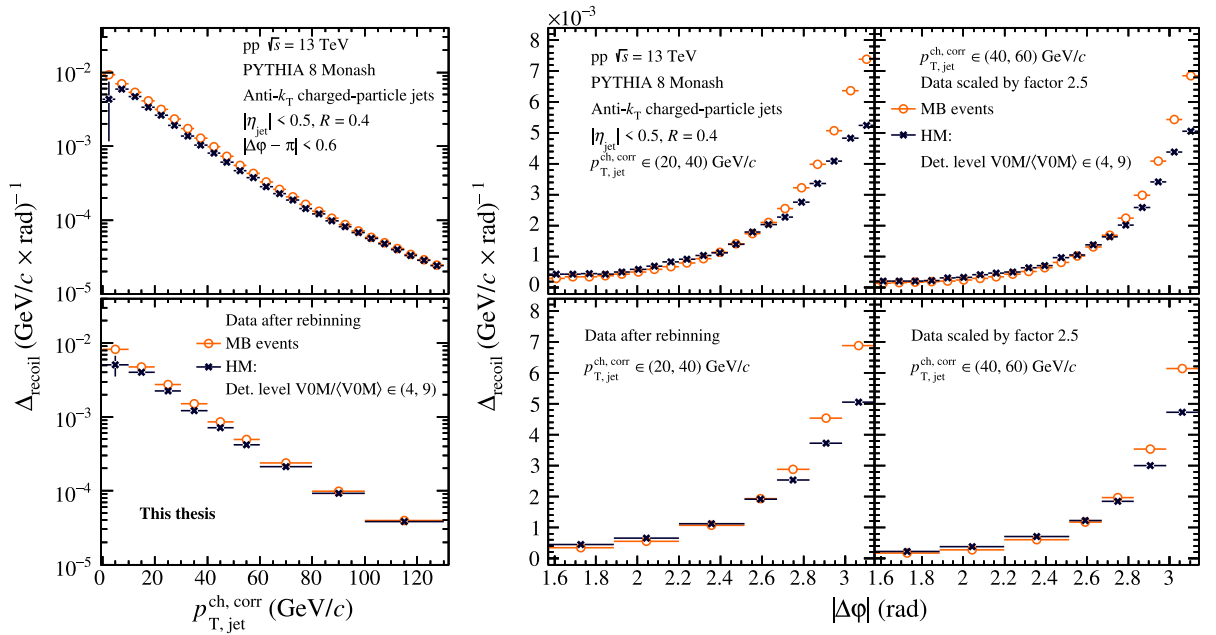


Figure 5.12: Prior Δ_{recoil} distributions for MB and HM pp events at $\sqrt{s} = 13$ TeV simulated by the PYTHIA 8 Monash [64]. The distributions are projected onto the $p_{T,\text{jet}}^{\text{ch, corr}}$ and $|\Delta\varphi|$ axes for restricted ranges in the other kinematic variable. Top panels: data before rebinning. Bottom panels: data after rebinning according to the scheme (5.4). Acoplanarity distributions for jets with $p_{T,\text{jet}}^{\text{ch, corr}} \in (40, 60)$ GeV/c are scaled by a factor of 2.5 for better visibility. The HM prior was obtained by applying a constraint on the detector-level $V0M/\langle V0M \rangle \in (4, 9)$ signal.

The Δ_{recoil} prior spectrum for HM events was defined by constraining the detector-level $V0M/\langle V0M \rangle$ signal. In this case, one properly accounts for the contribution of secondary generated particles, which affect the measured $V0M$ signal (see Fig. 5.4). These particles

are not present in the particle-level simulations.

5.5.3 Closure test

The unfolding procedure must be verified using a closure test. In this test, the simulated detector-level Δ_{recoil} distribution is unfolded by a regularized inversion of Eq. (5.3). The response matrix, which appears in Eq. (5.3), is the same as described in Sect. 5.5.1. Since the true Δ_{recoil} distribution is known in advance, the resulting unfolded distribution can be directly compared with it. The completed closure test is considered passed if the unfolded distribution is consistent with the true one within statistical uncertainties.

The raw data have limited statistical precision, directly affecting the unfolding results. In the case of PYTHIA 8 simulations, high- p_T jets, which emerge from rare events with large Q^2 transfer, can be efficiently generated by means of phase space constraints, known as $\hat{p}_T$ hard-bins [199]. However, in this case, the generated spectrum has an unrealistic size of statistical uncertainty. Therefore, it is requisite to adjust the size of statistical uncertainties in the detector-level spectrum to reflect the statistical precision of the raw data. To this end, the trigger-normalized 2D distributions of recoil jets for MB and HM events at the detector level, associated with TT{20, 30} and TT{6, 7}, were re-scaled by the corresponding number of trigger tracks from the raw MB and HM data,

$$N_{\text{trig}}^{\text{real data}} \times \left[\frac{1}{N_{\text{trig}}} N_{\text{jets}} \left(p_{T, \text{jet}}^{\text{ch, corr}}, |\Delta\varphi| \right) \right]^{\text{detector level}} \Big|_{p_T \in \text{TT}\{x,y\}}. \quad (5.6)$$

The detector-level HM events in the closure test were required to pass a cutoff on the detector-level $V0M/\langle V0M \rangle \in (4, 9)$. This choice provides sufficiently smooth input distributions. The number of trigger tracks available in real MB and HM events is presented in Tab 5.1. The obtained detector-level distributions were subsequently randomized bin-by-bin by sampling from the Poisson distribution with the mean value set to the actual bin content. The corresponding statistical error was then evaluated by taking the square root of the sampled bin content. Further, the randomized trigger-normalized distributions of recoil jets for MB and HM events were used to obtain the corresponding detector-level Δ_{recoil} spectra. Figure 5.13 compares the detector-level Δ_{recoil} distributions obtained from the original trigger-normalized detector-level spectra and the corresponding spectra after randomization.

The randomization yielded 20 instances of the randomized detector-level Δ_{recoil} spectra both for MB and HM events. The closure test was repeated for each of these instances. Prior distributions for MB and HM events were the same as described in Sect. 5.5.2. These prior spectra were concurrently the true spectra for the closure test. Unfolding of each of the generated distributions was carried out utilizing the 2D Bayesian unfolding

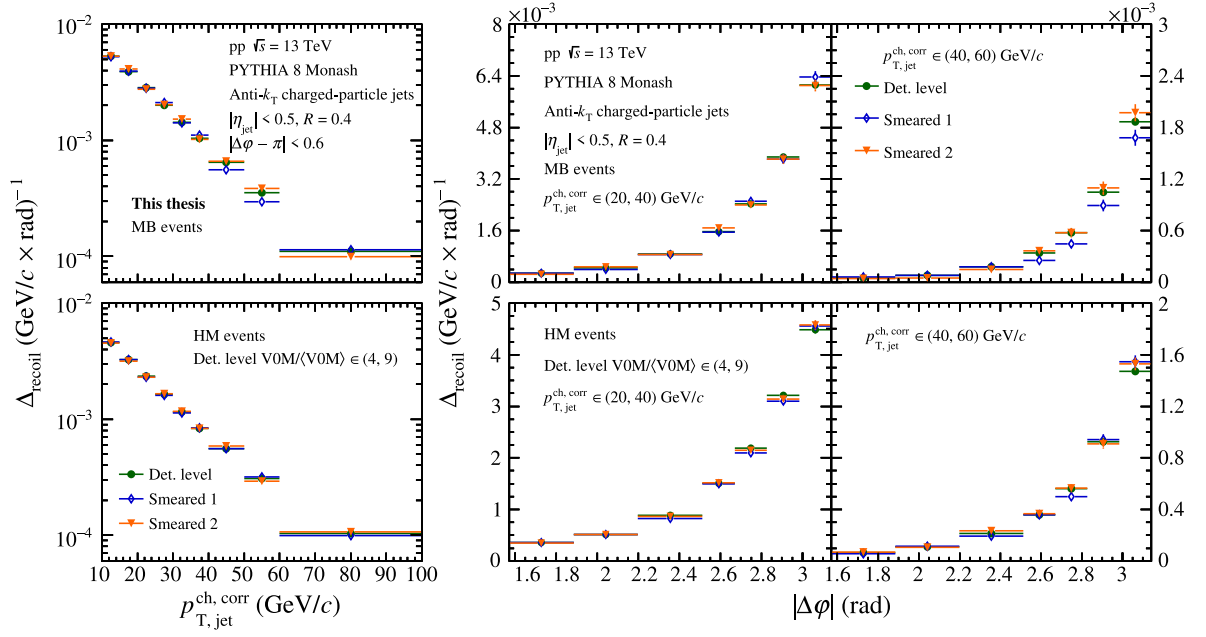


Figure 5.13: Projections of the detector-level Δ_{recoil} distributions as obtained from randomized per-trigger normalized spectra. The original detector-level spectrum is shown with dark green markers for comparison. Top panels: MB events. Bottom panels: HM events. See text for further details.

method. The number of iterations varied from 2 to 10. The resulting unfolded distributions for MB and HM events at each iteration step were obtained as the median of the 20 corresponding unfolded distributions. These spectra were then compared with the corresponding true spectra. All distributions involved in the closure test were rebinned according to the scheme from Eq. (5.4). The outcome of the closure test for MB and HM Δ_{recoil} distributions is shown in Fig. 5.14 and Fig. 5.15. The unfolding solutions had a rapid convergence to a stable result. Thus, the 5th iteration was considered to be the optimal choice both for MB and HM analyses. The gray area represents the relative statistical uncertainty of the median spectrum (1σ) obtained with five iterations of Bayesian unfolding. The unfolded spectra are comparable with the true distribution within the statistical uncertainties. Additionally, the unfolding results were verified using the folding procedure, where the obtained unfolded spectra were convoluted with the response matrix in accordance with Eq. (5.3) to get back the input spectrum. The folding results demonstrated the consistency of the folded spectra with the input distribution within the statistical uncertainties, see Appendix A. By summarizing the results of the closure test, one can confidently conclude that it was successful. Thus, the generated MC data sample can be used to unfold raw data.

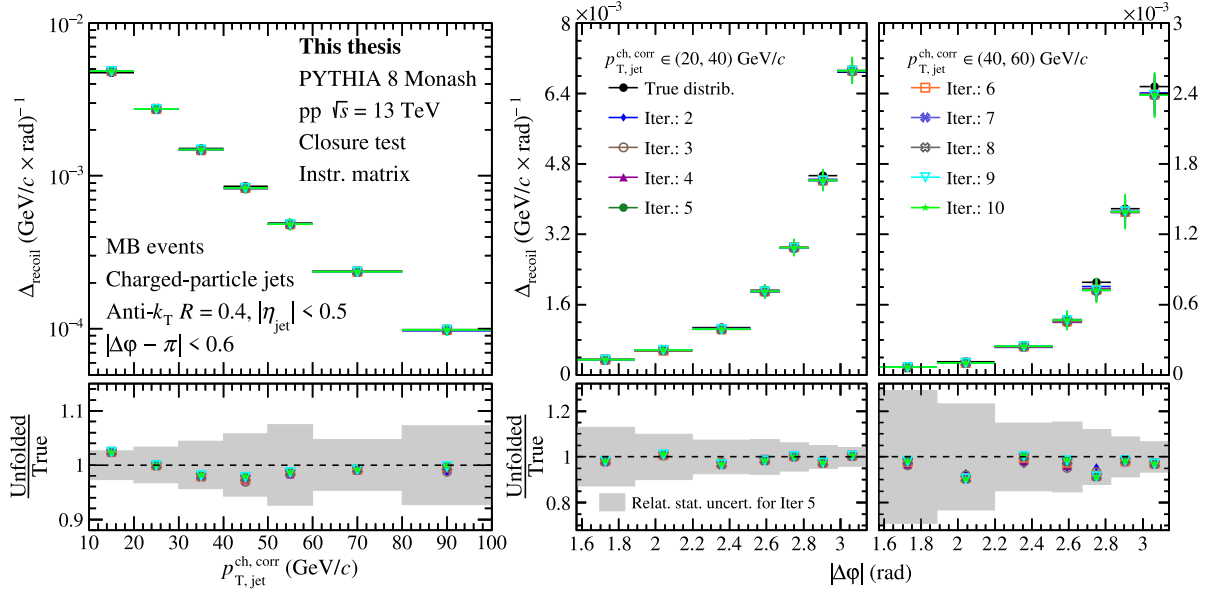


Figure 5.14: Results of the closure test for Δ_{recoil} distributions obtained in MB PYTHIA 8 events. Color markers show the unfolded solutions obtained at different iteration steps. The bottom panel of each plot shows a ratio of the unfolded distributions to the true spectrum. The gray bands represent the relative statistical uncertainty of the unfolded spectrum obtained in the 5th iteration. Left plot: projections of the Δ_{recoil} distribution onto the $p_{T, \text{jet}}^{\text{ch, corr}}$ axis for $|\Delta\phi - \pi| < 0.6$. Middle and right plots: projections of the Δ_{recoil} distribution onto the $|\Delta\phi|$ axis in two exclusive $p_{T, \text{jet}}^{\text{ch, corr}}$ intervals.

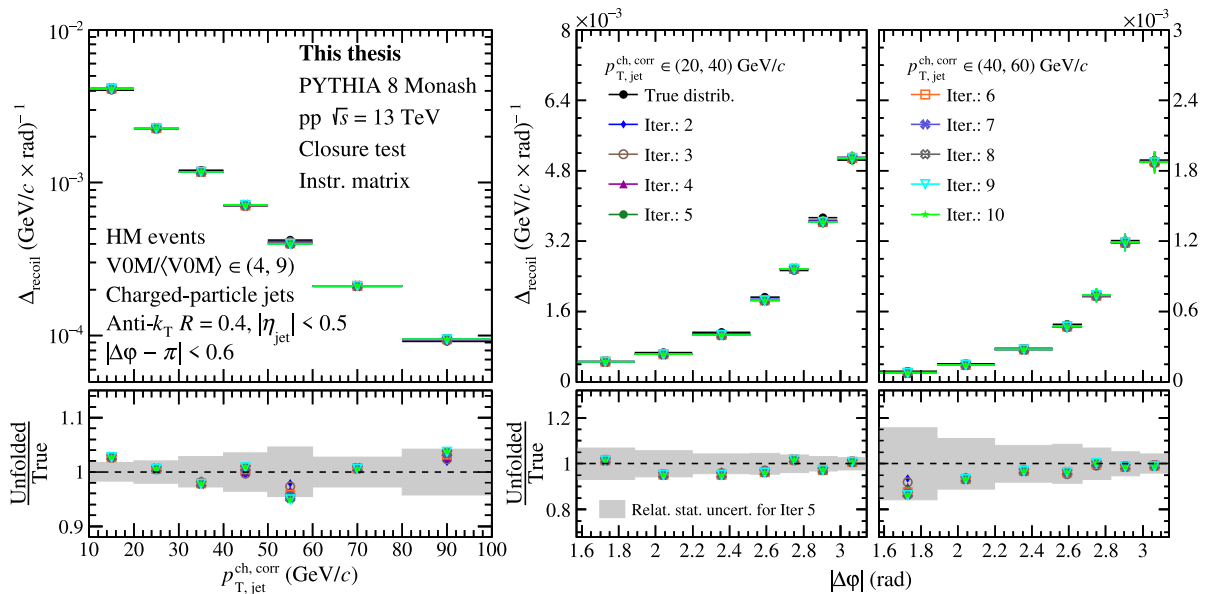


Figure 5.15: Same as Fig. 5.14, for HM PYTHIA 8 events.

5.5.4 Correction for background fluctuations

The background subtraction approach discussed in Sect. 2.2.2 does not account for jet p_T smearing due to local background fluctuations. Therefore, in addition to instrumental effects, the Δ_{recoil} distributions should also be corrected for background fluctuations. To simultaneously correct data for background fluctuations and instrumental effects, the unfolding procedure is performed utilizing the combined response matrix R_{comb} , which is supposed to factorize as a product of two independent response matrices describing background fluctuations and instrumental effects [36],

$$R_{\text{comb}} = R_{\text{bkgd}}(p_{T,\text{jet}}^{\text{reco}}; p_{T,\text{jet}}^{\text{det}}) \otimes R_{\text{instr}}(|\Delta\varphi|^{\text{det}}, p_{T,\text{jet}}^{\text{det}}; |\Delta\varphi|^{\text{part}}, p_{T,\text{jet}}^{\text{part}}), \quad (5.7)$$

where $p_{T,\text{jet}}^{\text{reco}}$ is jet p_T smeared by background fluctuations and instrumental effects. To investigate local background fluctuations in an unbiased way, the random cone method (RC) was employed [91]. A cone with a fixed size $R = 0.4$ was placed in each event with TT{20, 30} at random η and φ . The cone centroid was constrained to be within the ALICE fiducial volume $|\eta| < 0.5$. Further, the sum of tracks p_T , aggregated inside the placed cone, was compared to the mean energy density of the underlying event ρ ,

$$\delta p_T = \sum^{\text{RC}} p_{T,\text{track}} - \rho \times \pi R^2, \quad (5.8)$$

where the δp_T distribution characterizes the jet p_T smearing due to local background fluctuations. To eliminate the contribution of hard-scattering processes on the estimate of δp_T , the cone did not overlap with the leading and sub-leading jets, and the TT. Figure 5.16 shows the δp_T distributions obtained in MB and HM pp events at $\sqrt{s} = 13$ TeV simulated by PYTHIA 8 Monash + GEANT3. The R_{bkgd} matrices for MB and HM events were constructed by arranging the corresponding δp_T distributions along the main diagonal with the 0 GeV/ c bin on the diagonal, see Fig. 5.17. Finally, the combined matrices constructed according to Eq. (5.7) were used in the statistical closure test discussed in Sect. 5.5.3. Figures 5.18 and 5.19 show the outcomes of the closure test for MB and HM events, respectively. The unfolded Δ_{recoil} distributions are consistent with the true spectra within the statistical uncertainty. However, the distributions appear to be shifted by a constant factor when compared to the results of the closure test with the instrumental matrix, see Fig. 5.14. Figure 5.20 shows the ratio of the Δ_{recoil} distributions unfolded using the combined and instrumental matrices. The ratios are permanently below unity. Thereby, one may conclude that the closure test with the combined matrix provides worse results when compared to the analogous test with the instrumental matrix. Hence, in the primary analysis, the data will be unfolded using the instrumental matrix

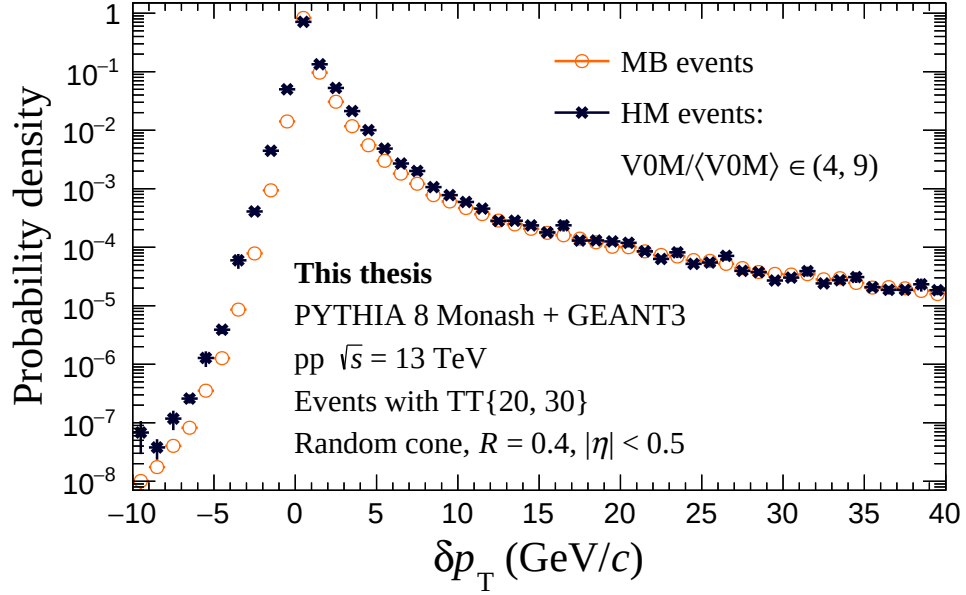


Figure 5.16: δp_T smearing due to local background fluctuations measured using the random cone method in MB and HM pp collisions at $\sqrt{s} = 13$ TeV simulated by PYTHIA 8 Monash + GEANT3.

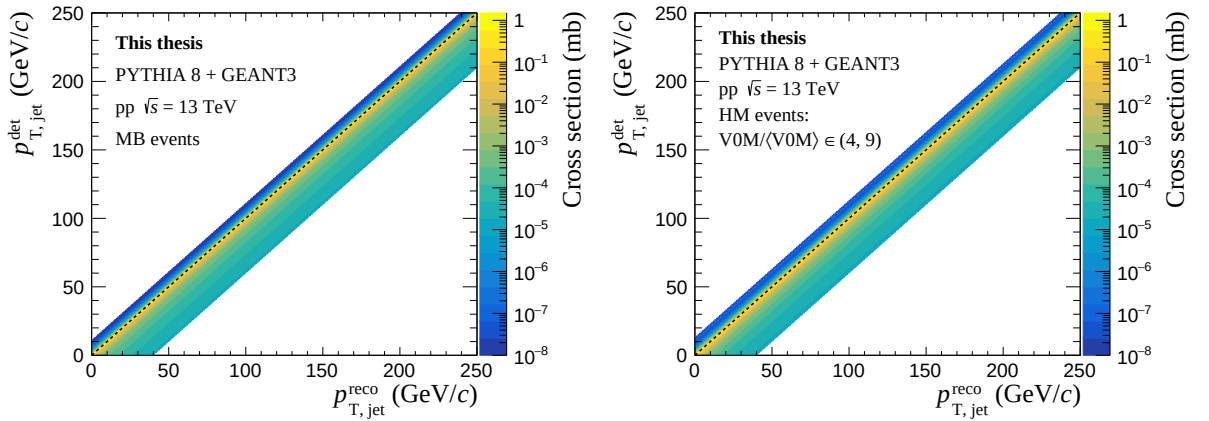


Figure 5.17: R_{bkgd} response matrices for MB and HM pp collisions at $\sqrt{s} = 13$ TeV simulated by PYTHIA 8 Monash + GEANT3. The dashed line represents the main diagonal.

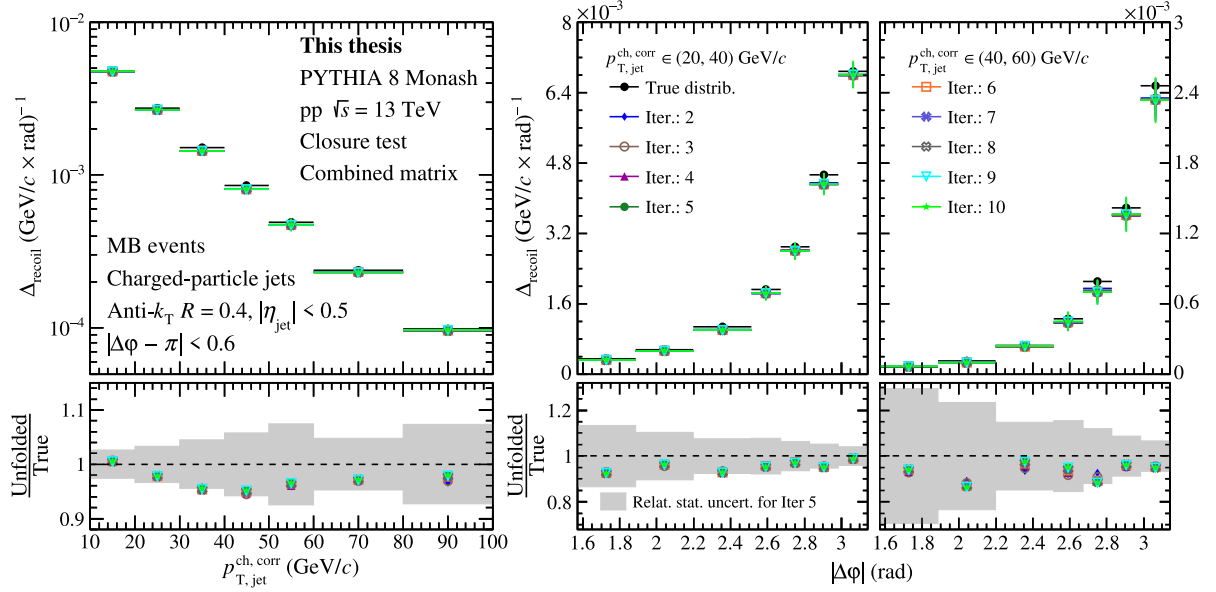


Figure 5.18: Results of the closure test with the combined matrix R_{comb} for Δ_{recoil} distributions obtained in MB PYTHIA 8 events. The meaning is the same as for Fig. 5.14.

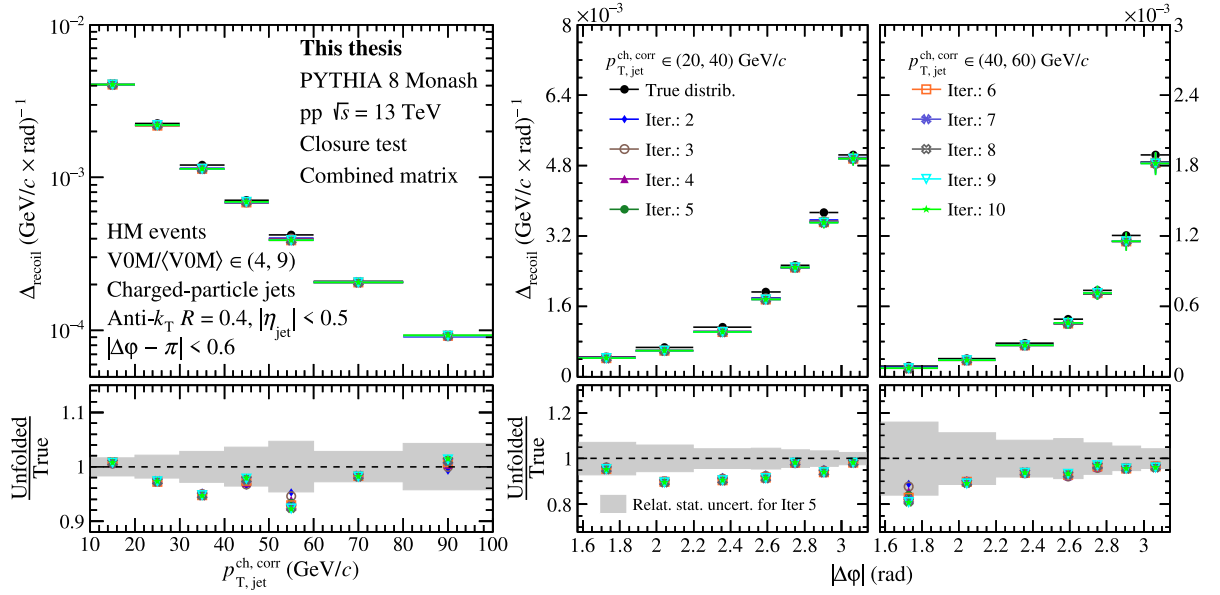


Figure 5.19: Same as Fig. 5.18, for HM PYTHIA 8 events.

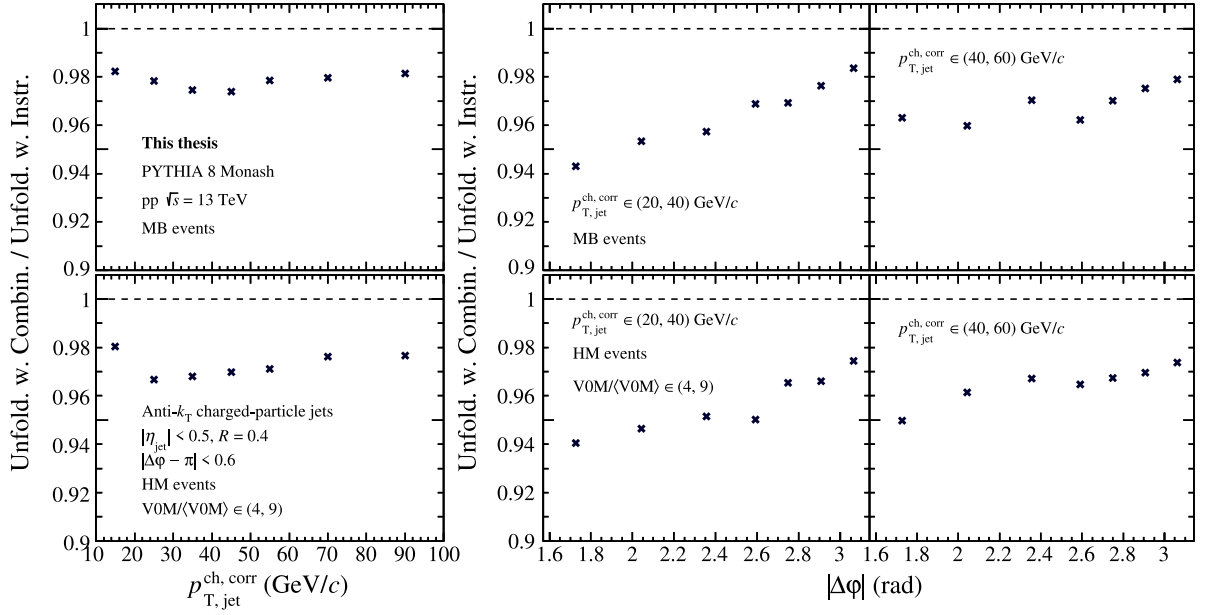


Figure 5.20: Ratio of the Δ_{recoil} distributions for MB (top) and HM (bottom) PYTHIA 8 events, which were unfolded using the combined matrix R_{comb} and the instrumental matrix R_{instr} . Statistical error bars are not shown.

R_{instr} , while the data unfolding with the combined matrix R_{comb} will be treated as a systematic variation, see Sect. 5.6.5.

5.5.5 Data unfolding

The raw Δ_{recoil} spectra for MB and HM events are shown in Fig. 5.8. The response matrix and prior spectra, which are requisite inputs to unfolding, were discussed in Sect. 5.5.1 and 5.5.2, respectively. The raw spectra were unfolded utilizing the 2D Bayesian unfolding method. The outcomes of the unfolding are shown in Fig. 5.21 and 5.22. Error bars in the unfolded distributions were calculated automatically by the RooUnfold package. The folding results are shown in Appendix B. The ratio plot explicitly demonstrates the convergence of the obtained unfolded solutions. The optimal number of iterations was found to be 5 for both the MB and HM analyses based on a requirement that unfolded distributions from successive iterations exhibit a mean change of less than 2% when averaged over all bins.

5.6 Systematic uncertainties

The corrected Δ_{recoil} spectra for MB and HM events can be affected by the following sources of systematic uncertainty:

- uncertainty on track reconstruction efficiency;

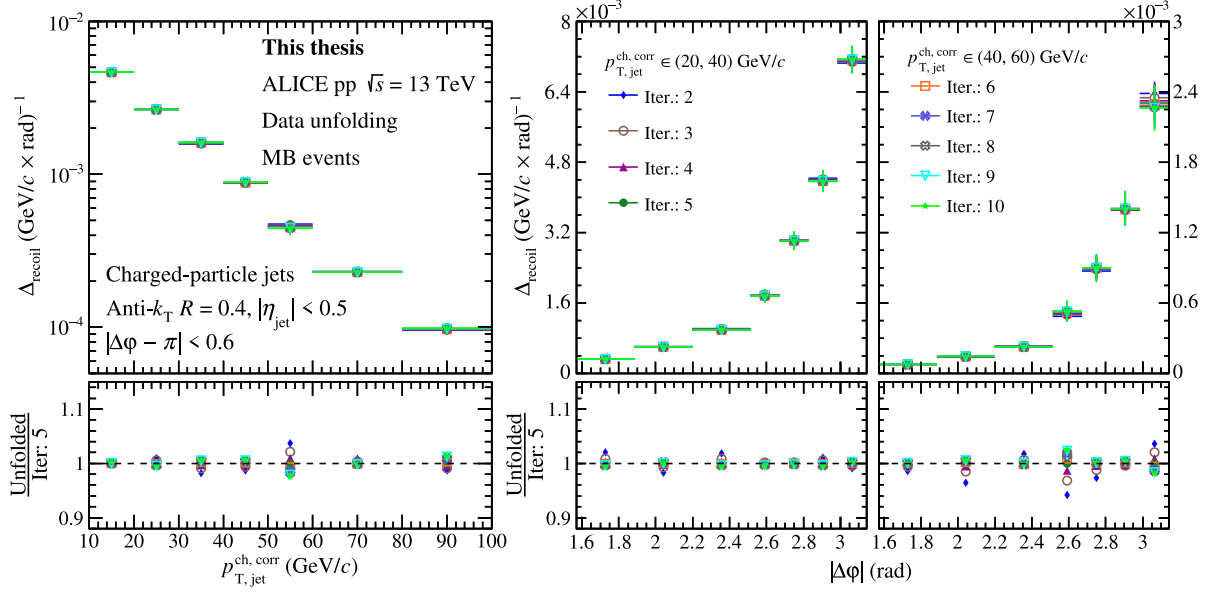


Figure 5.21: Unfolded Δ_{recoil} distributions measured in MB pp collisions at $\sqrt{s} = 13$ TeV. The unfolded distributions are shown for different iteration steps. The bottom plots contain the ratio of these distributions to the spectrum obtained for the 5th iteration. Left plot: projections of the Δ_{recoil} distribution onto the $p_{T, \text{jet}}^{\text{ch, corr}}$ axis for $|\Delta\phi - \pi| < 0.6$. Middle and right plots: projections of the Δ_{recoil} distribution onto the $|\Delta\phi|$ axis in two exclusive $p_{T, \text{jet}}^{\text{ch, corr}}$ intervals.

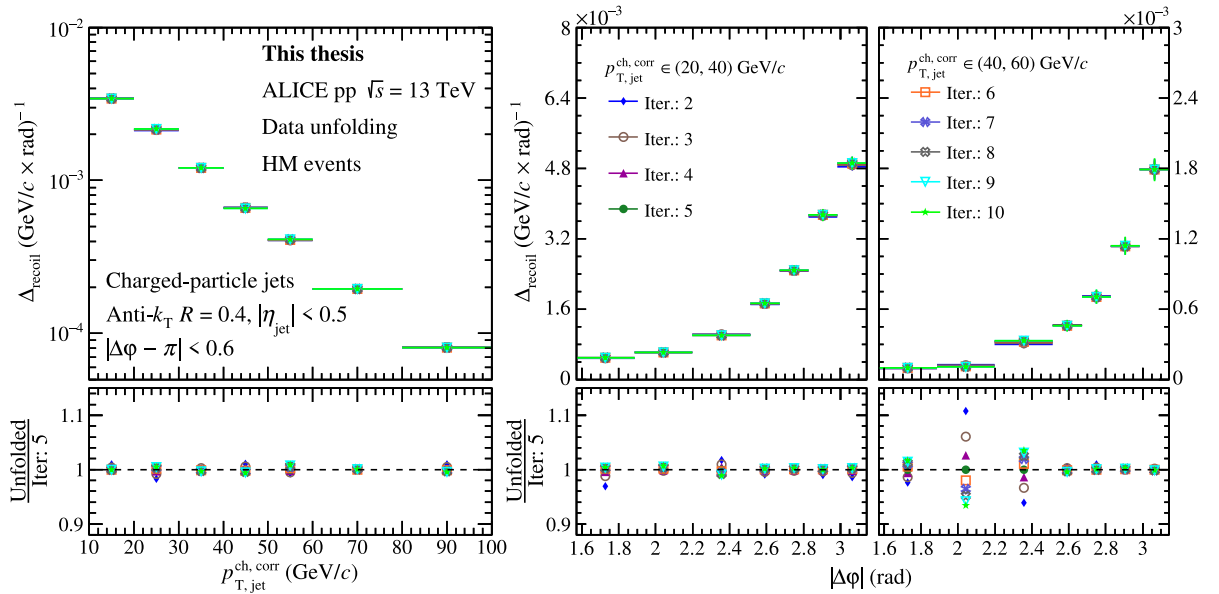


Figure 5.22: Same as Fig. 5.21, for HM events.

- uncertainty on track p_T resolution;
- unfolding settings;
- method for estimation of the scaling factor c_{Ref} ;
- residual background fluctuations.

Subsequent sections will discuss each item in more detail. The systematic uncertainty due to each source was determined by varying the corresponding parameters and rerunning the entire analysis chain. The considered sources were assumed to be statistically independent; therefore, the total systematic uncertainty was obtained by summing up the contribution from each source in quadrature.

To suppress the influence of statistical fluctuations on the estimated systematic uncertainties, the raw Δ_{recoil} spectra for MB and HM events were randomized similarly to what was done in Sect. 5.5.3. In total, nine statistically independent copies of the initial Δ_{recoil} spectra were generated. Each randomized distribution was also unfolded with the setting of the *primary analysis* (PA) and *varied settings*. The ratio of the spectra unfolded in the analysis with varied settings (Var., shorthand from “Varied”) of a given source to the unfolded spectra from the PA was calculated. Then, systematic uncertainty for this source was taken as the median of this set of ratios. Tables 5.2 and 5.3 show representative values of the systematic uncertainty in $\Delta_{\text{recoil}}(p_{T,\text{jet}}^{\text{ch, corr}})$ and $\Delta_{\text{recoil}}(|\Delta\varphi|)$ measurements for MB and HM-selected event populations, respectively.

Projection	Relative systematic uncertainty (%)					
	$\Delta_{\text{recoil}}(p_{T,\text{jet}}^{\text{ch, corr}})$ $ \Delta\varphi - \pi < 0.6$		$\Delta_{\text{recoil}}(\Delta\varphi)$ $p_{T,\text{jet}}^{\text{ch, corr}} \in (20, 40) \text{ GeV}/c$		$\Delta_{\text{recoil}}(\Delta\varphi)$ $p_{T,\text{jet}}^{\text{ch, corr}} \in (40, 60) \text{ GeV}/c$	
	10–20 GeV/ c	60–80 GeV/ c	$2\pi/3$	π	$2\pi/3$	π
Bin						
Tracking efficiency	0.2	7.1	3.7	1.9	8.4	5.6
Track p_T resolution	0.3	0.3	0.2	0.2	~ 0	0.3
Unfolding procedure	0.6	0.7	3.9	0.3	3.4	0.9
c_{Ref} variation	(−1.7, 1.9)	(−0.2, 0.2)	(−0.6, 0.7)	(−0.4, 0.4)	(−0.5, 0.5)	(−0.1, 0.2)
Bkgd fluctuations	(−1.6, 0)	(−2.4, 0)	(−4.7, 0)	(−1.6, 0)	(−5.0, 0)	(−2.7, 0)
Total uncertainty	(−2.5, 2.0)	(−7.6, 7.2)	(−7.2, 5.4)	(−2.5, 2.0)	(−10.4, 9.0)	(−6.3, 5.7)

Table 5.2: Main sources of systematic uncertainty and total uncertainty in $\Delta_{\text{recoil}}(p_{T,\text{jet}}^{\text{ch, corr}})$ and $\Delta_{\text{recoil}}(|\Delta\varphi|)$ in representative bins, for MB events.

5.6.1 Tracking efficiency

The tracking efficiency of inclusive charged tracks has a relative uncertainty of 3% [205], [206]. This uncertainty affects the energy scale of reconstructed jets. To estimate the

Projection	Relative systematic uncertainty (%)					
	$\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch, corr}})$ $ \Delta\varphi - \pi < 0.6$		$\Delta_{\text{recoil}}(\Delta\varphi)$ $p_{\text{T,jet}}^{\text{ch, corr}} \in (20, 40) \text{ GeV}/c$		$\Delta_{\text{recoil}}(\Delta\varphi)$ $p_{\text{T,jet}}^{\text{ch, corr}} \in (40, 60) \text{ GeV}/c$	
	10–20 GeV/c	60–80 GeV/c	$2\pi/3$	π	$2\pi/3$	π
Bin						
Tracking efficiency	0.6	7.6	4.2	1.9	7.1	5.2
Track p_{T} resolution	0.3	0.1	0.4	0.2	0.1	0.3
Unfolding procedure	0.9	0.7	5.7	1.6	2.6	1.2
c_{Ref} variation	(−4.5, 3.2)	(−0.3, 0.2)	(−1.5, 1.1)	(−0.6, 0.5)	(−1.3, 0.9)	(−0.3, 0.2)
Bkgd fluctuations	(−1.5, 0)	(−3.0, 0)	(−7.4, 0)	(−2.1, 0)	(−4.5, 0)	(−3.2, 0)
Total uncertainty	(−4.6, 3.7)	(−8.2, 7.6)	(−10.3, 7.2)	(−3.3, 2.5)	(−8.9, 7.6)	(−6.3, 5.4)

Table 5.3: Same as Tabel 5.2, for HM events.

corresponding uncertainty of the Δ_{recoil} distribution, a new R_{instr} was constructed, which accounted for 3% lower tracking efficiency, i.e. 3% of all detector-level tracks were excluded randomly from the jet reconstruction procedure. The ratios of the corrected Δ_{recoil} distributions obtained for the PA and the analysis with reduced tracking efficiency are shown in Fig. 5.23. The relative systematic uncertainty is represented by the dashed line, which is the median value of the 10 inputs in each bin. While generating a response matrix with 3% higher tracking efficiency is not feasible, this uncertainty is expected to be symmetric with respect to unity. It is worth stressing that the resulting uncertainty demonstrates a minor dependence on event activity.

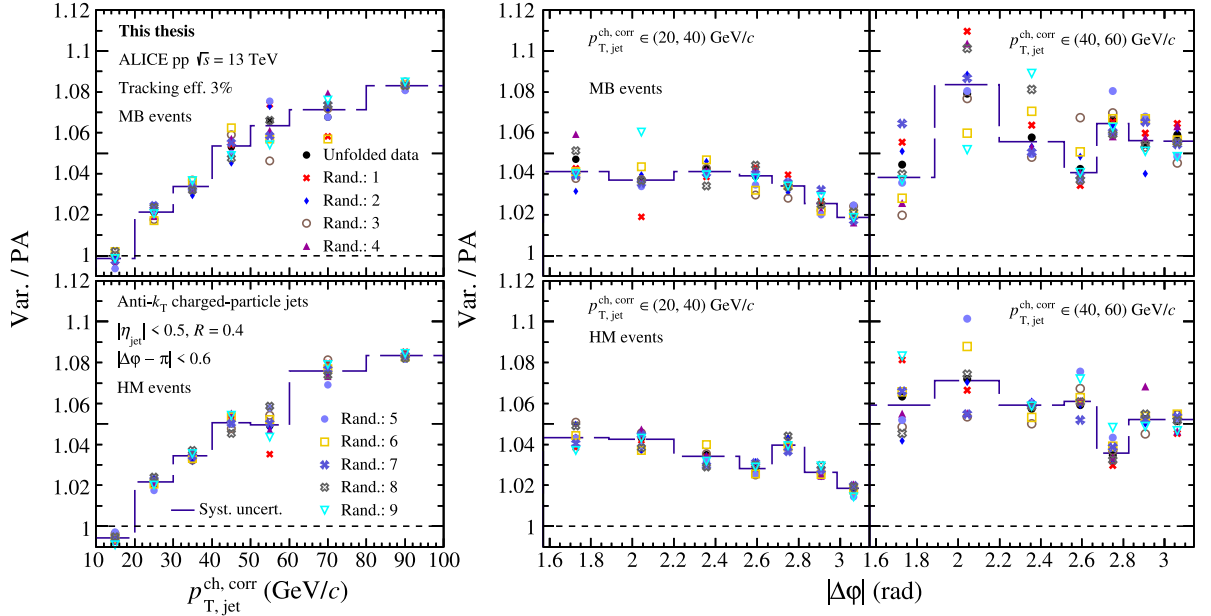


Figure 5.23: The ratio of the 10 unfolded Δ_{recoil} distributions obtained with 3% lower tracking efficiency and the PA. Top row: MB events. Bottom row: HM events. The blue dashed line represents the resulting relative systematic uncertainty.

5.6.2 Track- p_T resolution

The track momentum resolution is encoded in the tracking covariance matrix [189], which contains for each track the relative momentum uncertainty $\sigma_{p_{T, \text{trk}}}^{-1}/p_{T, \text{trk}}^{-1}$. The response matrix R_{instr} from the PA corrected the jet energy scale based on the p_T resolution of detector-level MC tracks. Therefore, the assessment of the systematic uncertainty due to the track- p_T resolution was based on the unfolding of the raw Δ_{recoil} spectra using different instances of R_{instr} that applied the same track reconstruction efficiency, but which differ in the track- p_T resolution. Track- p_T resolution corresponded to that of either reconstructed hybrid tracks or detector-level MC tracks. The corresponding response matrices were named “Hybrid matrix” and “MC matrix”. The resulting relative systematic uncertainty was defined as the ratio of the Δ_{recoil} spectra unfolded with the “Hybrid matrix” and the “MC matrix”. The response matrices were constructed using standalone PYTHIA 8 Monash simulations of pp collisions at $\sqrt{s} = 13$ TeV. The simulations included the following steps:

1. Only charged final-state particles were considered. The particle-level jets were reconstructed from accepted particles with $|\eta| < 0.9$. For the reconstruction of detector-level jets, particles were required to have $|\eta| < 0.9$ and $p_T > 0.15$ GeV/c. To emulate tracking efficiency that affects the reconstruction of detector-level jets, the generated particles were discarded randomly with the probability equal to track reconstruction efficiency for a given p_T in MB events, see Fig. 5.6;
2. The relative inverse momentum resolution $\sigma_{p_{T, \text{trk}}}^{-1}/p_{T, \text{trk}}^{-1}$ of hybrid and MC tracks was obtained as a function of $p_{T, \text{trk}}$, see Fig. 5.24. Both real data and MC simulations have rather poor statistics for tracks with $p_{T, \text{trk}} > 30$ GeV/c. Thereby, the mean value of $\sigma_{p_{T, \text{trk}}}^{-1}/p_{T, \text{trk}}^{-1}$ for each $p_{T, \text{trk}}$ was extracted and fitted with a smooth

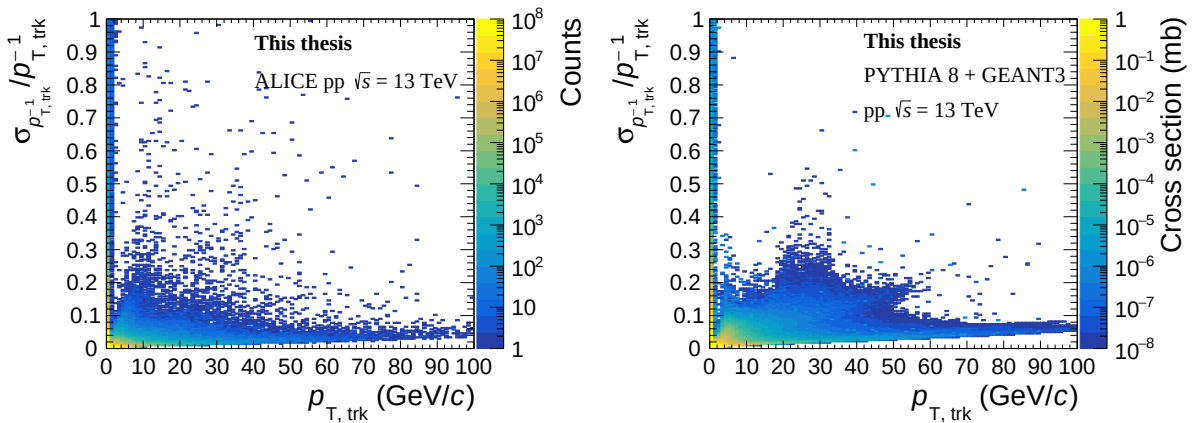


Figure 5.24: Relative inverse momentum resolution $\sigma_{p_{T, \text{trk}}}^{-1}/p_{T, \text{trk}}^{-1}$ for hybrid tracks (left) and MC detector-level tracks (right) as a function of $p_{T, \text{trk}}$.

function,

$$f(p_{T, \text{trk}}) = A + B \times \sqrt{p_{T, \text{trk}} + C}, \quad (5.9)$$

where A , B , and C represent fit parameters. Figure 5.25 shows the obtained dependence of $\langle \sigma_{p_{T, \text{trk}}}^{-1} / p_{T, \text{trk}} \rangle$ on $p_{T, \text{trk}}$. If a PYTHIA-generated particle had $p_T < 30$ GeV/ c then the $\sigma_{p_{T, \text{trk}}}^{-1} / p_{T, \text{trk}}$ value was sampled from the histograms shown in Fig. 5.24. Otherwise, the $\sigma_{p_{T, \text{trk}}}^{-1} / p_{T, \text{trk}}$ value was calculated based on the fit function from Eq. (5.9);

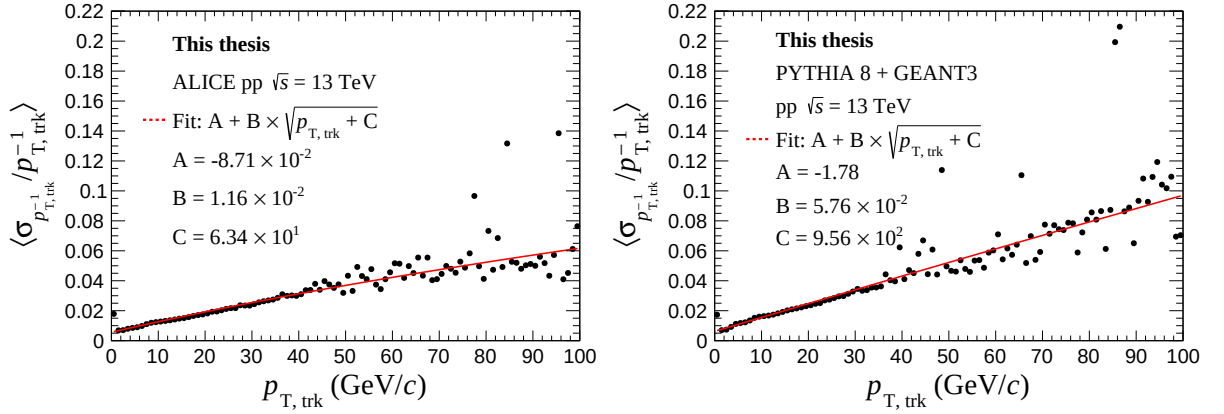


Figure 5.25: Relative mean inverse momentum resolution $\sigma_{p_{T, \text{trk}}}^{-1} / p_{T, \text{trk}}$ for hybrid tracks (left) and MC detector-level tracks (right) as a function of $p_{T, \text{trk}}$. The red line represents parameterization of the trend by Eq. (5.9).

3. Inverse p_T of the detector-level jet constituents was randomized utilizing a Gaussian function, which had its mean value equal to $1/p_T$ of PYTHIA-generated particle and width $\sigma_{p_{T, \text{trk}}}^{-1}$ for hybrid or MC tracks. As a result, two samples of the detector-level jets were obtained. Finally, two response matrices were constructed based on the geometric matching between the particle-level and detector-level jets.

Figure 5.26 shows the resulting systematic uncertainty on track- p_T resolution. Overall, this source has a negligible contribution to the total systematic uncertainty. Its contribution is considered to be symmetric.

5.6.3 Unfolding settings

The unfolding procedure has several parameters whose variations affect the corrected Δ_{recoil} distribution: the number of Bayesian iterations, choice of a prior spectrum, and binning of the input raw distribution. Each setting was varied independently of the others as follows:

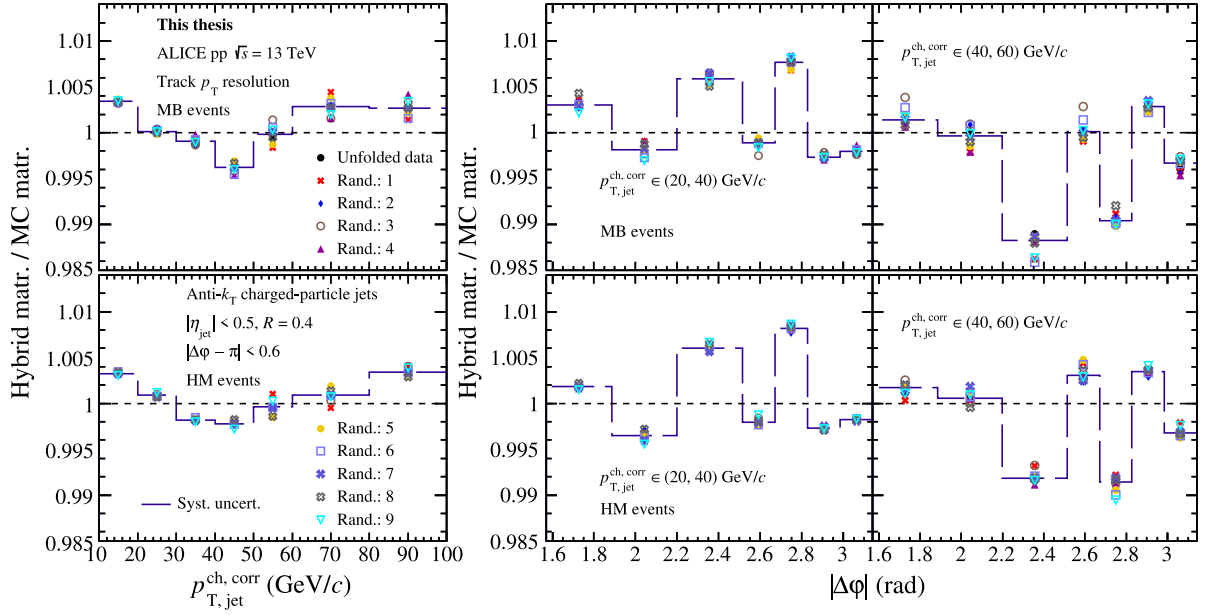


Figure 5.26: The ratio of the 10 Δ_{recoil} distributions unfolded with the “Hybrid matrix” and “MC matrix”. Top row: MB events. Bottom row: HM events. The blue dashed line represents the resulting relative systematic uncertainty.

- The number of Bayesian iterations was varied by ± 1 with respect to the chosen optimal value of 5. The associated uncertainty is relatively small since the unfolding solution converges rapidly to a stable result, see Sect. 5.5.5;
- Alternative prior spectra for MB events were obtained using the POWHEG simulations [207] of pp collisions at $\sqrt{s} = 13$ TeV. Three combinations of the renormalization (“Ren”) and factorization (“Fact”) parameters were considered, namely: Ren = 1 and Fact = 1; Ren = 1 and Fact = 0.5; Ren = 2 and Fact = 1. Instead of simulating HM POWHEG data, which is time-consuming, I used as alternative HM prior spectra the MB POWHEG Δ_{recoil} distributions scaled by the ratio HM/MB from PYTHIA simulations. The three generated prior spectra for MB and HM events were utilized to unfold the corresponding raw data;
- Binning of the input distribution was varied by shifting the boundaries of the inner bins by 1–2 GeV/ c and by changing the lower bound from 10 GeV/ c to 6 GeV/ c . The binning of the $|\Delta\phi|$ axis was kept unchanged since the smearing of $|\Delta\phi|$ due to instrumental effects is small.

The corresponding median distributions were calculated for each of the variations mentioned above. These distributions are shown in Fig. 5.27. The systematic uncertainty attributed to unfolding was defined as the maximum deviation from unity in each bin in Fig. 5.27. This source of systematic uncertainty was treated as symmetric.

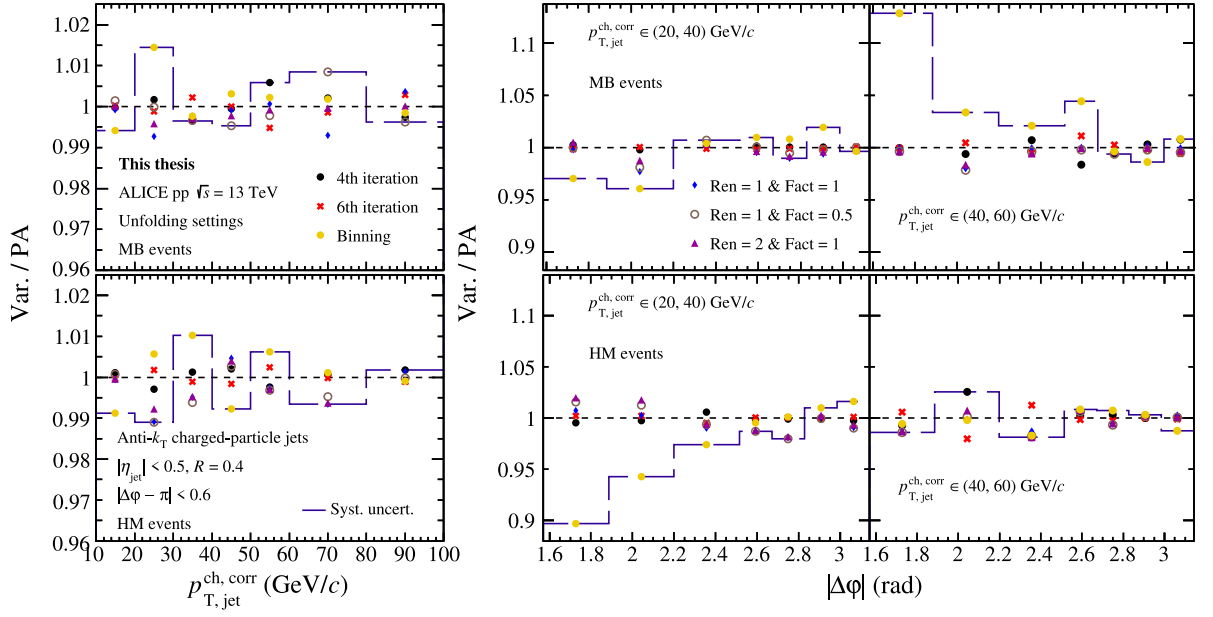


Figure 5.27: The ratio of the 10 unfolded Δ_{recoil} distributions obtained with varied unfolding parameters and the PA. Top row: MB events. Bottom row: HM events. The blue dashed line represents the resulting relative systematic uncertainty.

5.6.4 Variations of the c_{Ref} factor

In the PA, the c_{Ref} values (Fig. 5.7) were 0.95 for MB events and 0.94 for HM events. Note that these values were averaged over $|\Delta\varphi| \in (\pi - 0.6, \pi)$ (recoil acceptance). As a systematic variation, the value of c_{Ref} was varied in the range (0.9, 1), corresponding to the span of the c_{Ref} value considered as a function of $|\Delta\varphi|$ angle. Figure 5.28 shows the ratio of the unfolded Δ_{recoil} distributions for the c_{Ref} values of 0.9 and 1 to the Δ_{recoil} distribution from the PA. Each distribution represents the median value from the analyses with the randomized Δ_{recoil} spectra. The systematic uncertainty was defined as the value of the largest deviation from unity in each bin in Fig. 5.28. This source was considered asymmetric. Thus, the deviation was taken on both sides of unity. The variation of the c_{Ref} value leads to a relatively small systematic uncertainty. Its magnitude grows at low p_T , where two per-trigger normalized distributions of recoil jets are very close to each other. Furthermore, the systematic uncertainty is slightly larger for HM data. It is explained by a relatively larger population of underlying events at ~ 0 GeV/c in HM data.

5.6.5 Background fluctuations

As discussed in Sect. 5.5.4, the raw data were not corrected for the effect of local background fluctuations with respect to the median UE density because of the results of the closure test for the combined matrix R_{comb} . This systematic uncertainty reflects the dif-

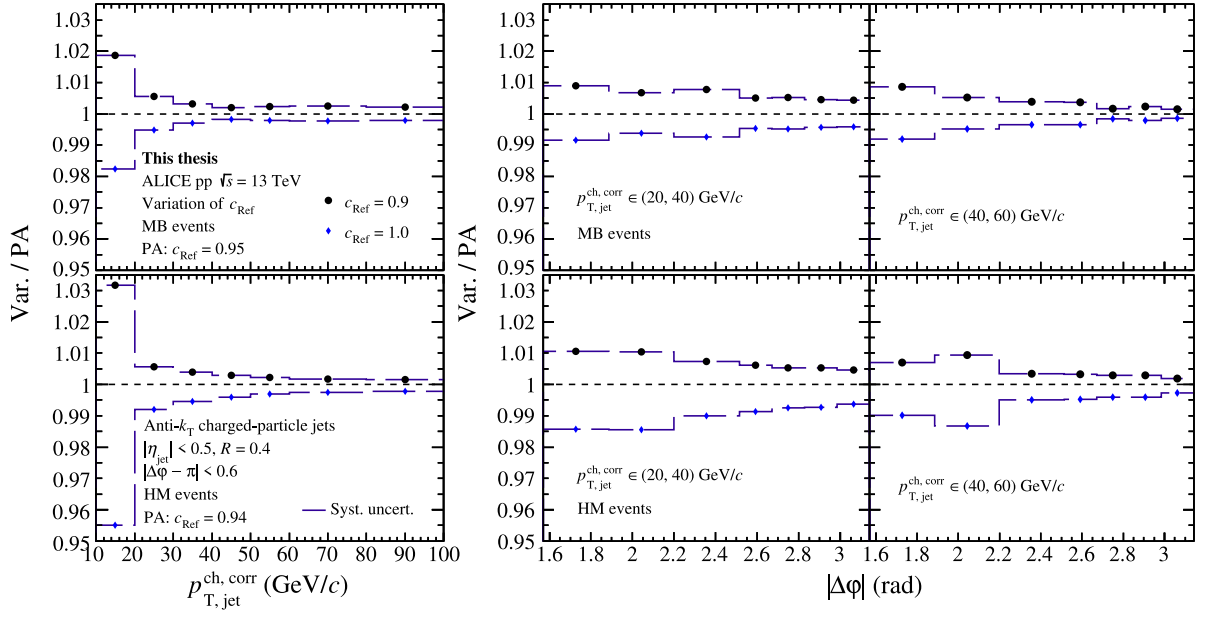


Figure 5.28: Ratio of the Δ_{recoil} distributions obtained for analyses with the varied c_{Ref} values and PA. Top row: MB events. Bottom row: HM events. The blue dashed line represents the resulting systematic uncertainty.

ference between the Δ_{recoil} distributions unfolded with R_{instr} and R_{comb} . To construct the combined matrix for raw data unfolding, one needs to estimate the δp_T smearing for the measured pp events at $\sqrt{s} = 13$ TeV containing a $\text{TT}\{20, 30\}$. The δp_T smearing was measured utilizing two methods: random cone and track embedding. The first method was employed and discussed in Sect. 5.5.4. In the track embedding approach, a track with uniformly generated p_T in the range (0, 20) GeV/c was placed in each event perpendicularly to the $\text{TT}\{20, 30\}$ in azimuth at the same η as the TT,

$$\varphi_{\text{track}}^{\text{emb}} = \varphi_{\text{TT}} + \pi/2, \quad \eta_{\text{track}}^{\text{emb}} = \eta_{\text{TT}}. \quad (5.10)$$

In the next step, the jet reconstruction procedure was launched, and δp_T was calculated utilizing the jets containing the embedded track,

$$\delta p_T = p_{T, \text{jet}}^{\text{ch}} - \rho \times A_{\text{jet}} - p_T^{\text{emb track}}. \quad (5.11)$$

Figure 5.29 shows the δp_T distributions obtained with the random cone and track embedding methods for MB and HM events. Both approaches provide close distributions with a prominent peak at 0 GeV/c and with a steep fall on both sides. Due to the requirement of the $\text{TT}\{20, 30\}$ presence in each event, the obtained distributions have limited statistical precision at high p_T , which may affect the data unfolding. Therefore, the δp_T

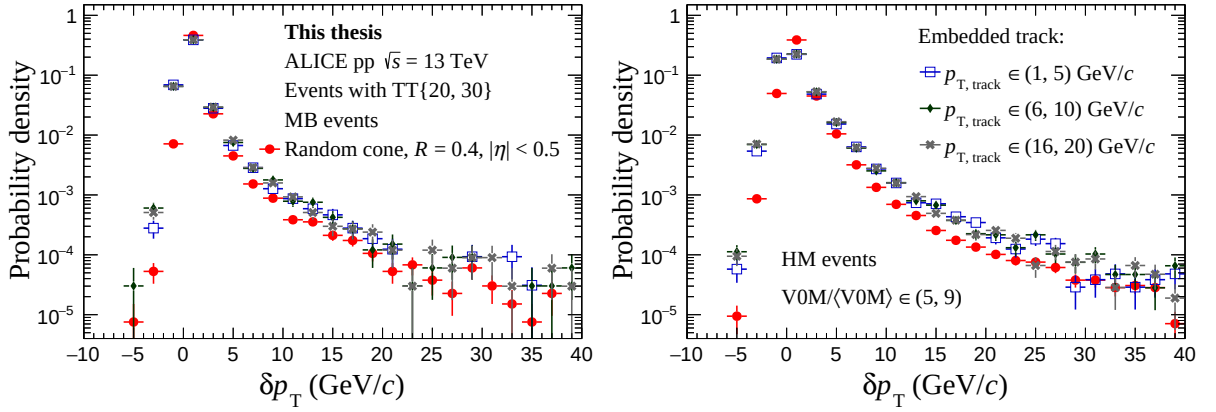


Figure 5.29: δp_T smearing due to local background fluctuations measured using the random cone and track embedding methods in MB and HM pp collisions at $\sqrt{s} = 13$ TeV.

distributions were fitted in the range (8, 40) GeV/c with a smooth function,

$$f(\delta p_T) = A \times \text{Exp}(B + C \times \delta p_T) + D \times (\delta p_T)^E, \quad (5.12)$$

where A, B, C, D , and E are parameters of the fit. The smoothed δp_T distributions are shown in Fig. 5.30. These distributions were used to construct the R_{bkgd} matrices and

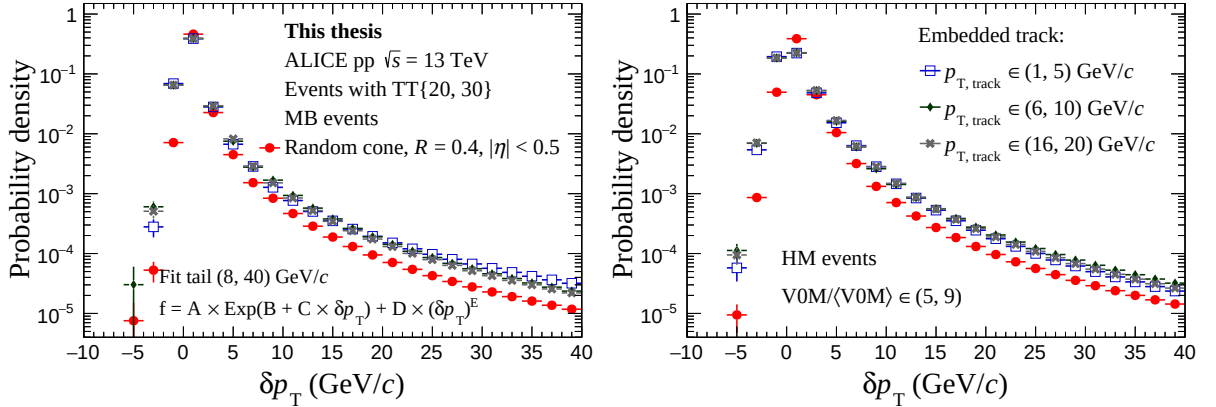


Figure 5.30: Same as Fig. 5.29 after the fit with Eq. 5.12.

then the R_{comb} matrices.

Figure 5.31 shows the ratio of the Δ_{recoil} distributions unfolded with the combined and instrumental matrices. The resulting systematic uncertainty was obtained as the maximum deviation from unity. This source was considered to provide asymmetric systematic uncertainty.

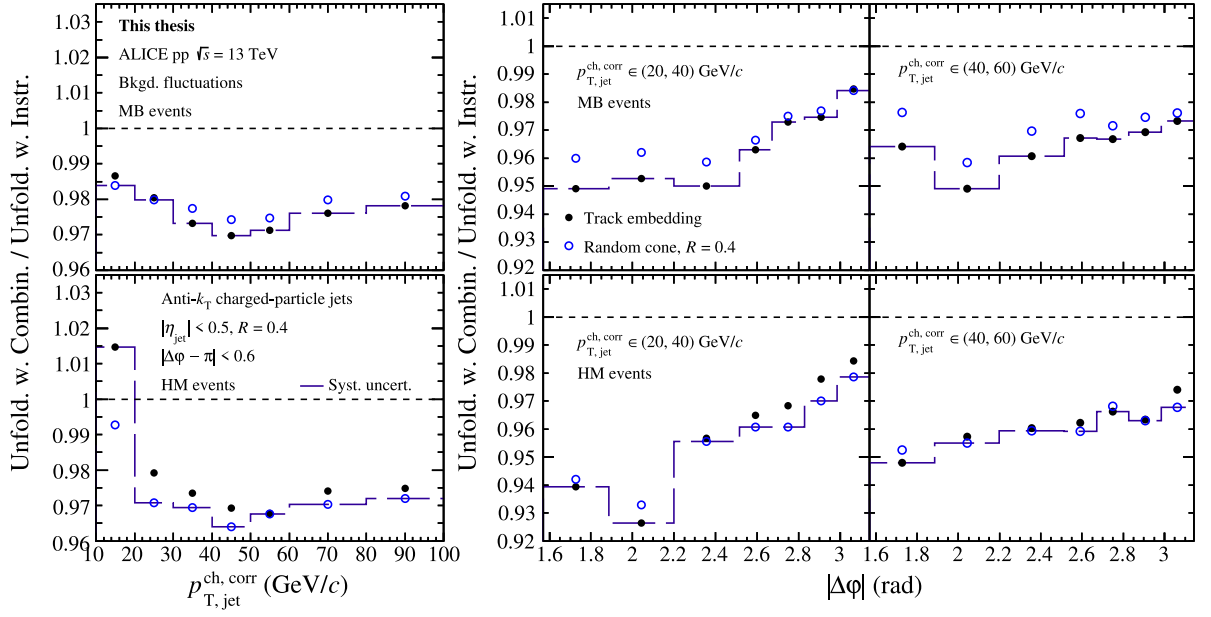


Figure 5.31: Systematic uncertainty due to the correction of the Δ_{recoil} distribution for local background fluctuations in MB (top) and HM (bottom) pp collisions at $\sqrt{s} = 13$ TeV. The δp_T smearing was obtained utilizing the random cone and track embedding approaches. The blue dashed line shows the resulting magnitude of the uncertainty.

5.6.6 Total systematic uncertainty

Figure 5.32 shows the total systematic uncertainties for the MB and HM analyses, which were obtained by summing up uncertainties from the abovementioned sources in quadrature. Those sources of systematic uncertainty that were considered symmetric are shown symmetrically on both sides with respect to the dashed line.

The systematic uncertainties from all sources, except the unfolding, are correlated in the HM and MB analyses. Thus, their contribution cancels each other to a large extent when taking the HM/MB ratio of the Δ_{recoil} distributions. To assess the systematic uncertainty of the HM/MB distribution, the same chain of steps as for the Δ_{recoil} spectra was followed. The obtained results are shown in Fig. 5.33.

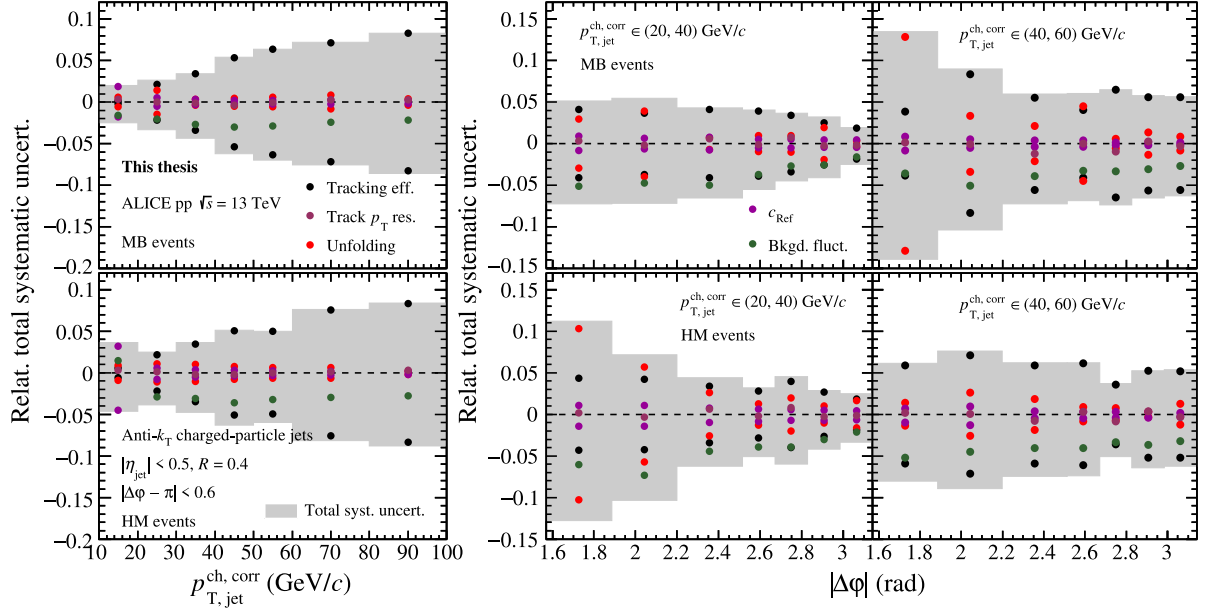


Figure 5.32: The total systematic uncertainty for the Δ_{recoil} distributions measured in MB and HM pp collisions at $\sqrt{s} = 13$ TeV. Top row: MB data. Bottom row: HM data.

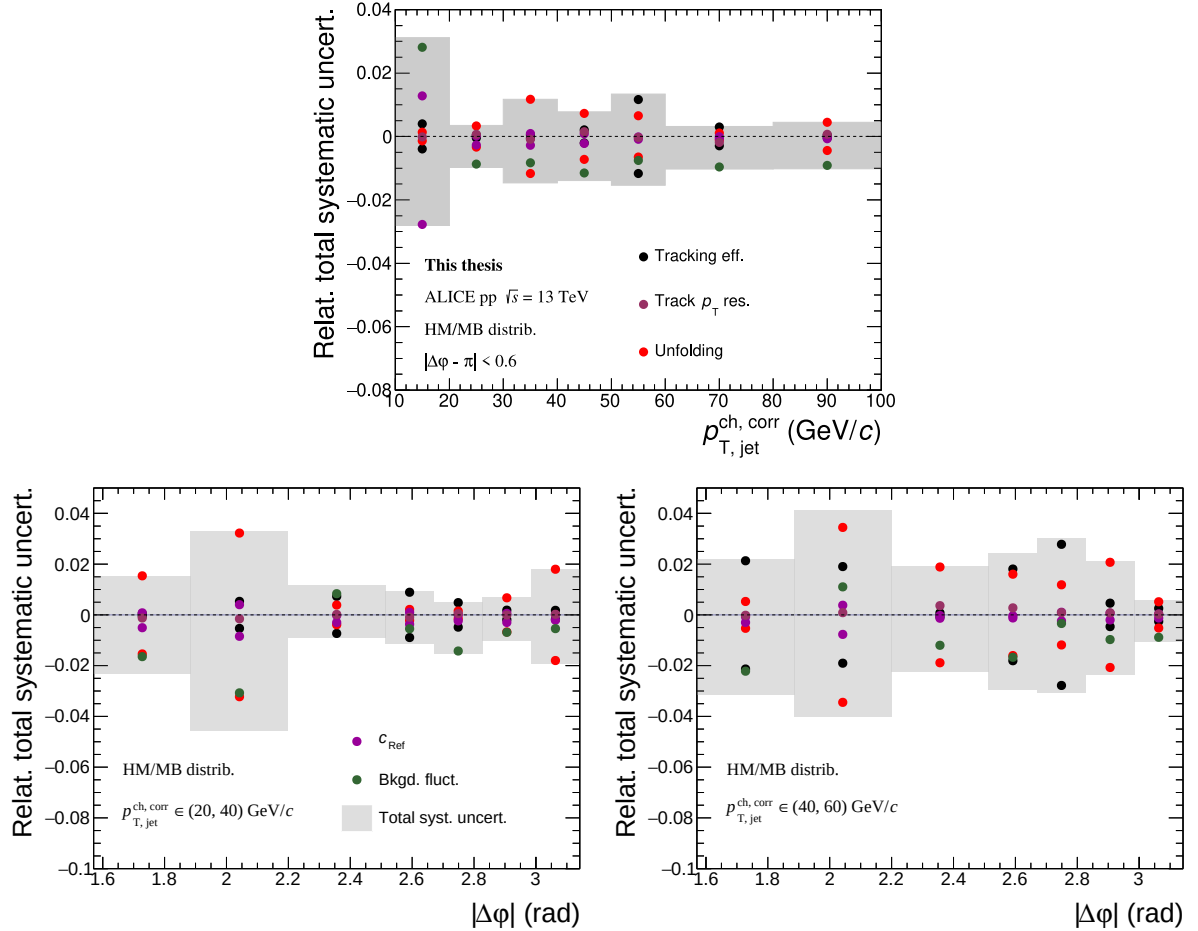


Figure 5.33: The total systematic uncertainty for the HM/MB ratios measured in pp collisions at $\sqrt{s} = 13$ TeV.

Chapter 6

Analysis results

6.1 Fully corrected data

The fully corrected $\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch,corr}})$ and $\Delta_{\text{recoil}}(|\Delta\varphi|)$ distributions measured in MB and HM pp collisions at $\sqrt{s} = 13$ TeV are shown in Fig. 6.1. In addition, simulations based on PYTHIA 8 Monash and pQCD calculations at LO with Sudakov broadening [38], [208],

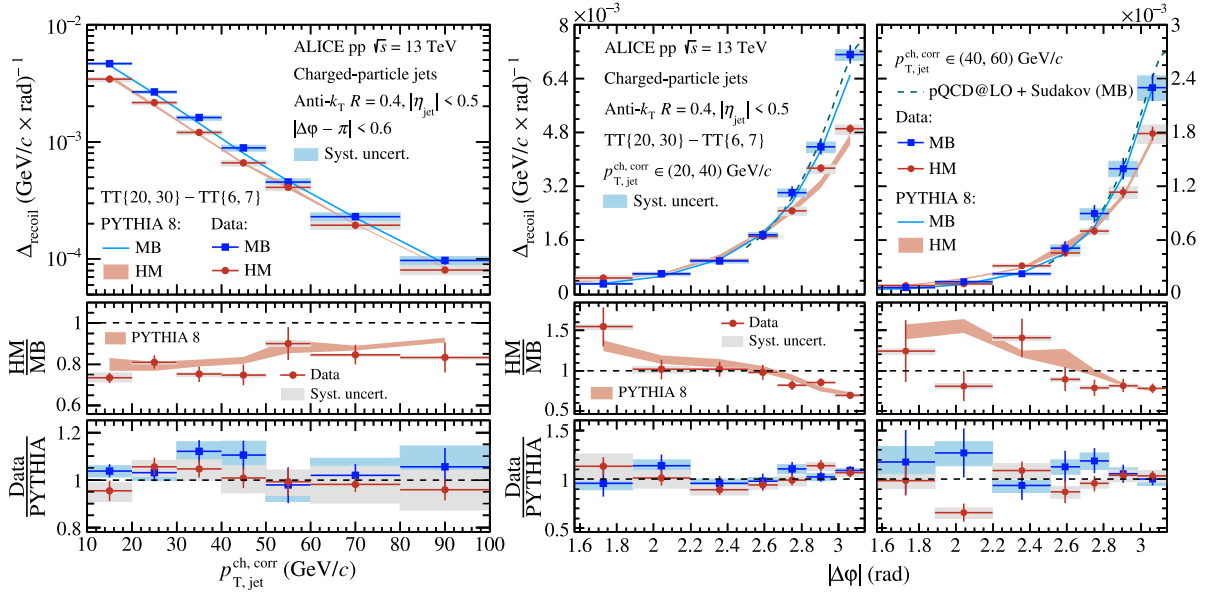


Figure 6.1: Fully corrected Δ_{recoil} distributions measured in MB and HM pp collisions at $\sqrt{s} = 13$ TeV. Left panel: $\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch,corr}})$ in $|\Delta\varphi - \pi| < 0.6$. Middle and right panels: $\Delta_{\text{recoil}}(|\Delta\varphi|)$ for $p_{\text{T,jet}}^{\text{ch,corr}} \in (20, 40)$ GeV/c and $p_{\text{T,jet}}^{\text{ch,corr}} \in (40, 60)$ GeV/c. The particle-level PYTHIA 8 Monash simulations are shown by the solid lines, where the line width represents the statistical errors. The PYTHIA-generated HM events are selected in the range $V0M/\langle V0M \rangle \in (4.4, 9)$. The pQCD calculations at LO with Sudakov broadening are available only for the shape comparison with MB $\Delta_{\text{recoil}}(|\Delta\varphi|)$. Middle rows: ratio of Δ_{recoil} distributions for HM and MB data. Bottom rows: ratio Data/PYTHIA separately for MB and HM event selections. Accessible in Ref. [39].

[209] are also presented. The latter data are only provided for $\Delta_{\text{recoil}}(|\Delta\varphi|)$ in MB events. Since the pQCD calculations are LO, their normalisation has large uncertainties. Thus, the $\Delta_{\text{recoil}}(|\Delta\varphi|)$ distributions from these calculations were scaled to the integrated yield of the data in the same $p_{\text{T,jet}}^{\text{ch,corr}}$ range to provide solely shape comparison. PYTHIA 8 Monash and the pQCD calculations (shape only) are consistent with the distributions measured in MB events within experimental uncertainties. Figure 6.2 shows a comparison of MB data with NLO POWHEG calculations [207]. The systematic uncertainty of POWHEG calculations was estimated by varying the regularization and factorization scale parameters. One can see that the POWHEG calculations likewise well describe the data within uncertainties.

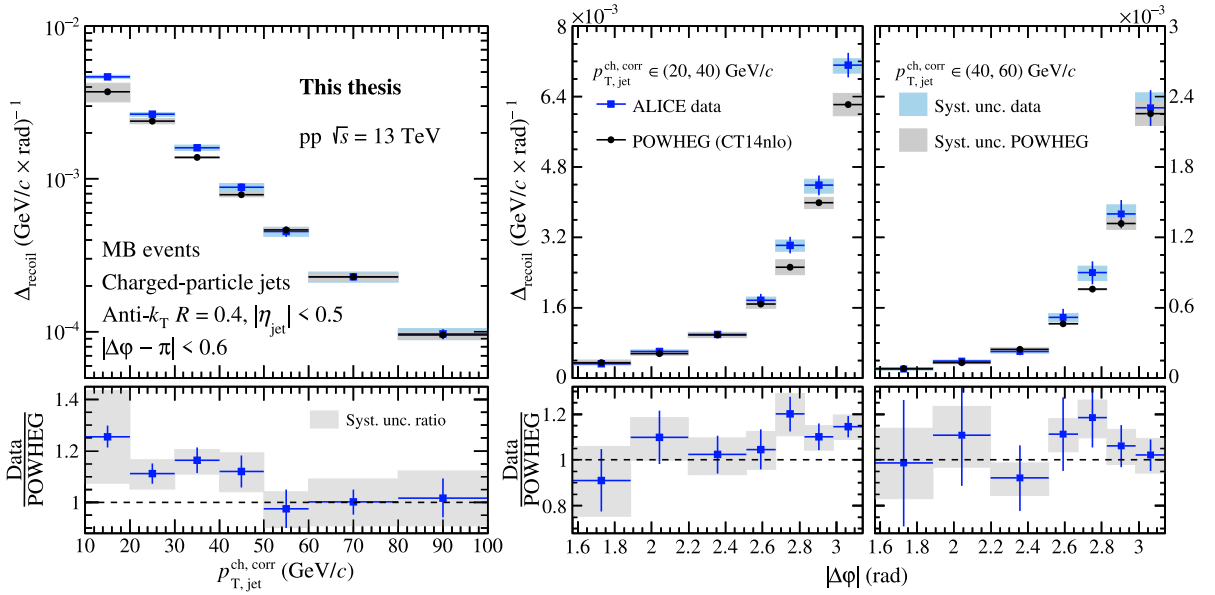


Figure 6.2: Fully corrected Δ_{recoil} distributions measured in MB pp collisions at $\sqrt{s} = 13$ TeV. Left panel: $\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch,corr}})$ in $|\Delta\varphi - \pi| < 0.6$. Middle and right panels: $\Delta_{\text{recoil}}(|\Delta\varphi|)$ for $p_{\text{T,jet}}^{\text{ch,corr}} \in (20, 40)$ GeV/c and $p_{\text{T,jet}}^{\text{ch,corr}} \in (40, 60)$ GeV/c. The data are compared with NLO POWHEG calculations. Bottom rows: ratio Data/POWHEG.

Comparison of the corrected $\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch,corr}})$ distributions demonstrates the yield suppression in HM events, which weakly depends on jet p_{T} . The mean value of the yield ratio HM/MB in the left-middle panel of Fig. 6.1 is about 0.8. The $\Delta_{\text{recoil}}(|\Delta\varphi|)$ acoplanarity distributions exhibit substantial yield suppression in HM events compared to MB events, which occurs predominantly in the back-to-back configuration at $|\Delta\varphi| \sim \pi$. The total yield of recoil jets is suppressed in HM events while the acoplanarity distributions are broadened. Such broadening can arise from jet quenching [38], [116], [182], i.e. the medium-induced parton scattering occurs in HM pp collisions. However, the PYTHIA 8 Monash simulation reproduces yield suppression of recoil jets and azimuthal broadening in HM events, accurately describing the measured data. Since the PYTHIA 8 event generator does not incorporate phenomenological models for jet quenching, this excludes jet

quenching as the predominant effect generating the broadening seen in the data.

6.2 Origin of the high-multiplicity induced hadron–jet acoplanarity in PYTHIA

To investigate the origin of the acoplanarity suppression and broadening for HM events, the PYTHIA 8 Monash simulations were further explored in detail. Figure 6.3 shows the $\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch,corr}})$ and $\Delta_{\text{recoil}}(|\Delta\varphi|)$ distributions for charged-particle jets with $R = 0.4$ reconstructed in a wide pseudorapidity range $|\eta_{\text{jet}}| < 5.6$, which also covers the acceptances of the V0 forward detectors. If jet yield suppression in Fig. 6.1 were caused by jet quenching, then the HM Δ_{recoil} distributions in Fig. 6.3 would also be suppressed. However, one can observe no signs of such suppression. Nevertheless, the HM $\Delta_{\text{recoil}}(|\Delta\varphi|)$ distributions remain broader than the corresponding MB distributions.

Figure 6.4 shows the pseudorapidity (η_{jet}) distribution of charged-particle jets with $R = 0.4$ that recoil from the $\text{TT}\{20, 30\}$ as a function of various lower thresholds in $p_{\text{T,jet}}^{\text{ch}}$ for MB and HM events. The recoil region in azimuth was defined by the opposite hemisphere with respect to the $\text{TT}\{20, 30\}$, $|\varphi_{\text{TT}} - \varphi_{\text{jet}}| > \pi/2$. Since the trigger-normalized recoil jet yield varies with $p_{\text{T,jet}}^{\text{ch}}$, the distributions were normalized to the unit integral

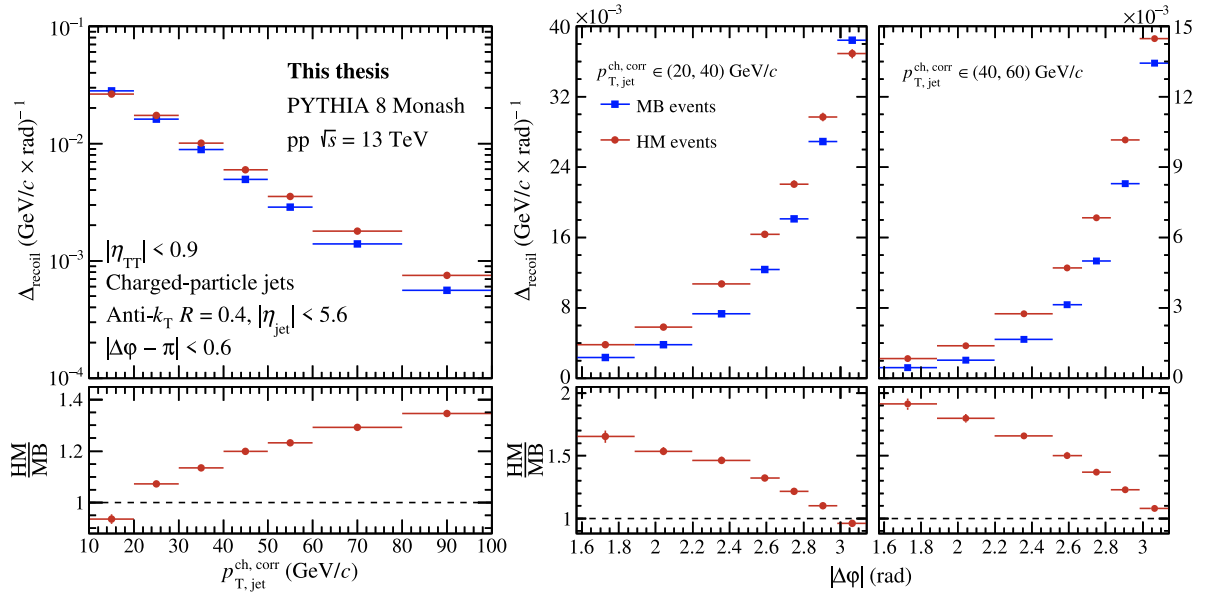


Figure 6.3: Δ_{recoil} distributions measured in PYTHIA 8 Monash simulation of MB and HM pp collisions at $\sqrt{s} = 13$ TeV. Jets are reconstructed in a wide pseudorapidity range $|\eta_{\text{jet}}| < 5.6$. Trigger track has $|\eta_{\text{TT}}| < 0.9$. Left panel: $\Delta_{\text{recoil}}(p_{\text{T,jet}}^{\text{ch,corr}})$ in $|\Delta\varphi - \pi| < 0.6$. Middle and right panels: $\Delta_{\text{recoil}}(|\Delta\varphi|)$ for $p_{\text{T,jet}}^{\text{ch,corr}} \in (20, 40)$ GeV/c and $p_{\text{T,jet}}^{\text{ch,corr}} \in (40, 60)$ GeV/c. HM events are selected in the range $\text{V0M}/\langle \text{V0M} \rangle \in (4.4, 9)$. Bottom row: ratio of Δ_{recoil} distributions for HM and MB data.

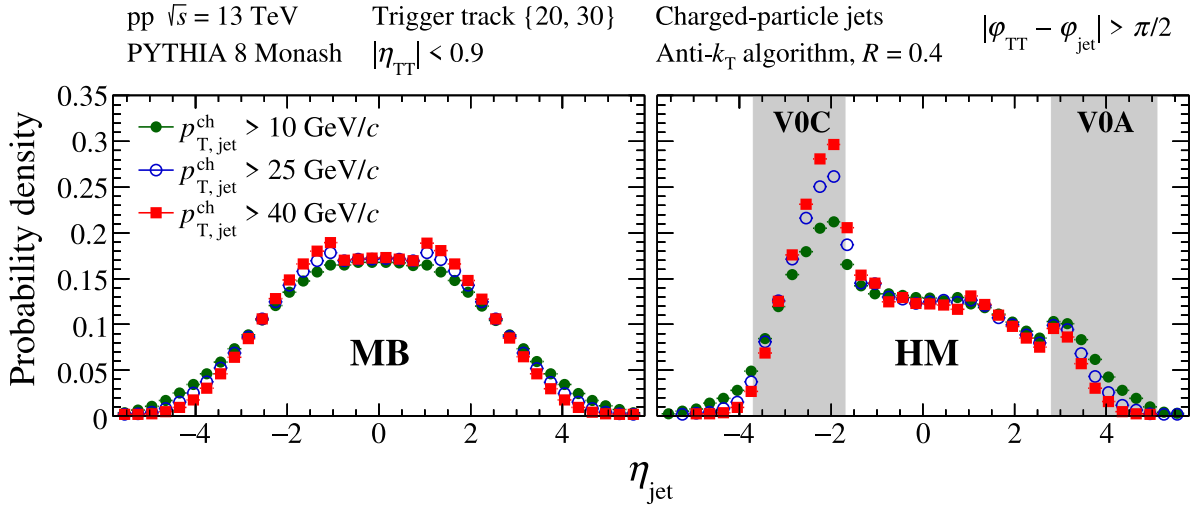


Figure 6.4: Pseudorapidity distribution of jets recoiling from $\text{TT}\{20, 30\}$ for various $p_{\text{T}, \text{jet}}^{\text{ch}}$ thresholds in pp collisions $\sqrt{s} = 13$ TeV simulated by PYTHIA 8 Monash. Left: MB data. Right: HM events, where $\text{V0M}/\langle\text{V0M}\rangle \in (4.4, 9)$. The acceptances of the V0A and V0C detectors are also shown. Accessible in Ref. [39].

to enable a comparison of their shapes. One can note that the η_{jet} distributions for MB events are symmetric with respect to $\eta_{\text{jet}} = 0$, whereas the analogous distributions for HM events are highly asymmetric. Remarkably, the relative yield of recoil jets is enhanced more in the V0C acceptance since the V0C detector is located closer to the ALICE interaction point than the V0A detector, see Fig. 4.1. Moreover, the enhancement grows with $p_{\text{T}, \text{jet}}^{\text{ch}}$.

Figure 6.5 shows the similar η_{jet} distributions for recoil jets with $p_{\text{T}, \text{jet}}^{\text{ch}} > 25$ GeV/ c , but as a function of different intervals of event multiplicity $\text{V0M}/\langle\text{V0M}\rangle$. A prominent enhancement is observed in the per-trigger normalized yield of recoil jets within the V0C acceptance, which grows when event multiplicity increases. Since HM events are selected based on the value of $\text{V0M}/\langle\text{V0M}\rangle$, one may conclude that the HM selection induces a bias which enhances the yield of high- p_{T} recoil jets in the V0C acceptance.

Supplementary insights into the bias induced by the HM event selection are provided by the distribution, which gives the probability to observe a specific number of recoil jets with $p_{\text{T}, \text{jet}}^{\text{ch}} > 15$ GeV/ c and $R = 0.4$ reconstructed in the ALICE fiducial volume ($|\eta_{\text{jet}}| < 0.5$) for MB and HM events. Figure 6.6 shows these distributions for PYTHIA 8 Monash simulations and ALICE data. The probability of measuring two or more recoil jets in both MB and HM events is suppressed with respect to the LO di-jet production channel. Both data and simulations suggest that the probability of reconstructing a single high- p_{T} recoil jet in HM events is suppressed relative to MB events by 2–3%. At the same time, the probability of observing multiple jets is substantially enhanced.

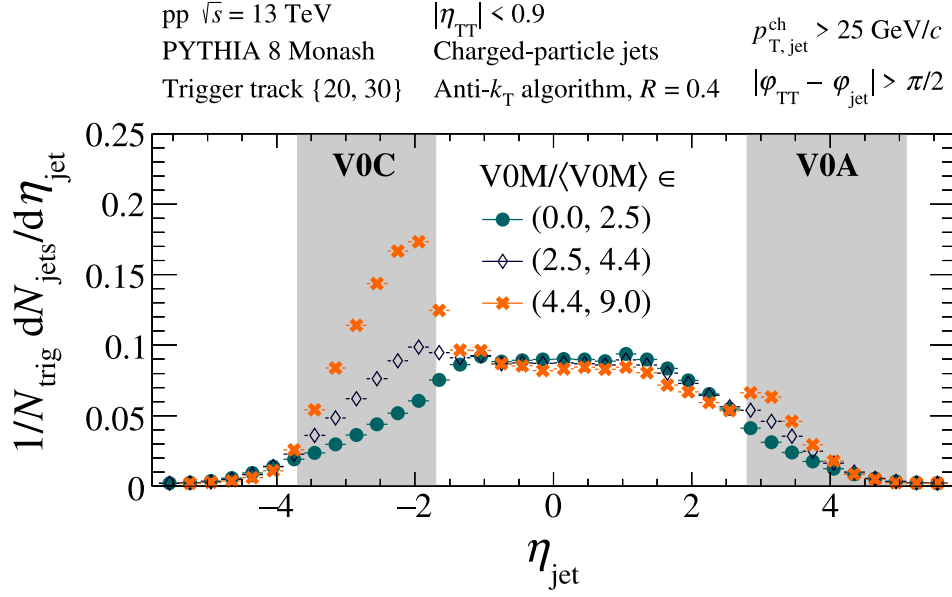


Figure 6.5: Same as Fig. 6.4, but for recoil jets with $p_{\text{T,jet}}^{\text{ch}} > 25 \text{ GeV}/c$ for various intervals in $V0M/\langle V0M \rangle$. Accessible in Ref. [39].

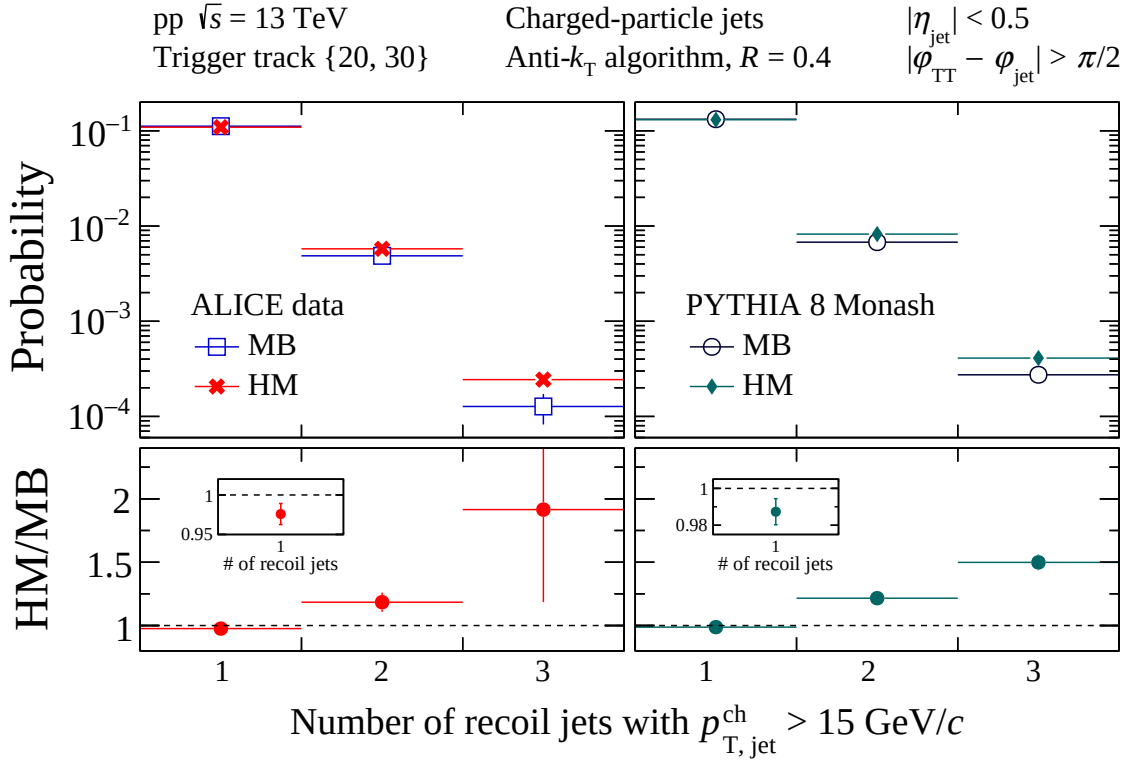


Figure 6.6: Probability distribution of the number of recoil jets with $p_{\text{T,jet}}^{\text{ch}} > 15 \text{ GeV}/c$ reconstructed in the ALICE central barrel ($|\eta_{\text{jet}}| < 0.5$) for MB and HM events with TT{20,30}. Left panels: ALICE data. Right panels: particle-level PYTHIA 8 Monash simulations, where $V0M/\langle V0M \rangle \in (4.4, 9)$. The bottom panels show the ratio HM/MB. The insert in the lower panels has a magnified vertical scale to show the ratio of probabilities to observe a single recoil jet. Accessible in Ref. [39].

6.3 Discussion

The effect of the HM bias discussed in Sect. (6.2) on the hadron–jet acoplanarity distributions shown in Fig. 6.1 can be interpreted as follows. The imposed HM requirement preferentially selects those events, where a high- p_T recoil jet is located in the V0C acceptance. Therefore, for dijet production at LO, in which the TT and recoil jet are back-to-back in azimuth, this leads to the suppression of both the trigger-normalized yield of recoil jets at $|\Delta\varphi| \sim \pi$ measured in the central barrel ($|\eta_{\text{jet}}| < 0.9 - R$) and the relative abundance of single recoil jets in the acceptance. The enhanced relative abundance of multiple jets arising from higher-order QCD processes in HM events broadens the TT–jet acoplanarity distribution. These arguments are supplemented by the same analysis in PYTHIA 8 Monash events selected in a wide pseudorapidity range, $|\eta_{\text{jet}}| < 5.6$, see Fig. 6.3. In that case, no back-to-back yield suppression is observed.

The yield suppression and acoplanarity broadening seen in Fig. 6.1 arise from the effect of different phase space being used for event-activity characterization (forward rapidity V0A and V0C) and recoil jet measurement (central barrel), combined with the HM selection bias towards events with a high- p_T jet in the V0C acceptance. Therefore, it cannot be attributed uniquely to jet quenching in HM pp events.

This conclusion about interpretability in terms of jet quenching of the effects observed in Fig. 6.1 is more general than this analysis. The bias of the HM event selection identified in this analysis is generic. It must also be considered for the interpretation in terms of QGP formation of other phenomena observed in small collision systems.

Chapter 7

Conclusions

Collisions of small systems have been considered for a long time as control experiments for quark–gluon plasma discovery in ultra-relativistic heavy-ion collisions. However, pp and p–Pb collisions exhibit collective flow and enhanced production of strange hadrons, which are typically attributed to quark–gluon plasma formation in heavy-ion collisions. At the same time, there is no persuasive evidence of jet quenching in small collision systems within experimental uncertainties. A standard observable for jet quenching measurements in heavy-ion collisions is the nuclear modification factor, which compares the yield of high- p_T hadrons or jets measured in the nucleus–nucleus collisions to that expected from binary-collision scaled yield measured in pp collisions. The nuclear modification factor cannot, however, be measured in the case of a search for jet quenching carried out in pp collisions because Glauber modelling is undefined for pp collisions.

This thesis is devoted to a search for jet quenching effects in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV collected by the ALICE experiment in 2016–2018. Since inclusive observables cannot be used in this measurement, a novel approach based on the semi-inclusive acoplanarity distribution of charged-particle jets recoiling from a high- p_T trigger track was developed and studied. Event activity was characterized by charged-particle multiplicity measured in forward rapidity by the ALICE V0 detector. This detector has a different phase-space coverage than the ALICE central barrel, where the trigger track is measured, which mitigates autocorrelation bias. Substantial suppression and broadening of the acoplanarity distributions measured in high-multiplicity events were observed relative to the corresponding distributions for minimum-bias events. These features are expected to arise from jet quenching effects. However, simulations based on the PYTHIA 8 event generator with the Monash tune exhibit the same features. Therefore, the observed acoplanarity suppression and broadening cannot be genuinely attributed to jet quenching since the PYTHIA event generator does not incorporate models for jet quenching calculations.

An elaborate analysis of the data corrected for instrumental effects and simulations revealed that the ALICE high-multiplicity trigger preferentially selects events with a high- p_T recoil jet in the acceptance of the V0 forward scintillator detector. Consequently, it biases the acoplanarity distribution measured in the ALICE central barrel region ($|\eta| < 0.9$). Therefore, the proposed observable is not uniquely sensitive to hypothetical jet quenching signals in small collision systems. The bias on high-multiplicity selection due to high- p_T jets, however, is generic, and it should likewise be considered in the interpretation of other measurements, which probe the formation of quark–gluon plasma in small collision systems.

Chapter 8

ALICE ITS upgrade

During my PhD, I was involved in the ITS2 and subsequent ITS3 projects. In this section, I provide a general overview of the ITS2 project and describe my contribution.

Before Run 3 of the LHC, the ALICE detector underwent a significant upgrade to cope with the increased luminosity of Pb–Pb collisions [210]. A crucial part of the upgrade was the installation of a new silicon vertex tracker called the Inner Tracking System (ITS2) [211]. The ITS2 can measure Pb–Pb collisions at rates up to 100 kHz and pp collisions at rates up to 400 kHz. Figure 8.1 shows a schematic view of the upgraded ITS detector. It consists of 7 coaxial layers¹. The innermost layer was installed closer to the interaction point at a radial distance of 22 mm compared to 39 mm for the old ITS. The first three layers form the Inner Barrel (IB). The Outer Barrel (OB) incorporates the middle and outer layers. The layers are equipped with the ALPIDE Monolithic Active Pixel Sensors (MAPS), which have a thickness of 50 μm in the IB and 100 μm in the OB [211]. The

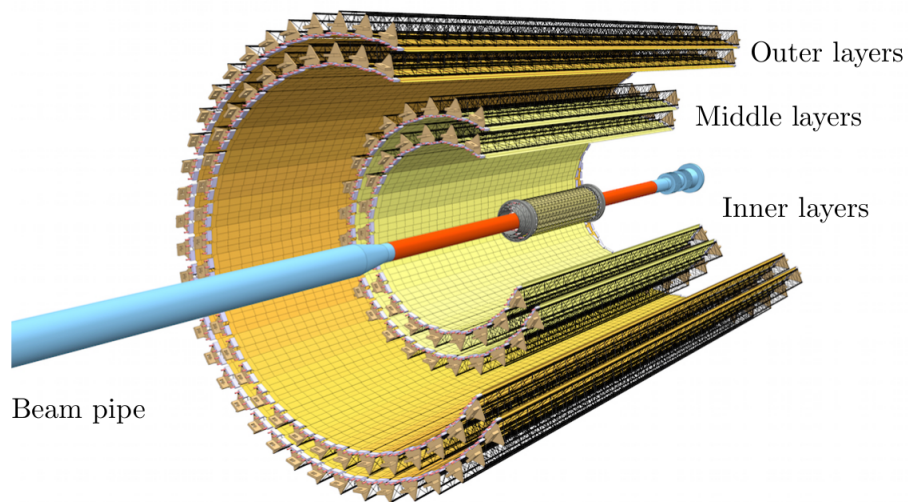


Figure 8.1: Layout of the new ITS detector for Run3. Taken from Ref. [211].

¹Previous ITS had 6 layers of silicon detectors, see Sect. 4.1.

new ITS has a material budget of $0.3\% X_0$ per layer and $1\% X_0$ per layer² for the IB and OB layers, respectively. In the azimuthal direction, each layer is segmented into mechanically independent structures called *staves*. The stave includes all structural and functional components, making it the smallest operable part of the detector. The new ITS covers a pseudorapidity range of $|\eta| < 1.22$. Overall, the detector surface amounts to 10.3 m^2 incorporating about 12.5×10^9 pixels with binary readout.

The upgraded ITS improves the pointing resolution by about a factor of 3 in the transverse plane ($r\phi$) and by about a factor of 5 along the beam axis (z) compared to the ITS detector in Run 1 and Run 2 (Fig. 8.2a). It also improves the tracking efficiency (see Fig. 8.2b) and p_T resolution of tracks with low p_T compared to the previous ITS [211]. The upgraded ITS allows for precise measurements of various observables, including low- p_T heavy-flavor hadrons, anti-nuclei, and hypernuclei, see Ref. [211] for further details.

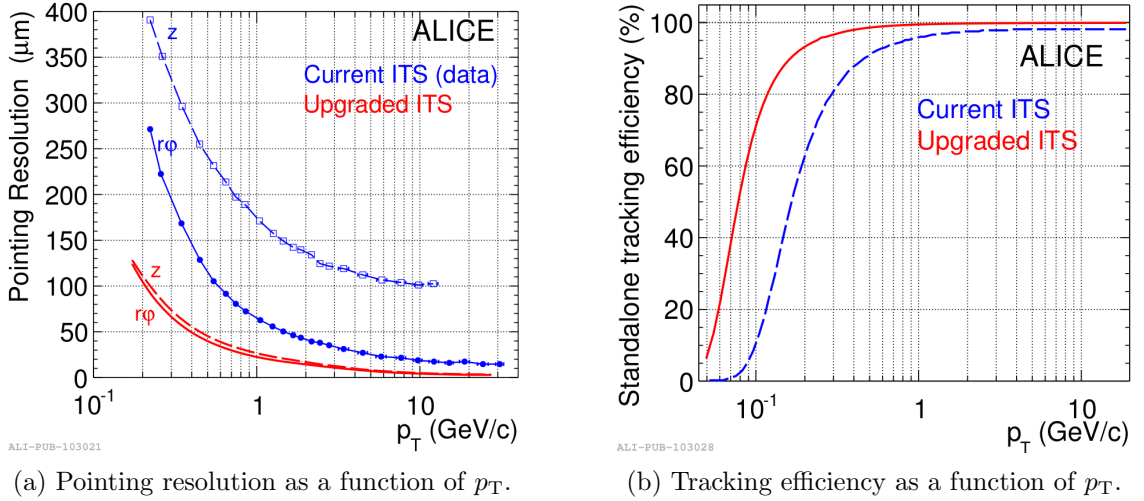


Figure 8.2: The stand-alone performance of the old ITS (blue data) compared to the simulated performance of the ITS2 (red data). The blue distributions correspond to Pb–Pb 2011 data. See text for further details. Taken from Ref. [211].

8.1 ALPIDE sensor

The design objectives, together with the operational environment of the new ITS, imposed requirements on the pixel chip, as shown in Table 8.1. After thorough research, the ALICE Collaboration concluded that the ALPIDE sensors fully satisfied the Technical Design Report (TDR) requirements. Figure 8.3 shows a schematic cross section of the ALPIDE sensor. The sensor production is based on TowerJazz 180 nm CMOS Imaging Process [212]. The ALPIDE sensor integrates detection volume and readout circuitry on the same

²Each layer of the old ITS had a material budget of $1.1\% X_0$ [211].

Parameter	Inner Barrel	Outer Barrel
Chip dimensions (mm $\times$ mm)		15×30
Sensor thickness (μm)	50	100
Spatial resolution (μm)	5	10
Detection efficiency (%)		> 99
Fake-hit rate (hits/pixel/event)		$< 10^{-6}$
Operation temperature ($^{\circ}\text{C}$)		20 to 30
TID radiation hardness (krad)	2700	100
NIEL radiation hardness (1 MeV $n_{\text{eq}}/\text{cm}^2$)	1.7×10^{13}	10^{12}

Table 8.1: Sensor requirements for the ITS upgrade for the Inner and Outer Barrels. Taken from Ref. [211].

die. A high-resistivity ($\rho > 1\text{k}\Omega \times \text{cm}$) p^- epitaxial layer serves as a sensitive volume. N-type implants act as collecting diodes for generated electrons, while p-type implants host the circuitry. The n-well of the PMOS transistor is isolated from the epitaxial layer via a deep p-well layer in order to avoid charge collection [213]. The depletion region of the pn-junction formed by the collection n-well diode and the p-type epitaxial layer can be expanded by applying a reverse bias voltage to the diode via highly doped p^{++} substrate. However, most of the sensitive volume remains partially depleted.

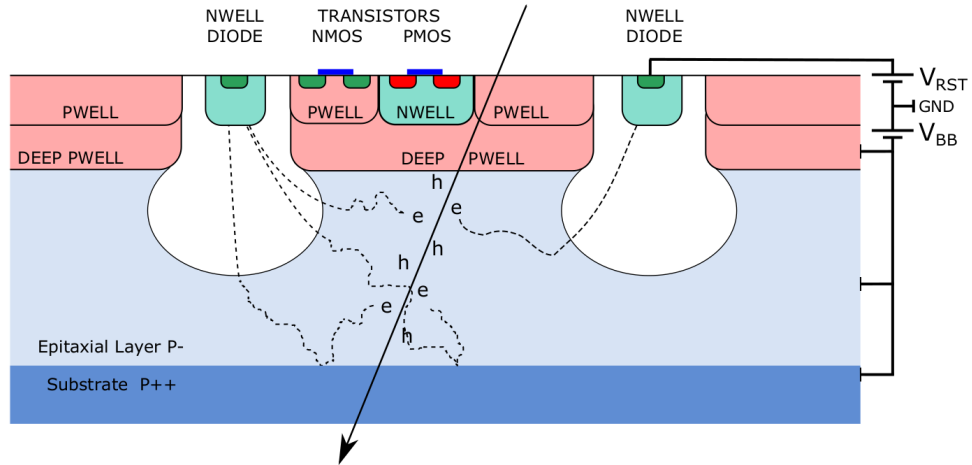


Figure 8.3: Schematic cross section of the ALPIDE sensor. The arrow represents a charged particle passing through the detector, and the dashed lines represent trajectories of the released electrons. Taken from Ref. [211].

8.2 ITS data quality control

During Run 3 and Run 4, the data throughput from the ALICE detector in Pb–Pb collisions is expected to be 3 TB/s, about two orders of magnitude more than in Run

1. Acquiring and analyzing such a large amount of data is challenging for online and offline computer systems. Therefore, the ALICE Collaboration developed a new computing model within the newly developed O² framework to reduce the data volume coming from the detector as early as possible during data taking [214]. The model incorporates two stages of data reconstruction. In the first stage, data reconstruction happens *synchronously* (or online) with data taking. Lossless algorithms are applied to compress the detector data. All non-physical data, such as noise, background from δ -rays, and beam-gas collisions, are rejected, thus further reducing the data volume and keeping only signal events. The second reconstruction stage is performed *asynchronously* (or offline) with data taking. Asynchronous reconstruction aims to provide the final detectors' calibrations and event reconstruction. More details about the ALICE strategy for data recording in Run 3 and Run 4 can be found in Ref. [214].

The O² framework contains several modules, including the global data Quality Control (QC) system. This thesis discusses only the ITS QC system. The ITS is a crucial detector for the ALICE data taking. The QC system quickly and automatically assesses the quality of collected data. It can operate in synchronous and asynchronous modes. Based on user-defined algorithms called *tasks*, the QC system generates Monitoring Objects (MOs) such as histograms, graphs, and values. Data are marked with a *quality flag* set based on comparing MOs with a reference. The automatic check system, which assigns a quality flag to data, particularly helps the central ALICE shifters during data acquisition, when it is necessary to act quickly to fix an issue. During an offline analysis, MOs are checked more thoroughly by detector experts and the Quality Assessment group. In the final stage of the QC, MOs are stored in the Condition and Calibration Data Base (CCDB), where a shifter or a detector expert can access them to check the data quality.

My service work for the ALICE Collaboration was devoted to developing and maintaining several asynchronous ITS QC tasks, particularly the Cluster, Track, and Post-Processing trending Front-End electronics (FEE) tasks. I also developed the post-processing Cluster Occupancy and Average Cluster Size tasks within the QC-analysis framework (not a part of O²) developed by Ivan Ravasenga [215] for an offline data quality analysis. Each task is briefly discussed below.

8.2.1 Asynchronous Cluster and Track tasks

The Cluster and Track QC tasks provide quality assessment for the ITS clustering procedure and stand-alone track reconstruction performance. The asynchronous versions of these tasks supplement the corresponding synchronous (online) tasks written by Artem Isakov [216]. The tasks were designed to obtain MOs generated by the synchronous tasks during the offline data quality analysis. The asynchronous Cluster and Track tasks read

ROOT files stored in the CCDB using the same JSON configuration files as the synchronous tasks. The source code of the asynchronous Cluster and Track tasks is available via the links in Refs. [217], [218]. Figure 8.4 shows some selected examples of plots provided by the tasks.

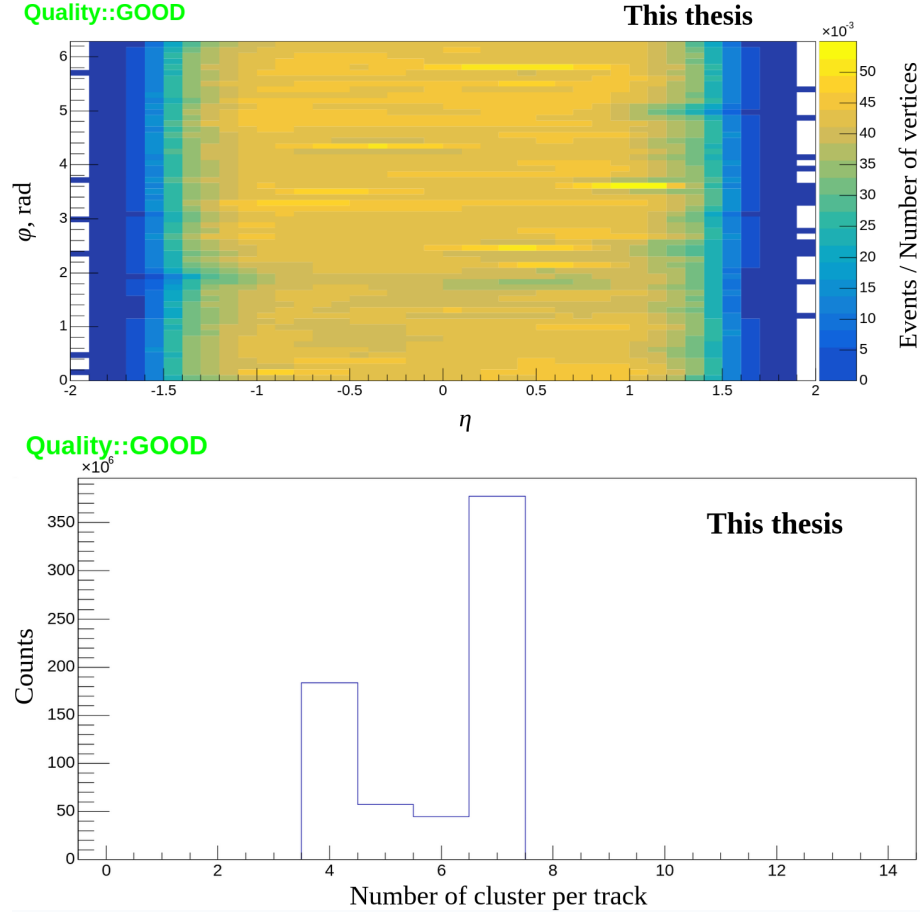


Figure 8.4: Example of MOs reproduced by the ITS asynchronous Cluster and Track tasks in Pb–Pb collisions at $\sqrt{s_{NN}} = 5.36$ TeV (run number: 544013). Top plot: distribution of azimuthal angle vs. pseudorapidity for tracks reconstructed by the ITS. Bottom plot: distribution of the number of clusters per track reconstructed in the ITS. Most reconstructed tracks have 7 clusters, i.e. hit in each ITS layer. Tracks with 4 clusters have 3 hits in the IB layers and 1 hit in Layer 3 (numeration of the layers starts from 0).

8.2.2 Post-processing trending FEE task

Front-end electronics QC task provides comprehensive information about the ITS readout performance, including the number and type of trigger flags, the payload size of detector links, and the number of links in the **WARNING/ERROR/FAULT** state. The data stored in the CCDB can be used to track long-term trends of the detector parameters. The developed post-processing trending FEE task allows tracking changes in the number of links in **WARNING/ERROR/FAULT** state for the IB, OB, and the whole ITS. Figure 8.5 shows an

example of a plot generated by the post-processing trending FEE task. The source code of the task is available via the link in Ref. [219].

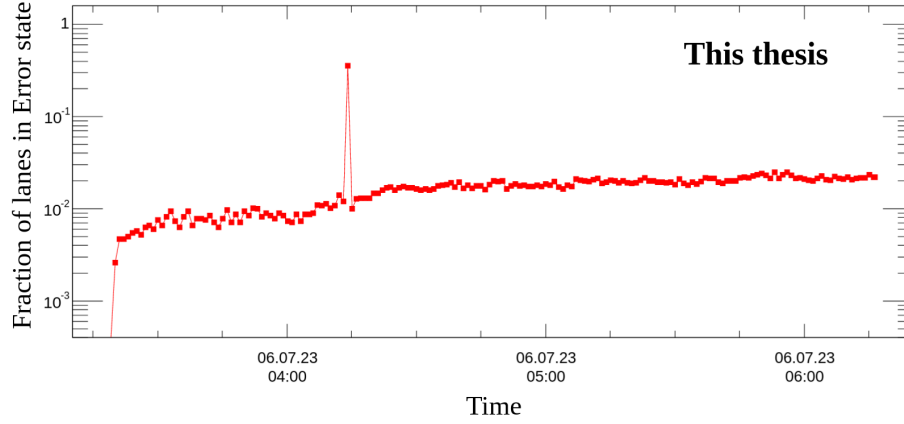


Figure 8.5: Example of MO produced by the ITS post-processing trending FEE task in pp collisions at $\sqrt{s} = 13.6$ TeV (run number: 539331). The plot shows the fraction of lanes in **ERROR** state for the whole ITS as a function of run time. The outlier corresponds to a beam splash event that produced an enormous number of particles, causing the electronics to saturate. As a result, 35% of ITS lanes moved into **ERROR** state. After the beam splash event, the lane recovery procedure was automatically triggered, and most of the lanes returned to working condition.

8.2.3 Post-processing Cluster tasks

These QC tasks provide a summary of the Cluster Occupancy (the number of clusters per event) and the Average Cluster Size for each stave in each ITS layer as a function of the run list. The obtained results are averaged over all ALPIDE chips in a stave. The QC data stored in the CCDB are used as input. Figure 8.6 shows some selected examples of plots provided by the tasks. The source code of the tasks is available via the links in Refs. [220], [221]

8.3 Conclusions

The upgrade of the ITS detector replaced 6 coaxial layers of position-sensitive silicon detectors (pixel, drift, and strip detectors) with 7 coaxial layers equipped with the ALPIDE Monolithic Active Pixel Sensors fabricated in the 180 nm CMOS process. The ALPIDE sensors demonstrated good performance and sufficient radiation hardness to meet the ITS project requirements. The performance of the ITS in Run 3, monitored by the Quality Control software, is stable, and the recorded data have good quality, which is crucial for physics analyses.

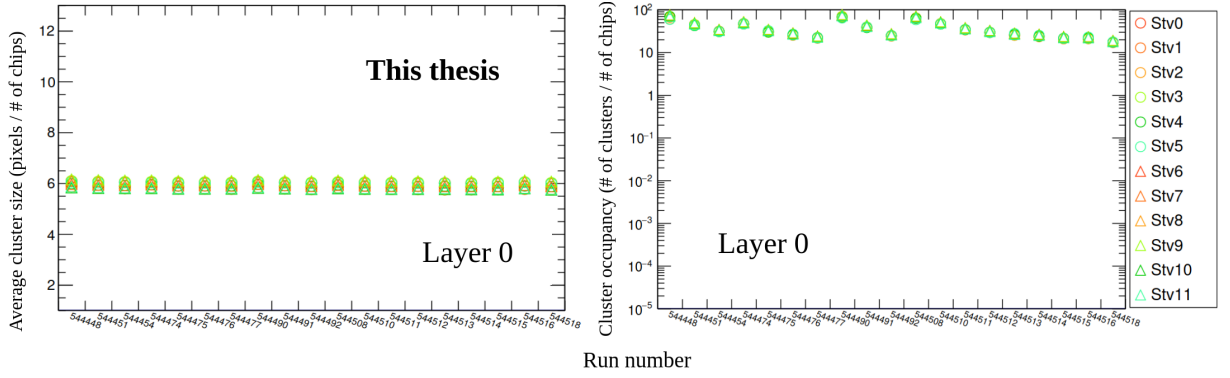


Figure 8.6: Example of MOs produced by the ITS post-processing Cluster tasks in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.36$ TeV (run list: 544448–544518). The data are shown for each stave (“Stv”) of Layer 0 (innermost ITS layer). Left: cluster size averaged over all chips in a stave as a function of run number. Right: cluster occupancy averaged over all chips in a stave as a function of run number.

Chapter 9

ALICE ITS3: first truly cylindrical inner tracker

During the next LHC Long Shutdown 3 in 2026–2028, the ALICE Collaboration plans to replace the three innermost layers of the ITS2 (Sect. 8) with truly cylindrical layers made of wafer-scale, thin ($< 50 \mu\text{m}$) MAPS that will be produced with the 65 nm CMOS imaging process [222]. The upgrade project is called ITS3. To fabricate the wafer-scale sensors, the stitching technology is utilized to “stitch” small reticles and build sensors up to 300 mm in length on a single silicon wafer. The upgrade aims to decrease the material budget of the innermost layers from the present $0.35\% X_0/\text{layer}$ to $0.05\% X_0/\text{layer}$, essentially reducing it to the silicon sensor contribution only. The ITS3 will improve the pointing resolution in the transverse plane by a factor of 2 and increase the reconstruction efficiency for low- p_T tracks [223]. The construction of the ITS3 encompasses many cutting-edge R&D efforts, for instance, production and characterization of the MAPS in the 65 nm CMOS process, fabrication of the stitched wafer-scale MAPS, and development of an ultra-light detector mechanics and a new air cooling system. This chapter provides an overview of the ALICE ITS3 detector, the R&D achievements mentioned earlier, and the author’s contribution to the ITS3 project.

9.1 Upgrade concept

One of the primary motivations for upgrading the current ITS2 detector is the reduction of the material budget of the three innermost layers that form the Inner Barrel (IB) part. Figure 9.1 shows the material distribution of the first IB layer of ITS2 averaged for tracks in $|\eta| < 1$ as a function of azimuthal angle. Note that, to a large extent, the material budget is formed by passive components, such as water cooling, carbon support structures, and flexible printed circuits. Silicon chips make up only about 15% of the

material budget. The distribution is highly non-uniform, containing many peaks that correspond to the overlap of the reinforced structures and the presence of water-cooling pipes.

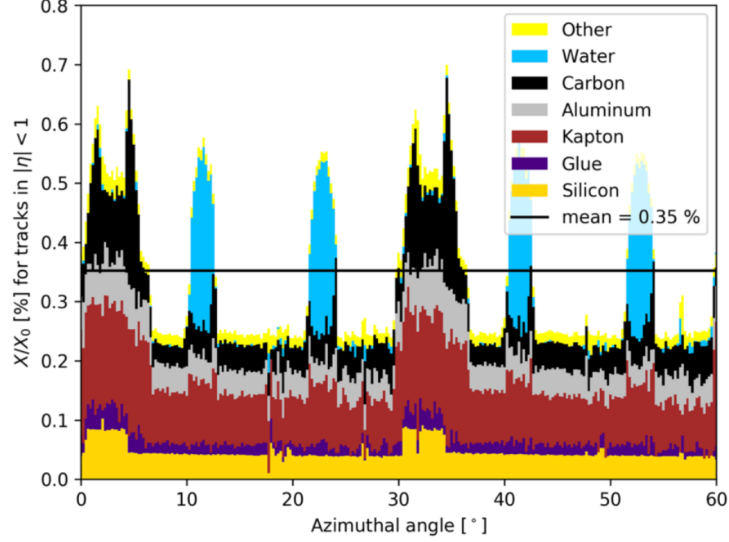


Figure 9.1: Material distribution in the first layer of the ITS2 detector as a function of azimuthal angle. The distribution was averaged for tracks in $|\eta| < 1$. Taken from Ref. [222].

The passive material worsens the decay-vertex reconstruction of low- p_T heavy-flavor hadrons owing to the enhanced rate of multiple scattering. To further improve the precision of the ALICE measurements in Run 4, the material budget of the ITS3 detector will be reduced due to the following design features:

- integration of the power and data buses into a silicon chip, hence the corresponding circuit boards can be removed;
- replacement of the water cooling system by air cooling, and reduction of the chip power consumption relative to the ITS2;
- development of the technology for producing large-area, bent silicon sensors, which form a self-supporting mechanical structure.

The IB of ITS3 will also have three layers, each consisting of only two wafer-scale MAPS bent to form a half-cylindrical shell, see the layout in Fig. 9.2. The sensors will be as large as $26 \times 9 \text{ cm}^2$, and their thickness will be less than $50 \text{ }\mu\text{m}$. The innermost layer will be placed at a radial distance of 18 mm from the interaction point, while for the ITS2, it is 22 mm . The layers will be held in place relative to each other utilizing light-carbon-foam longerons and half-rings.

The wafer-scale MAPS sensors will be produced using one-dimensional stitching technology. The spatial resolution of the future sensors is required to be about $5 \text{ }\mu\text{m}$. As the

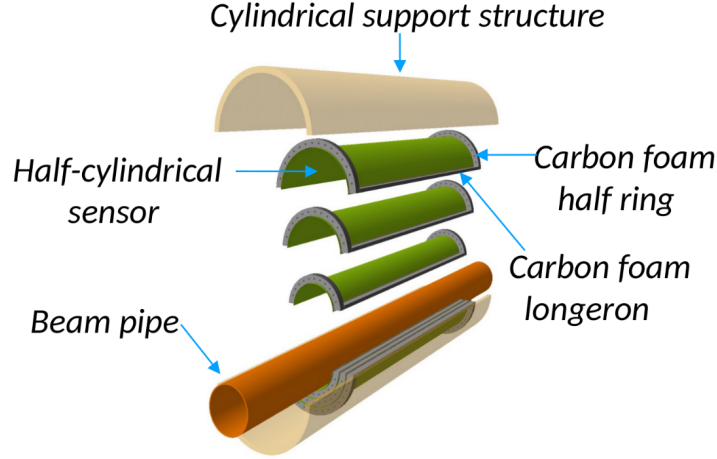


Figure 9.2: Layout of the ITS3 detector. Taken from Ref. [224].

first layer of the ITS3 will be closer to the interaction point, the ITS3 sensors will have to cope with higher radiation loads than the ITS2 sensors. The project requires that the future sensors should be radiation hard at least up to 10^{13} $1 \text{ MeV n}_{\text{eq}}/\text{cm}^2$ (NIEL) and 10 kGy (TID).

The ITS3 upgrade will enhance the ALICE capabilities to reconstruct open-heavy-flavor hadrons. For instance, the simulated significance of the signal from the decay of $B_s^0 \rightarrow D_s^- + \pi^+$ measured with the ITS3 in 0–10% central Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.5 \text{ TeV}$ is 1.4 times higher at $p_{\text{T}} = 10 \text{ GeV}/c$ than for the ITS2. Moreover, the ITS3 extends this measurement to lower p_{T} region, which is unattainable for the ITS2. The TDR [225] discusses the detector performance in more detail.

9.2 R&D status

The roadmap of the ITS3 R&D is split into several parallel activities: the characterization of bent sensors, the development of detector mechanics and air cooling system, and the design and characterization of new sensors. Each of the activities will be briefly discussed below.

9.2.1 Characterization of the bent ALPIDE sensors

To prove the concept of operating bent sensors, the ALICE Collaboration built and tested a setup called μITS3 , see the left photo in Fig. 9.3. The setup consists of 6 ALPIDE monolithic active pixel sensors with a thickness of $50 \text{ }\mu\text{m}$ [226], bent to the target radii of the ITS3. The right plot in Fig. 9.3 shows the detection inefficiency of the bent ALPIDE sensors for minimum-ionizing particles as a function of charge threshold. The project requires the detection inefficiency to be less than 10^{-2} (in other words, the detection

efficiency to be greater than 99%). One can see that at the nominal operating point of 100-electron threshold, the inefficiency is less than 10^{-4} . Furthermore, the detection inefficiency does not depend on the sensor curvature. The obtained results are consistent with the performance of a flat ALPIDE sensor. Complementary studies also showed that the bending does not affect the spatial resolution of the sensors. Similar studies were also carried out with a single bent ALPIDE sensor, see details in Ref [227].

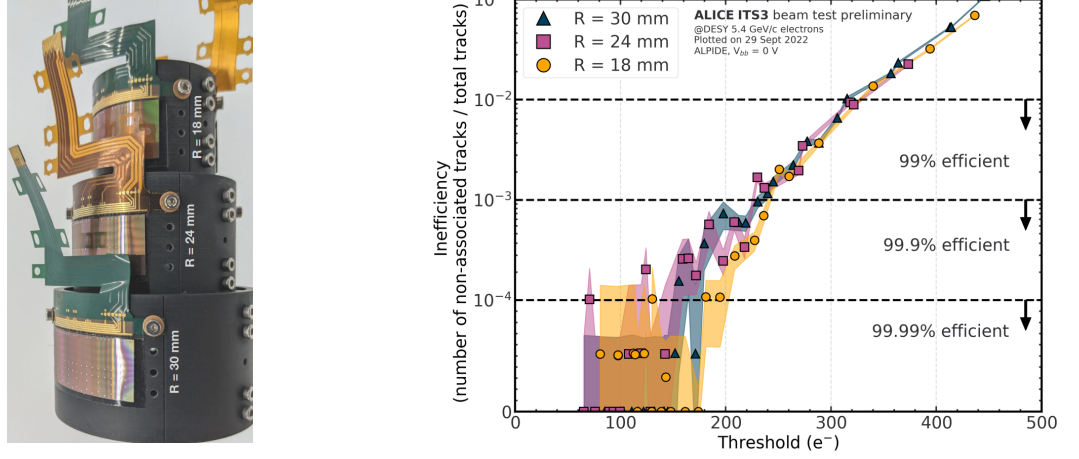


Figure 9.3: Left: μ ITS3 setup. Right: Inefficiency in detecting a minimum ionizing particle by the ALPIDE sensors bent to different radii as a function of charge threshold. Taken from Ref. [225].

9.2.2 Detector mechanics

The wafer-scale sensors are bent to a half-cylindrical shape of a given radius utilizing a mandrel. The half-cylindrical shape of the sensors will be fixed by carbon-foam half-rings and longerons glued to the sensors. There are two types of proposed carbon-foam materials, which primarily differ in density and thermal conductivity. The half-rings will be made of carbon foam with higher density ($0.2\text{--}0.26\text{ kg/dm}^3$) and higher thermal conductivity ($> 17\text{ W/m/K}$) since it is expected that most of the heat will be produced at the sensor endcaps. Thus, the half-rings act as a heat sink. Moreover, they provide mechanical support. In contrast, the longerons will be made of lower-density carbon foam ($\approx 0.045\text{ kg/dm}^3$) to reduce the material budget. This carbon foam also has lower thermal conductivity ($> 0.033\text{ W/m/K}$) since the sensor matrix generates a small amount of heat. Figure 9.4 shows a full-size breadboard model of the ITS3, assembled from dummy silicon wafers.

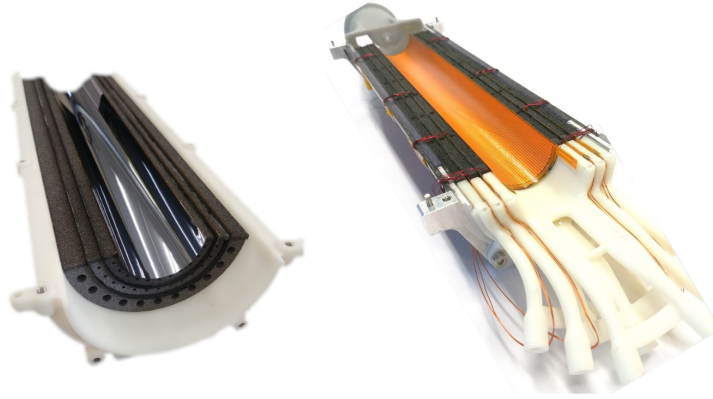


Figure 9.4: Left: One of the ITS3 breadboard models for mechanical studies. Right: The breadboard model equipped with an air distributor and thermal heater for mechanical and thermal tests in a custom wind tunnel. Taken from Ref. [225].

9.2.3 Air cooling system

Various breadboard models, similar to the one in Fig. 9.4, equipped with flow distributors and sensors for temperature and displacement measurements, are used in tests with a custom wind tunnel to determine their thermal and mechanical properties. The project requires an air cooling system that can keep the operating temperature of the sensors below 30 °C. It was shown that in order to achieve the stated performance, an airflow of 8 m/s is sufficient [228]. Measurements of mechanical vibrations caused by the airflow showed that the amplitude of these vibrations is less than 0.5 μm . This level of vibrations is negligible when compared to the expected spatial resolution of the sensor. The level of vibrations is expected to be further reduced by optimizing the geometry of holes in the carbon-foam half-rings.

9.2.4 65 nm sensor design and characterization

The 65 nm CMOS process is a novel technology for sensor production in high-energy physics. Producing a new pixel-sensor prototype involves a comprehensive design verification chain and simulation program. Generally, the characterization of new sensors incorporates tests in a laboratory and at a test beam facility. In a laboratory, many tests with a sensor are carried out, for instance, measurements of sensor power consumption and noise rate, optimization of sensor settings such as operation currents and voltages, and energy calibration with a radioactive source or X-ray tube. A test beam campaign measures the sensor's characteristics that can be solely determined at a high-energy beam. For instance, these characteristics are hit detection efficiency and spatial and time resolution. The hit detection efficiency can be studied for an entire sensor or as a function of hit coordinate, so-called in-pixel efficiency.

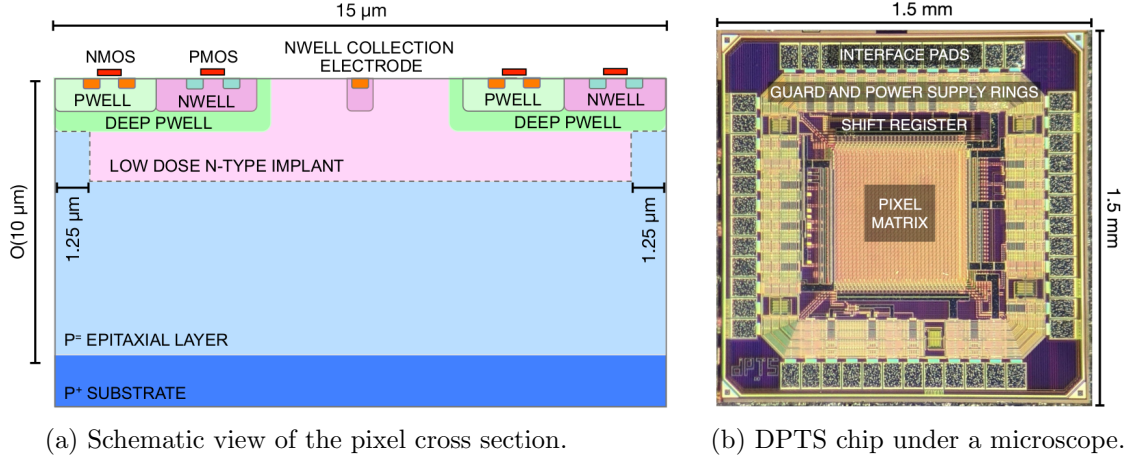


Figure 9.5: A schematic cross section of a DPTS pixel and a photo of the chip under a microscope. Taken from Ref. [231].

The design of the pixels for the ITS3 introduced a new feature, a low-dose n-type implant layer with a gap [229], which allows complete depletion of the sensitive volume and increases the radiation hardness of the sensors. The development of the ITS3 sensor prototypes commenced with the production of digital and analogue pixel test structures, DPTS and APTS, respectively. To optimize the new sensor design for the best particle-detection performance, the sensors were manufactured in 4 process splits [230], gradually modifying the doping levels of various implants (Fig. 9.5a). The DPTS sensor has a 32×32 pixel (Fig. 9.5b) matrix with an amplifier and a discriminator in each pixel and digital readout with time encoding [231]. The DPTS sensors were characterized in terms of charge collection efficiency, detection efficiency for minimum-ionizing particles, fake-hit rate, and radiation hardness [231]. The published results showed that the DPTS sensor, irradiated to the radiation levels expected for the ITS3, has a detection efficiency above 99% with a low but measurable fake-hit rate. The DPTS chips were irradiated with non-ionizing (NIEL), ionizing (TID), and combined doses. I and my colleagues carried out irradiation of the DPTS chips at the cyclotron of the Nuclear Physics Institute of CAS with a 30 MeV proton beam (TID + NIEL). My contribution to this measurement concerned the preparation and assembly of a setup for the chip irradiation and data monitoring during the irradiation.

9.2.4.1 Analogue Pixel Test Structure

The APTS is a prototype chip with a matrix of 4×4 pixels. Each pixel has an individually buffered analogue output connected to an output pad. The pixel matrix is surrounded by a two-pixel wide ring of dummy pixels to minimize the distortions of the electric field due to edge effects. The APTS can be operated at the reverse substrate voltages from 0 to -5 V.

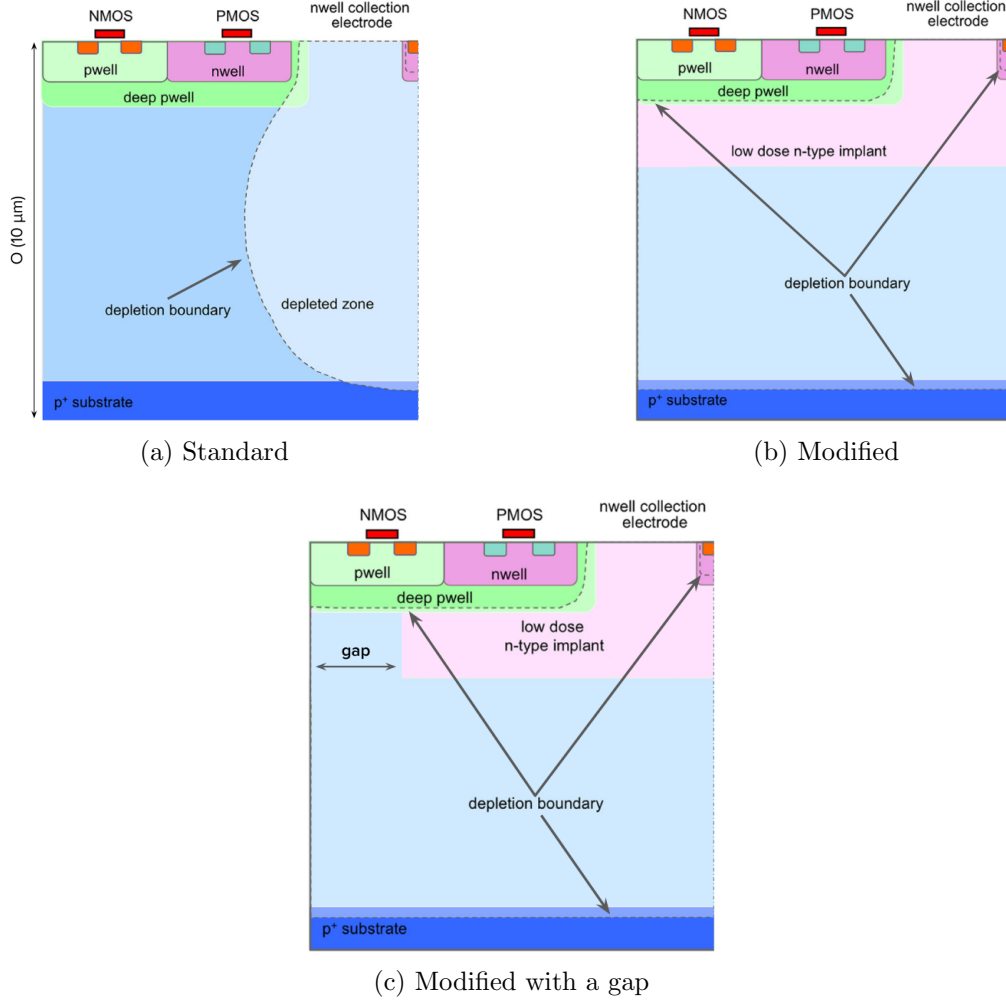


Figure 9.6: A schematic cross section of an APTS pixel produced in three different designs. Taken from [232].

Various prototypes of the APTS chips were manufactured, differing in the pixel pitch size (10, 15, 20, and 25 μm), geometry of the n-well/p-well, and sensor design. Three sensor designs were considered: *standard*, *modified*, and *modified with a gap* (Fig. 9.6) [229], [232]. In the *standard design* (ALPIDE-like design), the depletion zone has a “balloon” shape, see Fig. 9.6a. The regions close to the pixel edges are not depleted. The charge generated outside the depletion zone is primarily collected via charge diffusion, which is a relatively slow process compared to the charge drift. Trapping centres created by exposure to non-ionizing radiation most likely capture a slow-diffusing charge, leading to a lower charge collection efficiency. The *modified designs* incorporate a novel modification, a deep low-dose n-type implant. Implementing this layer creates a deep planar junction extending over the pixel area. As a result, by applying the reverse substrate voltage, the volume of the epitaxial layer is fully depleted. The dominant mechanism for charge collection is the charge drift; thereby, the charge collection efficiency is higher compared to the *standard design*. The *gap modification* contains a 2.5 μm wide gap introduced into the

deep n-type implant at pixel borders. Due to the gap, the lateral electric field increases, which drives the charges away from the pixel border towards the collection electrodes. Thus, charge sharing between neighbouring pixels is reduced.

I worked on the analysis of test beam data, where APTS sensors were tested at:

1. Mainz Microtron (MAMI) with a beam of electrons with an energy of 855 MeV;
2. CERN SPS with a beam of negative pions with an energy of 120 GeV.

The obtained results will be discussed in detail below.

Telescope setup Test beam measurements are carried out using a setup called a telescope. A scheme of a typical telescope the ITS project uses in test beams with APTS is shown in Fig. 9.7. The telescope setup is installed in a light-proof box to eliminate noise generated by ambient light. A sensor being tested, the so-called device under test (DUT), is placed in the telescope's centre. The ALPIDE sensors provide track reconstruction independently of a hit in the DUT. Generally, two scintillator detectors connected in a coincidence mode are added into the setup to suppress background events. The beam must pass through all detectors in the telescope. Since the size of the APTS chip is tiny compared to the size of the ALPIDE sensor and trigger detectors, the preliminary alignment of the setup relative to the beam is a non-trivial objective. Therefore, another APTS chip installed behind the DUT is commonly used as a trigger in test beams with an APTS sensor. The distance between the detectors is measured using a ruler. By default,

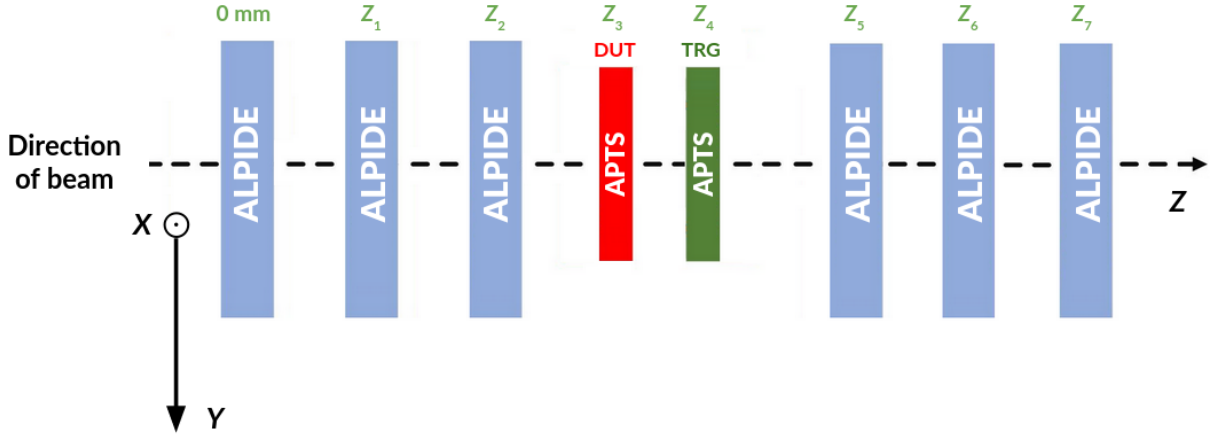


Figure 9.7: A scheme of a telescope with an APTS sensor as the DUT; top view. The z coordinate of each detector is determined relative to the ALPIDE sensor, which is the first in the setup along the beam direction.

the detector planes are assumed to be aligned on the z axis. The alignment on the x and y axes cannot be done manually since a misalignment by a fraction of a millimetre can

correspond to a shift of several pixel pitches. The ALPIDE sensors and DUT are aligned on the x and y axes using software, as discussed below.

Test beam data analysis Data from test beams are analyzed using the Corryvreckan reconstruction framework [233]. The Corryvreckan is written in the C++17 standard, allowing quick test-beam-data reconstruction. The framework is based on a modular concept. Each step in the reconstruction chain is performed by a corresponding module, which users can customize to suit their needs. Before starting the Corryvreckan analysis, one must prepare a configuration file with the telescope geometry. This file must contain information for each detector, such as its type, position coordinates in the telescope, angle orientation relative to the coordinate axes, etc. The transverse coordinates of each detector are initially set to $(x, y) = (0, 0)$ mm. The Corryvreckan will automatically update the coordinates once the telescope is aligned. For an unbiased DUT analysis, the Corryvreckan allows the selection of only those events containing one reconstructed track. A typical reconstruction chain consists of the following sequence of actions:

1. **Data loading and decoding** (`EventLoaderEUDAQ2` module): test beam data is collected by the EUDAQ2 software framework [234] and encoded in a binary format. The Corryvreckan performs binary-data decoding to obtain information about pixels that were hit;
2. **Masking of noisy pixels** (`MaskCreator` module): some sensor pixels constantly generate fake signals. When analyzing data, noisy pixels must be excluded (or masked) from consideration because their signal significantly distorts the measured signal. A frequently used approach for noisy pixel masking is based on a straightforward cut on a global pixel firing frequency, which masks pixels with a hit rate larger than `user defined cut  $\times$  the mean global hit rate`;
3. **Clustering**: a particle beam passing through a pixel detector fires not a single pixel but multiple neighbouring pixels. The Corryvreckan reads the hit pixel information and performs pixel clusterization. The cluster is a group of adjacent pixels inside a sensor matrix centred around the pixel with the largest signal amplitude. It should be noted that depending on whether a pixel detector has a digital or analogue output, an appropriate Corryvreckan module must be called. The ALPIDE sensor has a digital output, so one needs to use the `ClusteringSpatial` module. In the case of the APTS chip, the `ClusteringAnalog` module is utilized;
4. **Telescope prealignment** (`Prealignment` module): the Corryvreckan performs a preliminary alignment of the telescope planes (only translational alignment) based

on information from cluster correlations. It is assumed that the beam impinges perpendicularly on the sensors' surfaces. A cluster has a centre defined by the (x, y) coordinates on the detector plane, where it was reconstructed. The cluster correlations for each detector are considered relative to a reference ALPIDE sensor, whose "reference" role is assigned before the prealignment. In general, other telescope planes are misaligned relative to the reference plane. As a result, the centre positions of the corresponding clusters are shifted in the local coordinate frame associated with each detector plane. The Corryvreckan calculates the amount of the shift on x and y axes and applies the corresponding corrections to the position coordinates of the detectors in the global frame;

5. **Tracking and fine telescope alignment:** the Corryvreckan performs a fine telescope alignment using tracks reconstructed with the **Tracking4D** module. A straight line connects any cluster on the first ALPIDE sensor with any cluster on the last ALPIDE sensor. Knowing the coordinate, where the interpolated track intersects the plane of the subsequent detector, the nearest cluster is searched based on user-defined spatial and time cuts. Hits are required on each ALPIDE sensor to reconstruct high-quality tracks. The distribution of the spatial differences between the coordinates of the interpolated track and the centre of the nearest cluster is called track residuals. If the detectors are misaligned on the x and y axes, then the mean of track residuals is shifted relative to 0 μm , and the Corryvreckan will automatically adjust the detector position on the x and y axes in the global frame (**AlignmentMillepede** module). The distribution of track residuals has a finite width, which depends on the tracking resolution of the telescope $\sigma_{\text{telescope}}$. The $\sigma_{\text{telescope}}$ resolution depends on beam energy and detector material budget. In addition to the straight-line-track model, the Corryvreckan includes the General Broke Lines (GBL) model, which takes into account the deviation of a particle direction from a straight track due to multiple scattering within detector material [235]. By default, the GBL model is used in all analyses;
6. **DUT analysis:** after the telescope is aligned, one can study the DUT characteristics, such as spatial resolution (**AnalysisDUT** module) and hit detection efficiency (**AnalysisEfficiency** module). The DUT spatial resolution along the x and y axes is obtained from the RMS of the corresponding track residuals. It is important to note that the obtained resolution is a combination of both the telescope tracking resolution and the DUT spatial resolution,

$$\sigma_{x, y} = \sqrt{\sigma_{\text{telescope}}^2 + \sigma_{\text{DUT } x, y}^2}. \quad (9.1)$$

The $\sigma_{\text{telescope}}$ can be obtained from Monte Carlo-based tracking simulations, which account for a given telescope geometry (distance between detectors, material budget) and beam energy [236], [237].

Hit detection efficiency represents the probability of detecting a particle passing through the DUT,

$$\varepsilon = \frac{C}{T}, \quad (9.2)$$

where C is the number of tracks with an associated cluster on the DUT; T is the number of all reconstructed tracks. If the telescope tracking resolution is less than the size of DUT pixel pitch, then hit detection efficiency can be studied as a function of a position within a pixel in the DUT plane. This characteristic is called in-pixel detection efficiency.

MAMI test beam The beam test was carried out at the Mainz Microtron [238], which provided an electron beam with an energy of 855 MeV. The telescope setup differed slightly from that shown in Fig. 9.7. There was no external trigger in the setup. The DUT was utilized as a trigger. Thus, the DUT hit detection efficiency could not be measured, yet the spatial resolution could be determined. It is important to stress that a firmware threshold was imposed on the trigger during the test beam, affecting the data. This point will be discussed later. The DUT analysis has been done by associating clusters on the ALPIDE detectors to the reconstructed tracks passing within a search window of a radius of 75 μm around the track intercept point on the ALPIDE plane. Solely those tracks that passed the 4 central pixels of the APTS sensor were considered to avoid the impact of the border effects.

Different prototypes of APTS chips are labelled with abbreviations that encode the type of buffer (for more details, see Ref. [225]), pitch size, and the type of modification design. For instance, “AF15P” stands for a chip with a pixel pitch size of 15 μm and the gap modification. Figure 9.8 shows the spatial resolution for the AF15P chip as a function of the charge threshold for different values of the reverse substrate voltage V_{sub} . The presented spatial resolution is defined as $(\sigma_x + \sigma_y)/2$. The telescope tracking resolution was estimated using simulations [236] and found to be 4.12 μm for both x and y axes. The dashed horizontal line represents the binary resolution determined solely by the pixel size. The binary resolution is the worst upper limit on the position resolution in the case when no additional information about the hit has been utilized, such as cluster distribution or incidence angle of the particle track producing the hit [239]. The right y axis shows the dependence of the average cluster size on the charge threshold. The data show that the average cluster size decreases with the increasing charge threshold. This trend is expected since a high threshold value excludes adjacent pixels from the clusterization, thus reducing

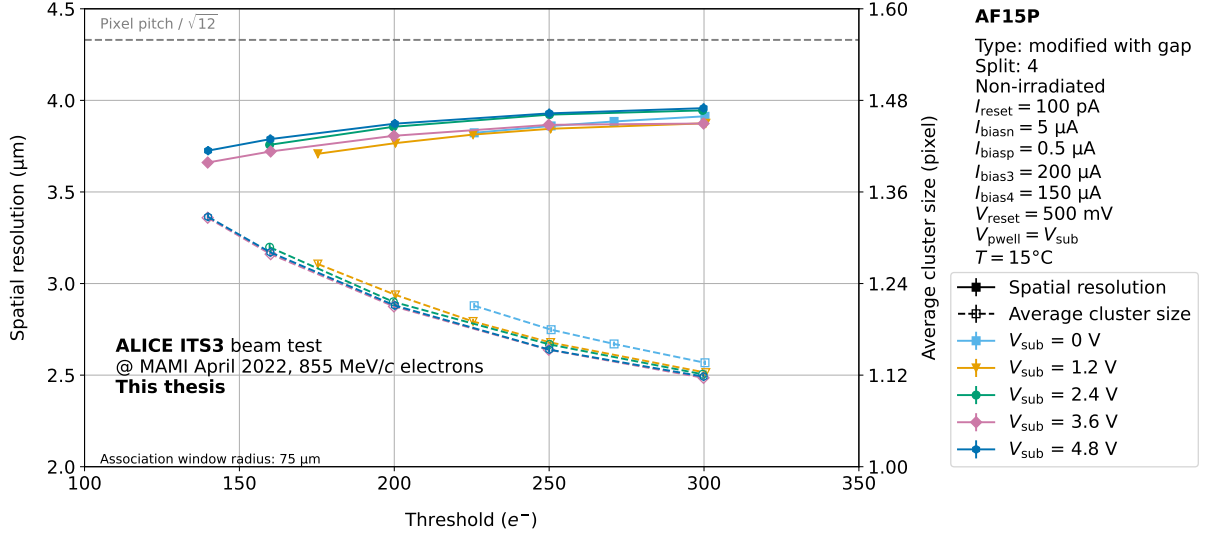


Figure 9.8: Mean spatial resolution (solid lines) and average cluster size (dashed lines) vs. charge threshold at different V_{sub} . The information about the chip and the operation setting are listed in the right-hand side legend. The dashed gray line represents the binary resolution limit; see text for further details.

the cluster size. The charge generated within the sensitive volume thermally diffuses, thus spreading the charge cloud transversely over neighbouring pixels (charge sharing effect) with RMS of $\sigma \sim 1/\sqrt{V_{\text{sub}}}$ [240]. Therefore, the average cluster size should decrease with increasing the reverse substrate voltage. The data support this trend. The largest difference is observed between the data for $V_{\text{sub}} = 0$ V and $V_{\text{sub}} = 1.2$ V because of depletion process. Work [241] demonstrated that estimated spatial resolution can be improved beyond the binary resolution due to charge transverse diffusion. In other words, the spatial resolution should grow with increasing the cluster size. The data demonstrate that the spatial resolution deteriorates as the average cluster size decreases. In addition, the spatial resolution deteriorates with increasing the reverse substrate voltage.

Figure 9.9 shows the comparison of the spatial resolution for the AF15P chip at $V_{\text{sub}} = 1.2$ V measured in the MAMI and SPS test beams. In the test beam at SPS, the particle beam consisted of negative pions with an energy of 120 GeV. The telescope tracking resolution was $2.11 \mu\text{m}$ for both x and y axes. Although the distributions of average cluster size are similar for both test beams, the distributions of the spatial resolution are noticeably different at relatively low charge thresholds. This observation implies that the spatial resolution depends not solely on the cluster size. Since the measurements in both experiments were carried out with the same chip, the difference may arise, for instance, from a different signal-to-noise ratio [241]. The same results for the AF20P chip are shown in Appendix C.

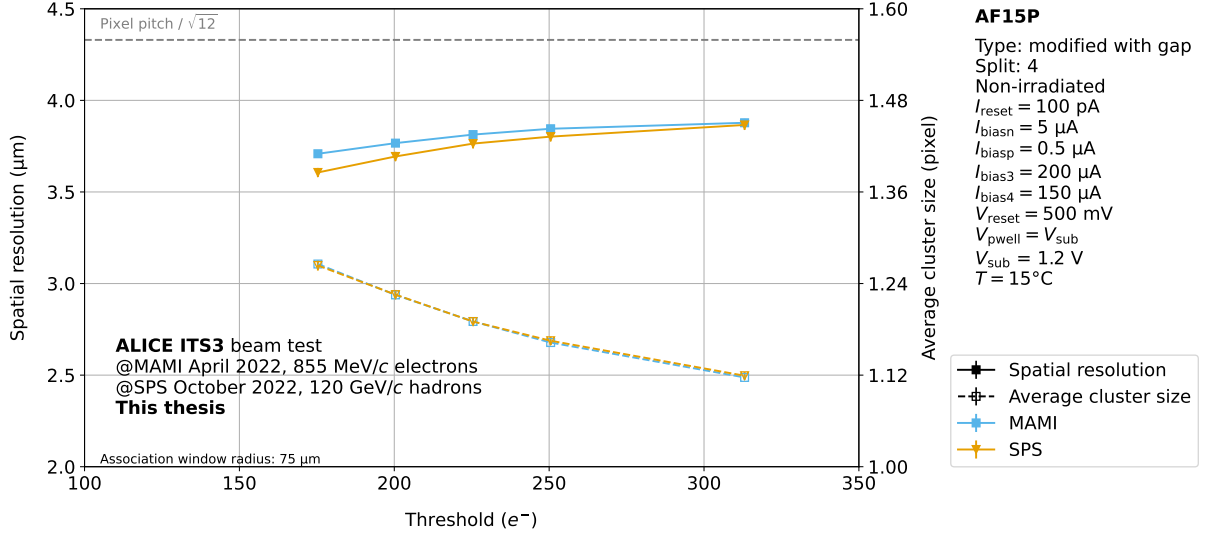


Figure 9.9: Mean spatial resolution (solid lines) and average cluster size (dashed lines) vs. charge threshold at $V_{\text{sub}} = 1.2 \text{ V}$ measured in the MAMI and SPS test beams. The information about the chip and the operation setting are listed on the right-hand side legend. The dashed gray line represents the binary resolution limit.

In-pixel detection efficiency measurements at SPS Figure 9.10 shows the detection efficiency of the AF25P sensor as a function of reconstructed track position relative to the centre of the nearest pixel for the charge threshold values of 120 e^- and 180 e^- . The ABC triangle in the left plot represents how detection efficiency changes moving from the pixel centre (point A) to the nearest border (point B) and then to the pixel corner (point C). The right plot in Fig. 9.10 shows these changes as a function of the distance along the path. By increasing the charge threshold, hit detection efficiency decreases, i.e. particle hit does not produce enough signal amplitude to exceed the threshold. One can see that the detection efficiency decreases when the particle track is further away from the collection electrode located at point A. The maximum decrease in detection efficiency is observed in the corner of the pixel, where the charge-sharing effect is maximum due to the decrease in the electric field with the distance from the pixel centre.

9.2.4.2 Monolithic stitched sensor

In 2023, the ITS3 project got the first stitched-sensor prototypes called the MOlonithic Stitched Sensor (MOSS) and MOlonithic Stitched sensor with Timing (MOST). Further, only the MOSS chip will be discussed. The MOSS is a 6.7-megapixel sensor, having a size of $1.4 \times 25.9 \text{ cm}^2$, see Fig. 9.11. It contains ten repeated sensor units (RSUs), which are basically the largest individual sensors that could fit in the design reticle and which are stitched together. Each RSU comprises two half-units (HUs), labelled top and bottom. Each half-unit contains four matrices, also called regions, with a pixel pitch of $22.5 \text{ } \mu\text{m}$

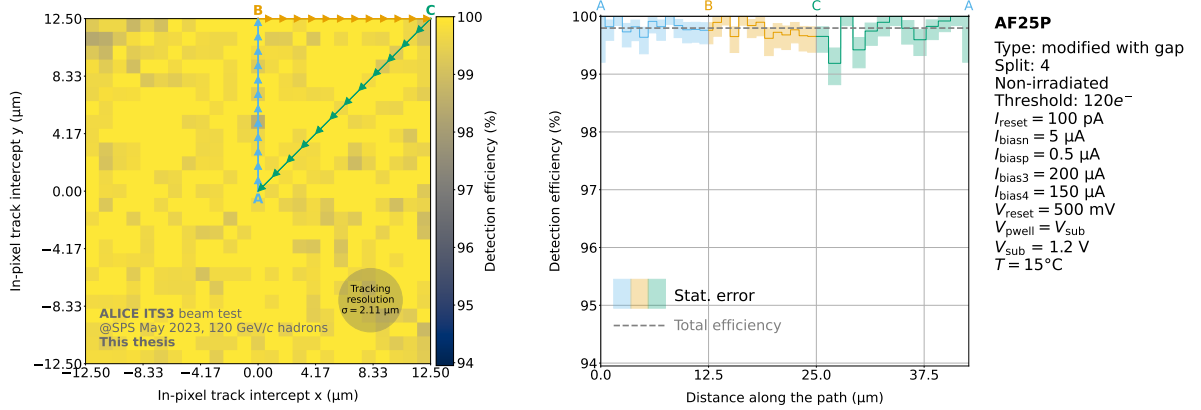
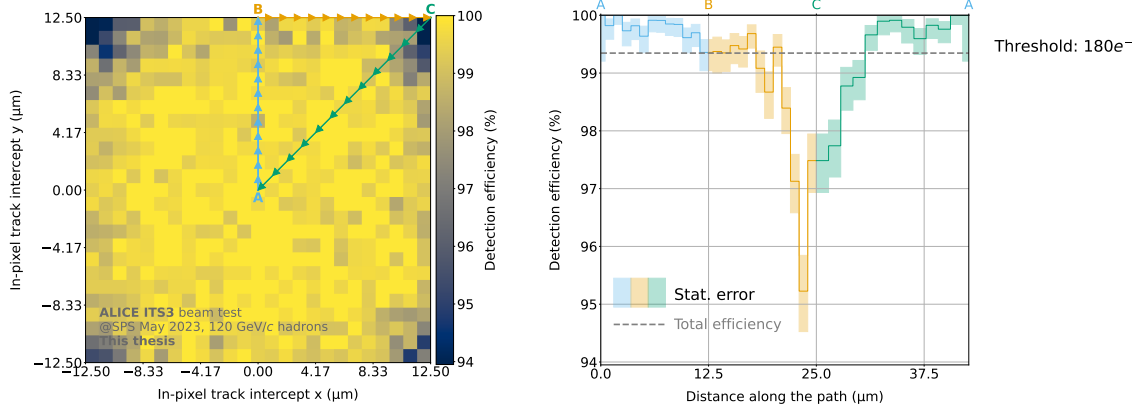

 (a) Charge threshold: 120 e^-

 (b) Charge threshold: 180 e^-

Figure 9.10: In-pixel hit detection efficiency of the AF25P sensor with the charge threshold set to 120 e^- and 180 e^- . Left plot: heat map of in-pixel detection efficiency. Right plot: change of detection efficiency as a function of distance along the path ABC . The information about the operation setting is listed on the right-hand side legend. The tracking resolution is represented by the gray circle with a radius of $\sigma_{\text{telescope}} = 2.11 \text{ } \mu\text{m}$.

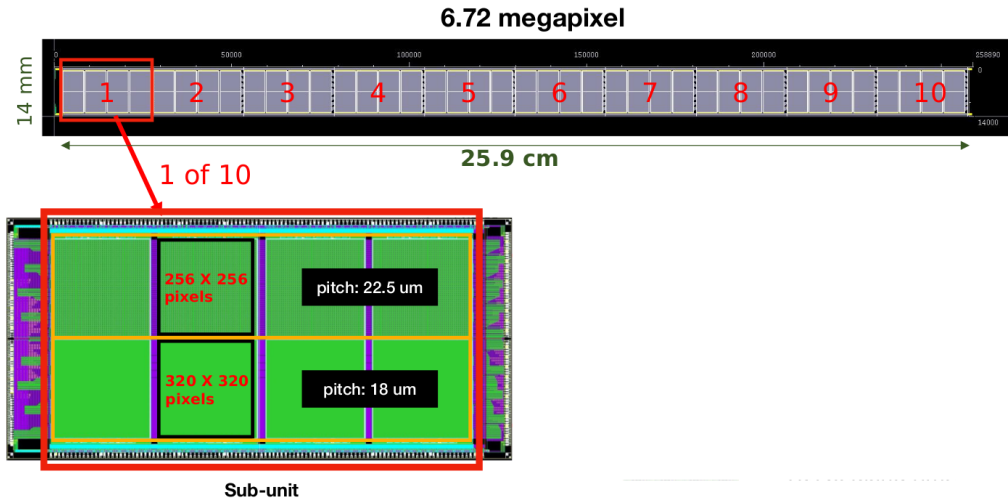


Figure 9.11: Structure of the MOSS chip. The enumerated rectangles represent the repeated sensor units (RSUs). Figure courtesy ITS3 project Work Package 2.

at the top and $18\text{ }\mu\text{m}$ at the bottom. The corresponding number of pixels in each of the top/bottom regions is $256 \times 256/320 \times 320$. The different pitch sizes for the top and bottom half-units are used to test two designs of pixel architecture.

The MOSS chip is operated by a setup, which sketch is shown in Fig. 9.12. The MOSS

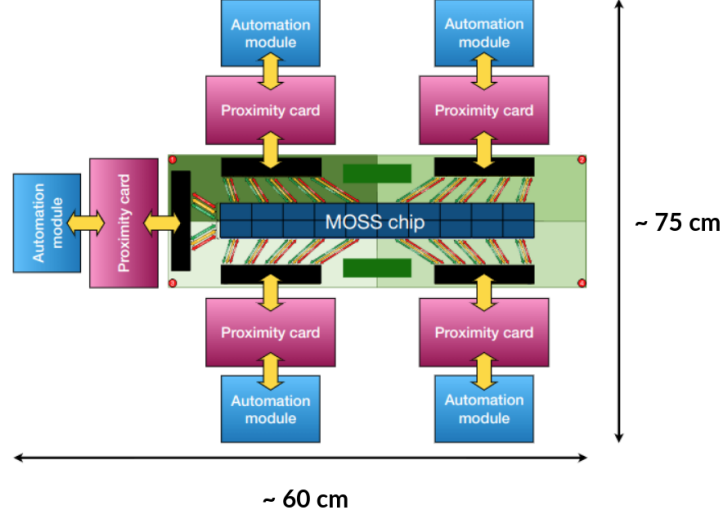


Figure 9.12: Sketch of MOSS test setup. Figure courtesy ITS3 project Work Package 3.

chip is bonded on a carrier board, allowing chip operation via five 560-pin connectors located along its perimeter. Four connectors operate four quadrants of the MOSS chip, each comprising five independently operable half-units. The fifth connector on the left short edge operates the chip as two halves via the stitched backbone. Five proximity cards are connected to the carrier board via the connectors to provide power supply, control, and readout. The proximity cards are steered by automation modules that interface the sensor control and readout with a computer.

The MOSS chips were tested in test beams to characterize the pixel structure. I will discuss below the results from the beam test at the CERN Proton Synchrotron (PS), which I participated in.

PS test beam The CERN Proton Synchrotron provided a beam of negative pions with an energy of 10 GeV. A photo of the telescope setup is shown in Fig. 9.13. Events were triggered by two scintillator detectors connected in a coincidence mode. The triggers were misaligned relative to each other to increase their overlap area, thus increasing the trigger efficiency. Due to the large size of the MOSS test system (Fig. 9.12), not all RSUs could be tested with the available telescope setup. The results for RSU T7 (shorthand notation for top half-unit 7) will be discussed in this section, while the results for RSU B4 (bottom half-unit 4) are presented in Appendix D.

The DUT analysis has been done by associating clusters on the ALPIDE detectors to the reconstructed tracks passing within a search window of a radius of $150\text{ }\mu\text{m}$ around

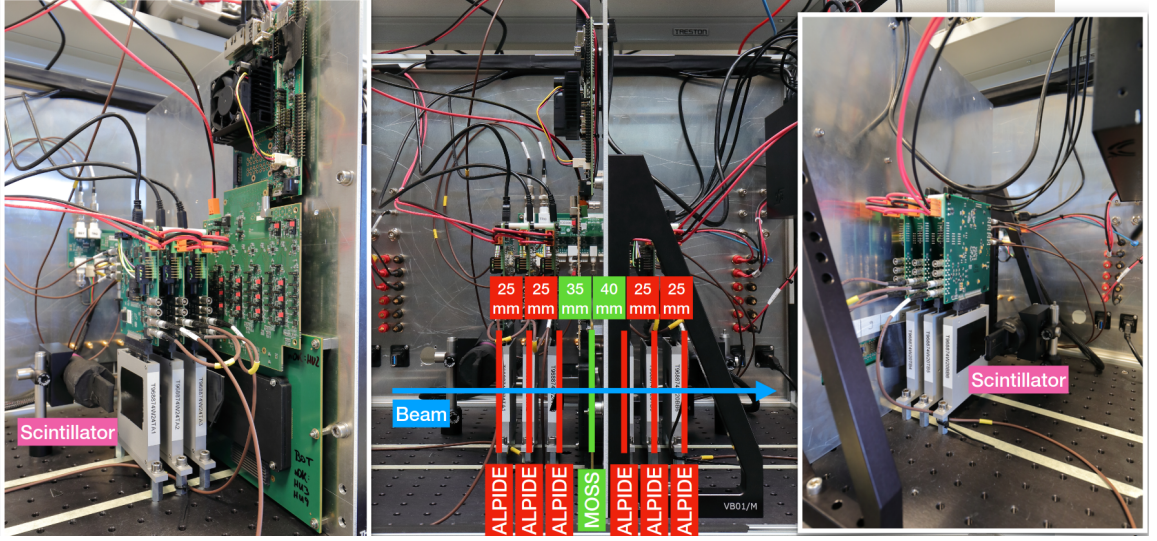


Figure 9.13: Photo of the telescope setup used in the PS test beam with the MOSS chip.

the track intercept point on the ALPIDE plane. Figure 9.14 shows the dependence of the hit detection efficiency of the four regions of the RSU T7 on V_{casb} . The voltage V_{casb} is inversely proportional to the charge threshold. The right y axis corresponds to the

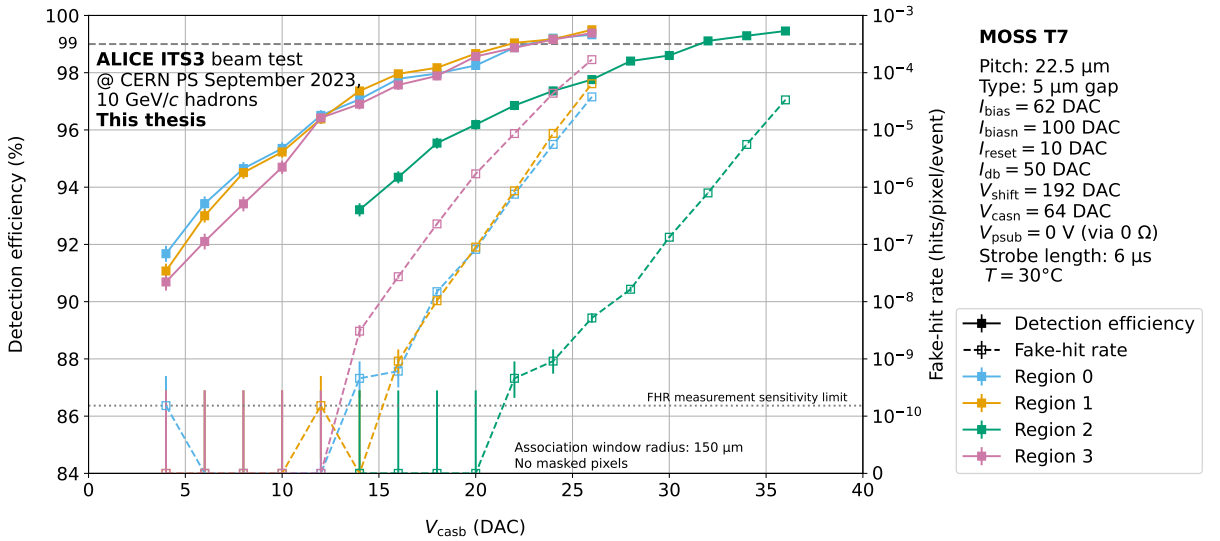


Figure 9.14: Hit detection efficiency (solid lines) and fake-hit rate (dashed lines) vs. V_{casb} for the RSU T7 of the MOSS chip. The information about the chip operation setting is listed in the right-hand side legend. Zero reverse substrate voltage was applied. The dashed line represents 99% detection efficiency limit required by the ITS3 project. The dotted line shows the fake-hit rate sensitivity limit determined by a single-fired pixel.

fake-hit rate. The fake-hit rate shows the number of matrix pixels fired without a particle hit per the number of pixels and events. The fake-hit rate was measured separately from the test beam in a laboratory for 100k events. It can be seen that the detection efficiency and the fake-hit rate data for the region 2 are shifted relative to the data for the other

regions by 10 DAC units of V_{casb} . This shift relates to the V_{casb} threshold for this region. During the fake-hit rate measurements, the V_{casb} threshold, at which more than 0.1% of the pixels (66 pixels) are fired, was determined for each region. The value of the V_{casb} threshold for the region 2 was about 10 DAC units higher than for other regions. Thus, for the region 2, the V_{casb} scan range was set to be 10 DAC units higher than for the other regions. It is worth noting that this feature is inherent only for the top half-units. The project requires 99% detection efficiency, which is achievable for all regions, however, in a very narrow range of V_{casb} but at the same time with the fake-hit rate higher than 10^{-6} hits/pixel/event.

Figure 9.15 shows the mean spatial resolution measured as a function of V_{casb} . The right y axis shows the average cluster size. Despite the V_{casb} shift for the region 2, the spatial resolution and average cluster size data are similar for all regions. Detailed discussion about spatial resolution was provided in Sect. 9.2.4.1.

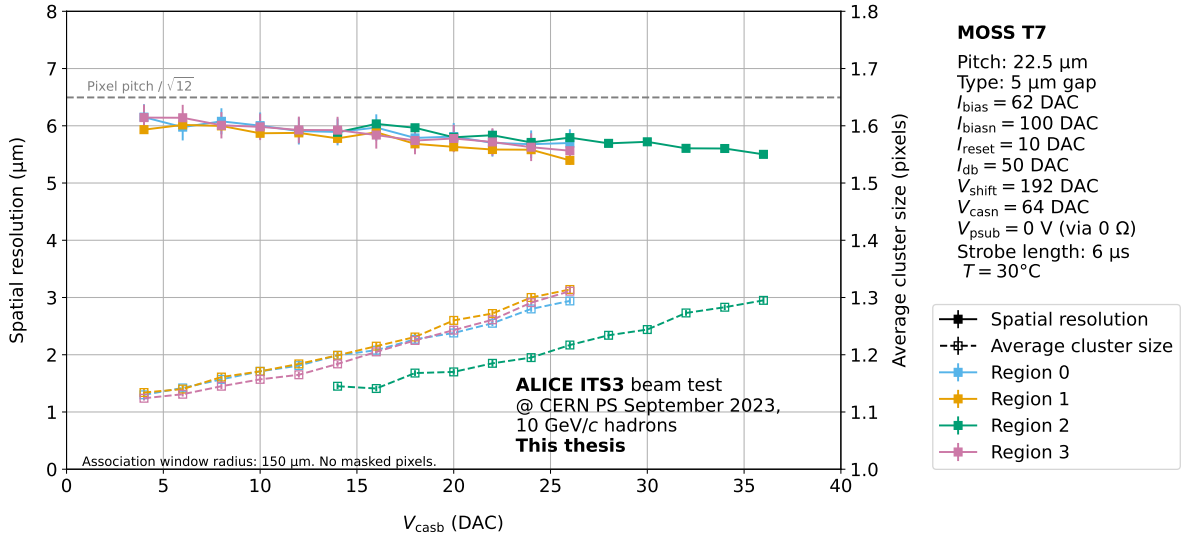


Figure 9.15: Mean spatial resolution (solid lines) and average cluster size (dashed lines) vs. V_{casb} for the RSU T7 of the MOSS chip. The information about the chip operation setting is listed in the right-hand side legend. Zero reverse substrate voltage was applied. The dashed gray line represents the binary resolution limit.

At this stage, no firm conclusions can be drawn since it was the second test beam with the same MOSS chip tested in the first test beam. Further studies are needed. However, the obtained results provide valuable information for subsequent test beams and for chip designers developing the final sensor.

9.3 Conclusions

The future upgrade of the ITS detector will replace the Inner Barrel layers of the ITS2 with truly cylindrical wafer-scale Monolithic Active Pixel Sensors fabricated in the 65 nm CMOS process utilizing the stitching technology. The possibility of operating bent silicon sensors without losing their performance was confirmed in the tests with bent ALPIDE chips [226]. The new 65 nm sensor prototypes from the first round of development demonstrated sufficient radiation hardness to meet the requirements of the ITS3 project. Currently, the R&D activities focus on characterizing the stitched sensors, designing the final sensors, and improving the detector mechanics and air cooling system. The project status of ITS3 was reviewed by LHCC, and the corresponding Technical Design Report was approved in 2024 [225].

Appendix A

Closure test: folding

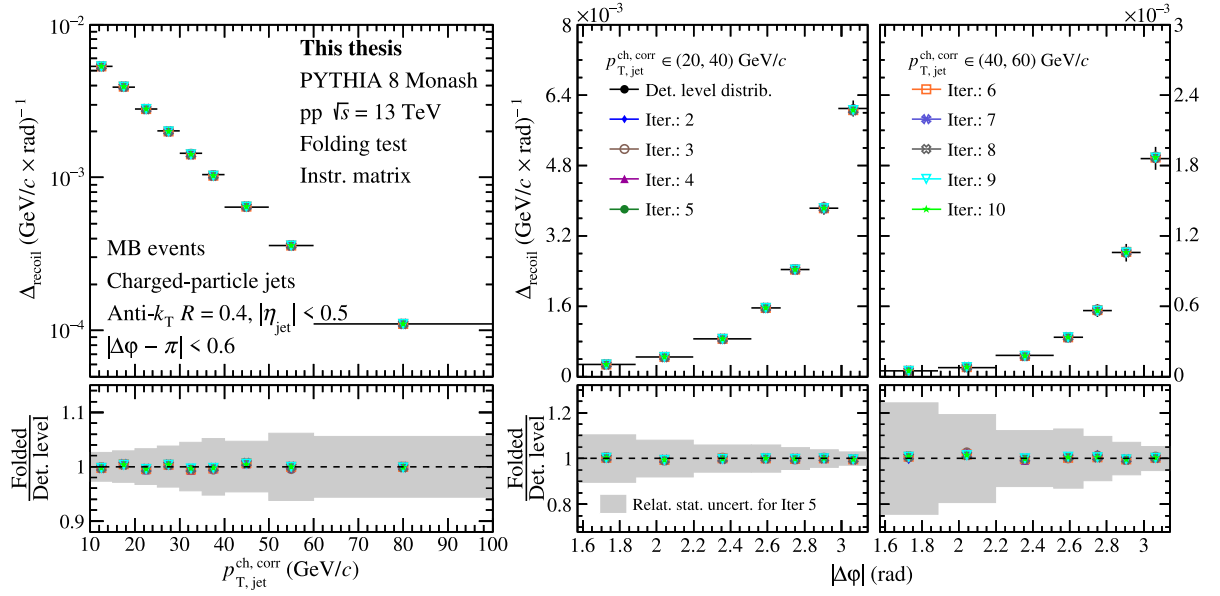
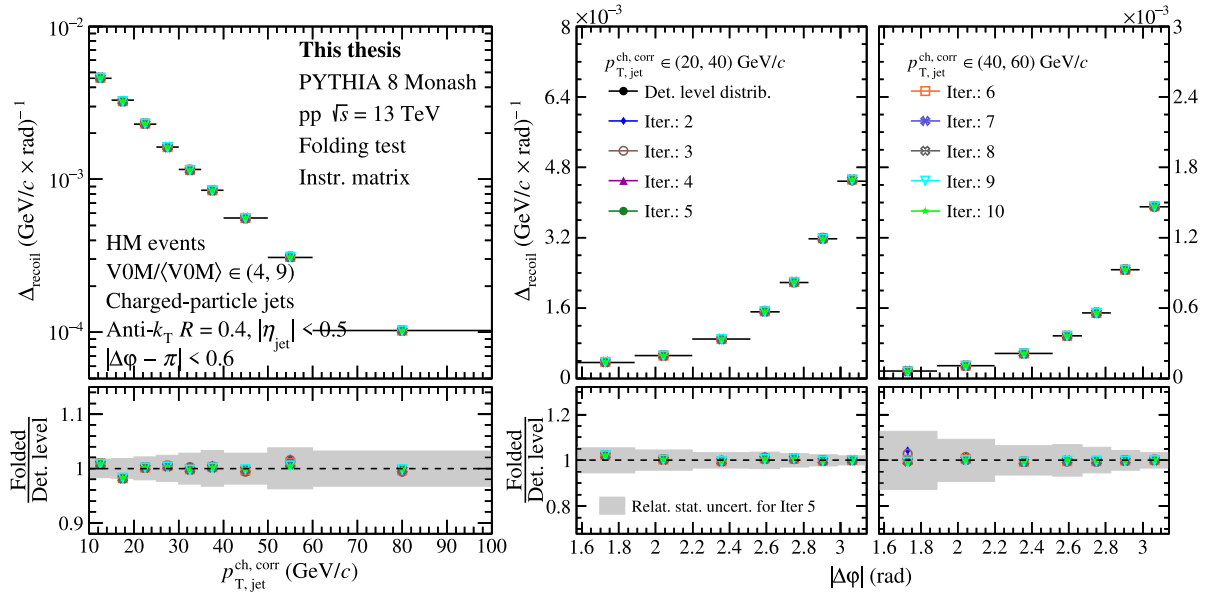


Figure A.1: Folding of the unfolded Δ_{recoil} distributions in Fig. 5.14.

Figure A.2: Folding of the unfolded Δ_{recoil} distributions in Fig. 5.15.

Appendix B

Data folding

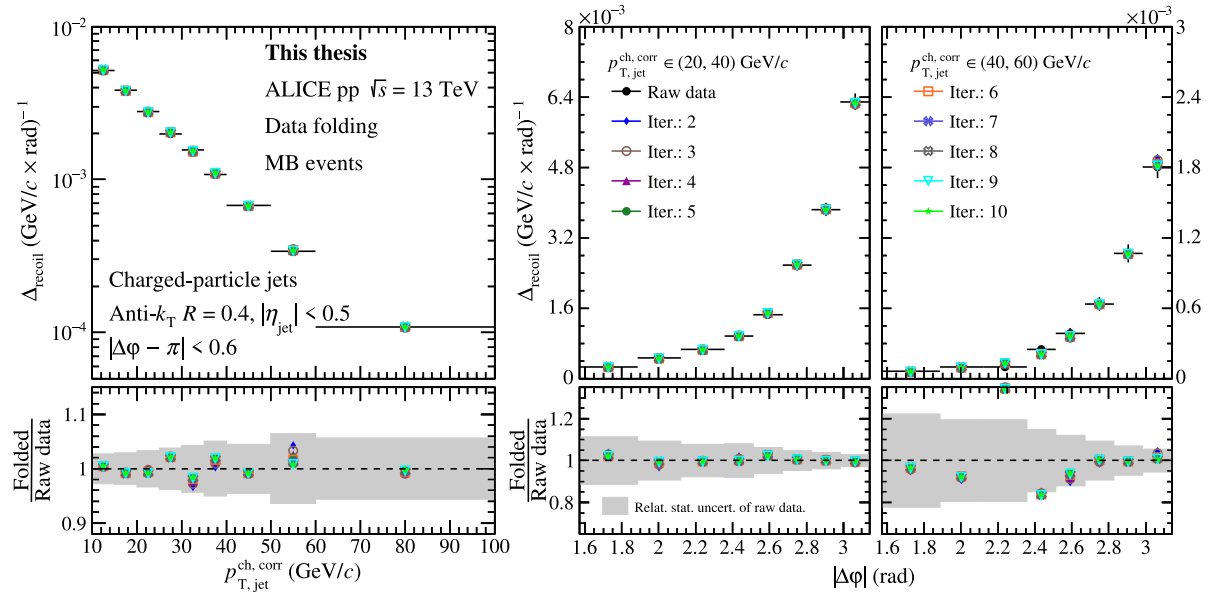
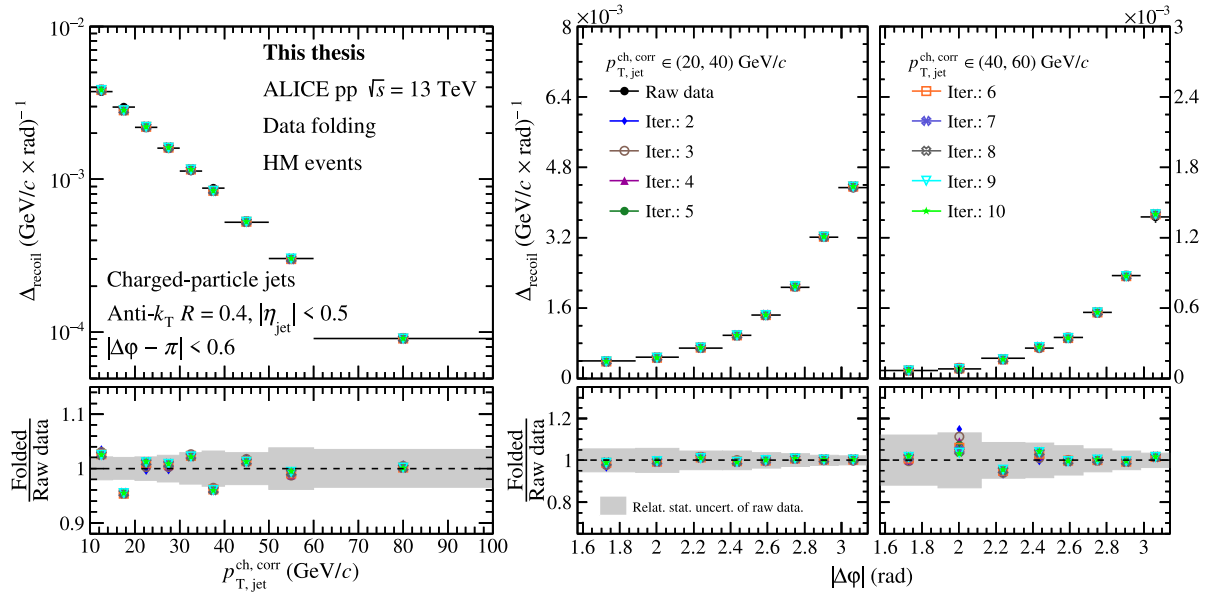


Figure B.1: Folding of the corrected Δ_{recoil} distributions in Fig. 5.21.

Figure B.2: Folding of the corrected Δ_{recoil} distributions in Fig. 5.22.

Appendix C

MAMI test beam with APTS sensor

This section contains the results of measuring the spatial resolution for the AF20P chip as a function of the charge threshold for different values of the reverse substrate voltage V_{sub} . Also, the comparison to the SPS results at $V_{\text{sub}} = 1.2$ V is shown. More details can be found in Sect. 9.2.4.1.

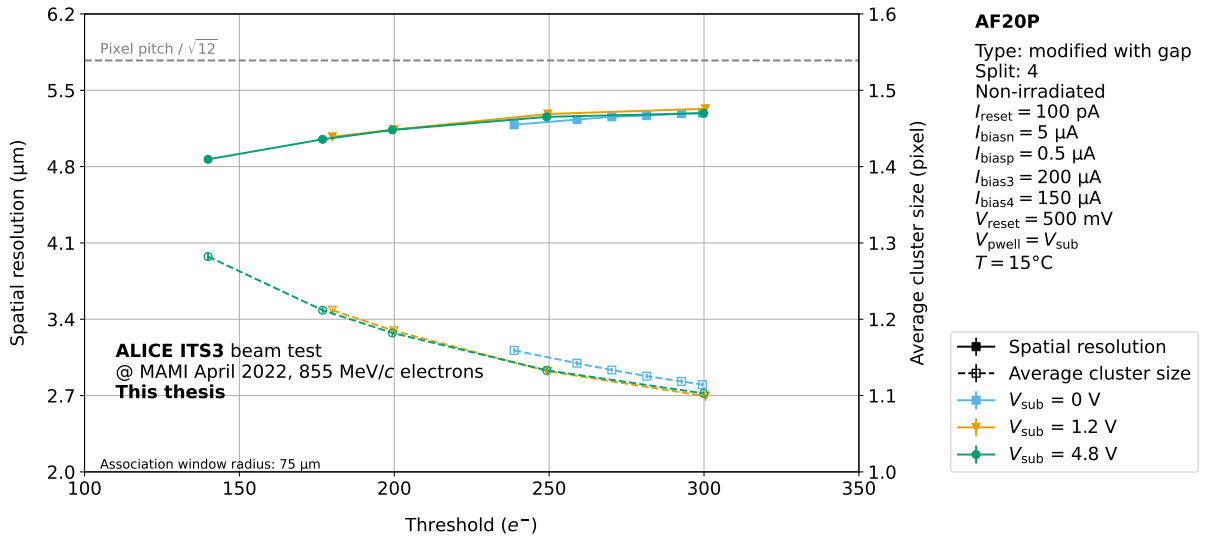


Figure C.1: Mean spatial resolution (solid lines) and average cluster size (dashed lines) vs. charge threshold at different V_{sub} . The information about the chip and the operation setting are listed in the right-hand side legend. The dashed gray line represents the binary resolution limit; see Sect. 9.2.4.1 for further details.

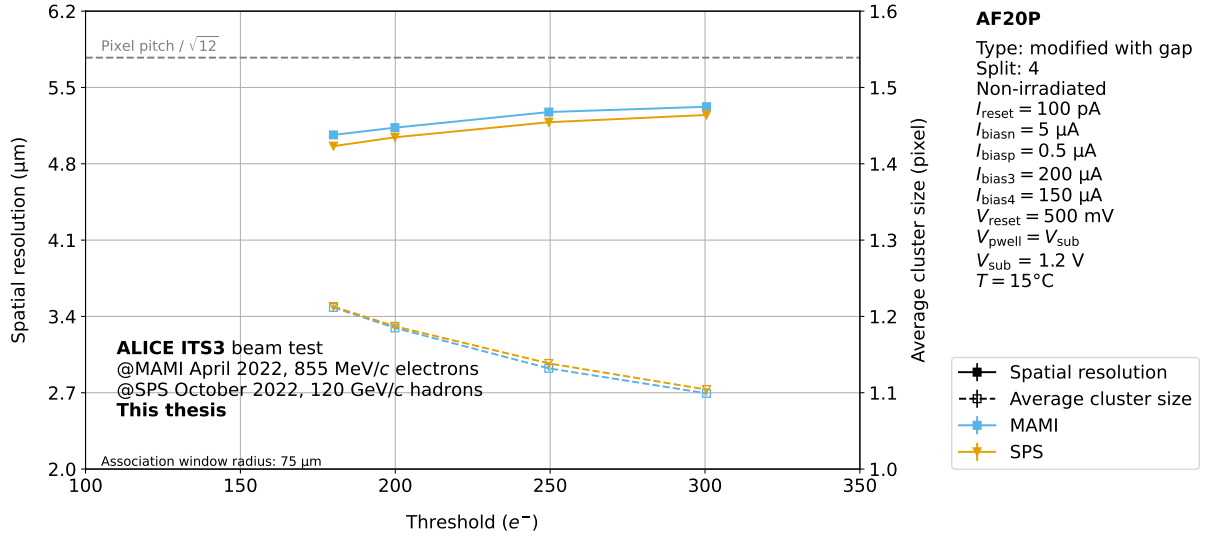


Figure C.2: Mean spatial resolution (solid lines) and average cluster size (dashed lines) vs. charge threshold at $V_{\text{sub}} = 1.2 \text{ V}$ measured in the MAMI and SPS test beams. The information about the chip and the operation setting are listed on the right-hand side legend. The dashed gray line represents the binary resolution limit.

Appendix D

PS test beam with MOSS sensor

Figure D.1 shows the dependence of the hit detection efficiency of the four regions of the RSU B4 on V_{casb} . No fake-hit rate measurements were available for the RSU B4 during the PS test beam. More details can be found in Sect. 9.2.4.2.

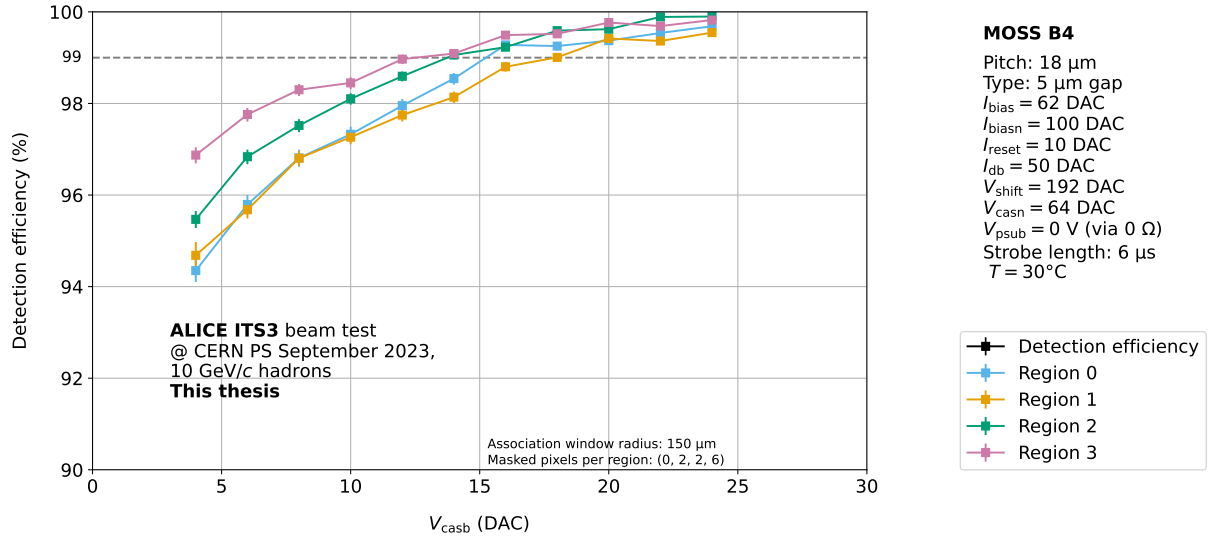


Figure D.1: Hit detection efficiency vs. V_{casb} for the RSU B4 of the MOSS chip. The information about the chip operation setting is listed in the right-hand side legend. Zero reverse substrate voltage was applied. The dashed line represents 99% detection efficiency limit required by the ITS3 project.

Figure D.2 shows the mean spatial resolution measured as a function of V_{casb} . The right y axis shows the average cluster size. The outlier for the region 3 arises from the noise contamination.

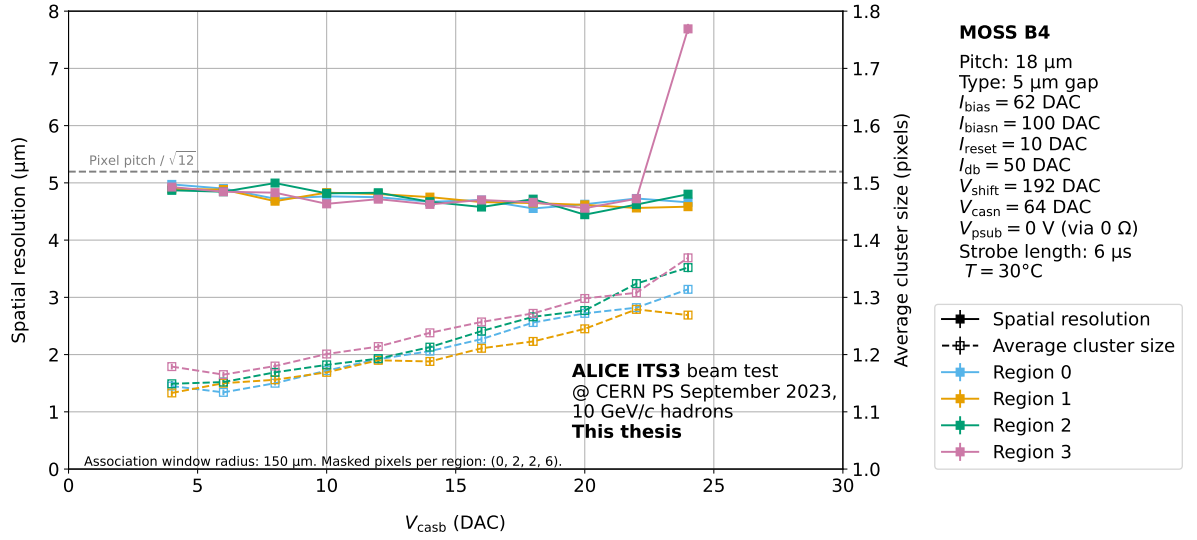


Figure D.2: Mean spatial resolution (solid lines) and average cluster size (dashed lines) vs. V_{casb} for the RSU B4 of the MOSS chip. The information about the chip operation setting is listed in the right-hand side legend. Zero reverse substrate voltage was applied. The dashed gray line represents the binary resolution limit.

Appendix E

List of publications

This section provides an overview of my publications related to the topics discussed in this thesis:

- ALICE Collaboration, *Search for jet quenching effects in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV via di-jet acoplanarity*, arXiv:2309.03788v1 [hep-ex]; **Member of the Paper Committee (among 3 contributors)**, submitted to JHEP journal.
- A. Kotliarov, *ALICE ITS3: the first truly cylindrical inner tracker*, EPS-HEP 2023, Published: January 19, 2024; DOI: <https://doi.org/10.22323/1.449.0538>.
- A. Kotliarov, *Searching for jet quenching effect using high-multiplicity inclusive jet and h+jet semi-inclusive jet in pp collisions with ALICE*, ICHEP 2022; Published: June 15, 2023; DOI: <https://doi.org/10.22323/1.414.0459>.
- A. Kotliarov, *Measurements of jet quenching via hadron+jet correlations in Pb–Pb and high-particle multiplicity pp collisions with ALICE*, PANIC 2021; Published: May 24, 2022; DOI: <https://doi.org/10.22323/1.380.0227>.
- A. Kotliarov, *Measurements of jet quenching via hadron+jet correlations in Pb–Pb and high-particle multiplicity pp collisions with ALICE*, EPS-HEP 2021; Published: February 17, 2022; DOI: <https://doi.org/10.22323/1.398.0293>.
- A. Kotliarov, *Search for jet quenching effects in high-multiplicity proton-proton collisions at $\sqrt{s} = 13$ TeV*, LHCP 2021; Published: November 17, 2021; DOI: <https://doi.org/10.22323/1.397.0183>.

Appendix F

List of conferences

This section provides an overview of my talks and posters related to the topics discussed in this thesis:

- The European Physical Society Conference on High Energy Physics (EPS-HEP 2023), **talk**: “ALICE ITS3: the first truly cylindrical inner tracker”, Hamburg, Germany, 20–25 Aug. 2023.
- The 41st International Conference on High Energy Physics (ICHEP 2022), **talk**: “Searching for jet quenching effect using high-multiplicity inclusive jet and h+jet semi-inclusive jet in pp collisions with ALICE”, Bologna, Italy, 6–13 July 2022.
- The XXIXth International Conference on Ultra-relativistic Nucleus-Nucleus Collisions (Quark Matter 2022), **poster**: “Search for jet quenching effects in high-multiplicity proton-proton collisions at $\sqrt{s} = 13$ TeV with ALICE”, Krakow, Poland, 4–10 April 2022.
- The Particles and Nuclei International Conference 2021 (PANIC2021), **talk**: “Measurements of jet quenching via hadron+jet correlations in Pb-Pb and high-particle multiplicity pp collisions with ALICE”, Online format, 5–10 Sept. 2021.
- The European Physical Society Conference on High Energy Physics (EPS-HEP 2021), **talk**: “Measurements of jet quenching via hadron+jet correlations in Pb-Pb and high-particle multiplicity pp collisions with ALICE”, Online format, 26–30 July 2021.
- 50th International Symposium on Multiparticle Dynamics (ISMD2021), **talk**: “Measurements of jet quenching via hadron+jet correlations in Pb-Pb and high-particle multiplicity pp collisions with ALICE”, Online format, 12–16 July 2021.

- 9th Annual Large Hadron Collider Physics (LHCP2021), **poster**: “Search for jet quenching effects in high-multiplicity proton-proton collisions at $\sqrt{s} = 13$ TeV”, Online format, 7–12 May 2021.
- The VI-th International Conference on the Initial Stages of High-Energy Nuclear Collisions, **poster**: “Search for jet quenching effects in high-multiplicity proton-proton collisions at $\sqrt{s} = 13$ TeV”, Online format, 10–15 Jan. 2021.
- ZIMÁNYI SCHOOL 2020, **talk**: “Search for jet quenching effects in high-multiplicity proton-proton collisions at $\sqrt{s} = 13$ TeV”, Online format, 7–11 Dec. 2020.

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