

박사 학위논문

LHC에서 ALICE를 이용한 $f_0(980)$ 의 특징의 이해

2023년 8월 22일

전 북 대 학 교 대 학 원

과학교육학과

김 준 이



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Understanding the nature of $f_0(980)$ with ALICE at the
LHC

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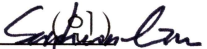
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ABSTRACT

Understanding the nature of $f_0(980)$ with ALICE at the LHC

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The $f_0(980)$ resonance had been observed years ago in $\pi\pi$ scattering experiments and is expected to be one of the scalar mesons. Since its first observation in the 1970s, the nature of light scalar mesons is far from the understanding, and no consensus on its internal structure has been reached, raising different suggestions regarding the structure of the $f_0(980)$, such as tetraquark, mesonic molecule, and conventional diquark structure.

The extreme environment, such as high temperature and density, deconfines quarks and gluons, freeing them and forming the Quark-Gluon Plasma (QGP), which is expected to be created immediately after the Big Bang. Such an early stage of our universe can be reproduced in relativistic heavy ion collisions, providing chances to study the properties of QGP. Extensive studies have been conducted on the geometry, evolution, and particle production of QGP and its modification, with many scientific acceptances for the existence of QGP. Surprisingly, phenomena exhibiting the existence of QGP are also observed in high-multiplicity proton-proton and proton-ion collisions, leading to further questioning QGP formation.

In this respect, the $f_0(980)$ resonance is measured with the ALICE detector via the $f_0(980) \rightarrow \pi^+\pi^-$ decay channel in relativistic nucleus-nucleus collisions. The present thesis describes the entire methodology to measure the invariant yield in different multiplicity classes at midrapidity. The particle yield ratios of the $f_0(980)$ resonance are measured to discuss the properties of the late hadronic phase throughout different collision systems and to explore the internal structure of the $f_0(980)$. Furthermore, the measurement of the nuclear modification factor and the elliptic flow and model predictions with different assumptions for

the $f_0(980)$ strengthen the physics messages in the present thesis.

Keywords : ALICE, resonance, $f_0(980)$, scalar meson

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I am obliged to the KoALICE (Korean ALICE) group for hugely providing me with a multitude of invaluable opportunities. The KoALICE has actively encouraged and supported my participation in various international conferences and research projects during the Ph.D. course. These experiences have enhanced my academic knowledge and practical communication skills with people around the globe.

Finally, best wishes to everyone around me for their academic achievement, health, and happiness.

Research activities

The present thesis consists of the following scientific articles and ongoing research activities, as well as introductory information. The research activities the author has achieved are summarized, and the details of contributions can be found throughout the thesis.

1. ALICE Collaboration, S. Acharya, et al., “Long- and short-range correlations and their event-scale dependence in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV”, JHEP 05 (2021) 290.

The author contributed from obtaining results to writing the paper as a paper committee chair, especially overcoming the limited acceptance of ALICE TPC and suggesting a new event classification using hard probes in pp collisions.

2. J. Kim, E.-J. Kim, S. Ji, and S. Lim, “Exploring the string shoving model in Pythia8 for collective behaviors in pp collisions”, J. Korean Phys. Soc. 79 no. 5, (2021).

The author largely contributed from obtaining results to writing the paper as a first author to further understand the string shoving model as a posterior study of item [1](#).

3. J. Kim, E.-J. Kim, Y. J. Kim, J. K. Ahn, and G. Y. Lim, “Simulation of angular resolution of a new electromagnetic sampling calorimeter”, Nucl.Instrum.Meth.A 1052 (2023) 168261.

The author contributed from obtaining results to writing the paper as a first author, especially developing the machine learning algorithm and systematic study to optimize the angular reconstruction and quantitatively understand the performance of the machine learning.

4. J. Kim, J. Seo, B. Hong, J. Hong, E.-J. Kim, Y. Kim, M. Kweon, S. H. Lee, S. Lim, and J. Park, “Model study on $\Upsilon(nS)$ modification in small collision systems”, Phys.Rev.C 107 (2023) 5, 054905.

The author contributed to the analysis and review processes as a main author, especially calculating the hydrodynamic evolution of p-Pb, pO, and OO collisions, systematically studying how param-

ters affect the result, and evaluating the nuclear modification factor from simulation with different setups.

5. Multiplicity and event-scale dependent flow and jet fragmentation in pp collisions at $\sqrt{s} = 13$ TeV and in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV,

<https://alice-publications.web.cern.ch/node/8859>

The author contributed to the analysis and obtaining the results as a paper committee member, especially proving the validity of the flow extraction method and providing the multiplicity and event scale dependent flow coefficients in small systems to understand the origin of the collectivity as a posterior study of item 1. The result is approved in internal paper proposals to be published in scientific journal.

6. Observation of abnormal suppression of $f_0(980)$ production in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV,

<https://alice-publications.web.cern.ch/node/8282>

The author contributed from obtaining results to writing the paper as a paper committee chair, where the full analysis process is described in the present thesis. The result is approved in internal paper proposals to be published in scientific journal.

7. Luminosity measurement in Run 3,

<https://indico.cern.ch/event/1207057>

The author contributed to the development of software that analyzes the vdM scan data for Length-scale calibration and non-factorization in Run 3 as a service work during the Ph. D course.

Chapter 1

Quark-Gluon Plasma and heavy ion collisions

The universe we are living in is created in an event known as Big Bang, around 13.8 billion years ago. Immediately after the Big Bang, around a timescale of $\mathcal{O}(10^{-6}\text{s})$, the universe is composed of extremely hot and dense plasma, in which no nucleons can exist in it. The plasma consists of free quarks and gluons and is called Quark-Gluon Plasma (QGP). QGP is an extreme matter possessing very high temperature and density that allow hadrons to deconfine, freeing quarks and gluons. Quantum ChromoDynamics (QCD), which describes the strong interaction, predicts the existence of QGP in an environment of high temperature and high density. Extensive researches on QGP are expected to lead to a better understanding of the creation of the universe and its evolution at the early stage of the universe [39]. Such an extreme environment is produced in laboratories operating high-energy particle accelerators by colliding with heavy ions having relativistic velocity, the so-called Little Bang [40].

The existence of QGP had been first widely accepted by the majority of scientists by measuring the anisotropic particle distributions conducted at Relativistic Heavy-Ion Collisions (RHIC) particle accelerator

in Brookhaven National Laboratory (BNL) [41] a few tens years ago. Among several experiments that precede the RHIC, the vast and significant researches have been conducted at the Large Hadron Collider (LHC) [42] at CERN, which presently operates the most powerful particle accelerator.

1.1 Quantum Chromodynamics

QCD, the most reputable quantum field theory describing the strong interaction, predicts nuclear matter that interacts with itself by the strong interaction. The elementary particles, quarks and gluons, compose the QGP, exhibit their color charges, and interact with the strong interaction by exchanging the color charges, which are mediated by the gluon. In comparison with the electromagnetic interaction, gluons, as a mediator of the strong interaction, can also interact by the strong interaction because gluons also exhibit color charges, unlike photons, having no electric charge, hence, not interacting by electromagnetic interaction.

The strength of the strong interaction is expressed as its coupling constant, α_s , which highly depends on the momentum transfer, Q^2 . The generic dependency of the α_s on the Q^2 is expected by the so-called renormalization group equation [43], in which its first-order perturbation leads to

$$\alpha_s(Q^2) = \frac{1}{B \ln(Q^2/\lambda^2)}, \quad (1.1)$$

where the $B = \frac{1}{12\pi}(11n_b - 4n_f)$ is a quantity related with the number of degrees of freedom for bosons (n_b) and fermions (n_f) and $21/(12\pi)$ in QCD, and the $\lambda \approx 200$ MeV is the QCD scaling parameter. The α_s decreases with increasing Q^2 , which is the behavior due to the self-coupling of gluons. At high $Q^2 \gg \lambda^2$, α_s correspondingly gets smaller, and quarks behave as free particles. Such a phenomenon is denoted as asymptotic freedom.

Another key feature of the gluon, unlike electroweak bosons, is no room to be experimentally observed as a free particle because of the non-neutral color charge, which is known as color confinement [44]. Instead,

quarks and gluons form color-neutral states, hadron which we can observe, by connecting quarks with gluons. In the system where quark and antiquark are connected with a gluon flux when the quark and the antiquark try to be separated, the formation of a new quark and antiquark pair is energetically favorable compared with the elongation of the flux, which is also derived phenomenon due to the self-interactions of gluons in $Q^2 \approx \lambda^2$ regime.

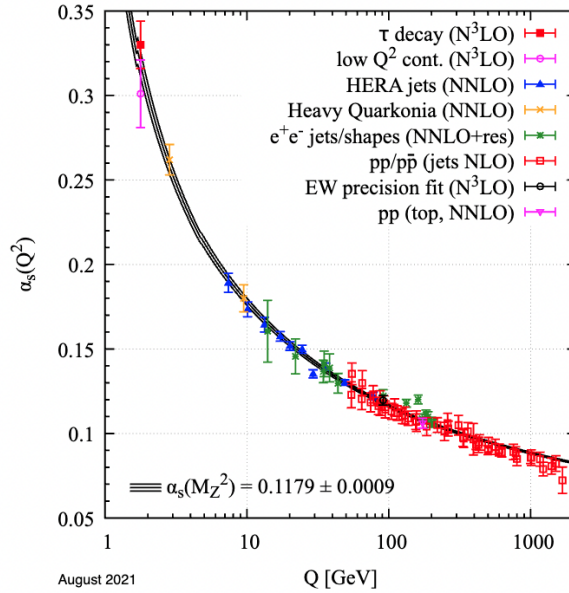


Figure 1.1: Summary plot of measurement for α_s as a function of Q [1]

Figure 1.1 shows the α_s as a function of Q measured by many different methods. As expected, the α_s decreases with increasing Q .

The effective potential considering the phenomenon [45], as well as the potential from the single color exchange, can be phenomenologically expressed as

$$V(r) \approx -\frac{\alpha_s}{r} + \kappa r, \quad (1.2)$$

where κ and r denote the string tension constant (≈ 1 GeV/fm) [45] and the distance between quarks, respectively.

1.2 Relativistic heavy ion collisions

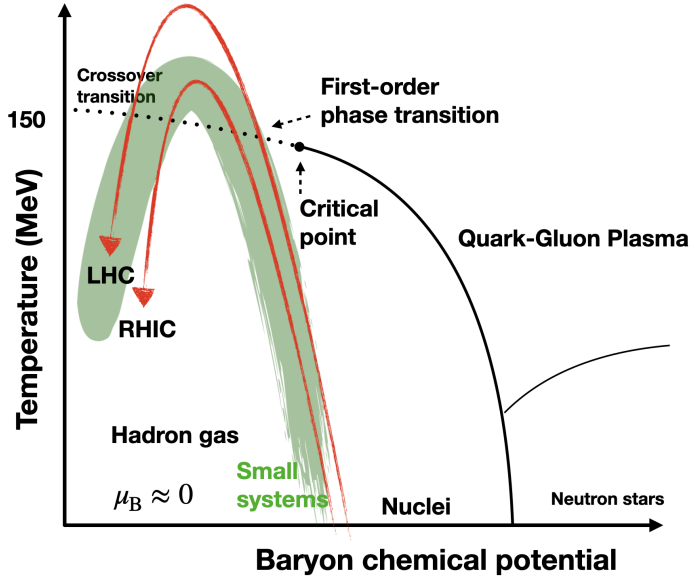


Figure 1.2: Phase diagram of QCD [2]

Figure 1.2 shows the phase diagram of QCD. The baryon chemical potential corresponds to the net-baryon density or imbalance between quarks and antiquarks. The collision of relativistic heavy ions is represented as red arrows for different collision energies. Ordinary matters lie in non-zero μ_B and low temperature, represented as nuclei. QCD predicts the state of QGP at high temperature for vanishing baryon chemical potential, and such an environment is produced by colliding with heavy ions accelerated up to relativistic velocity. The acceleration and collision of heavy ions are represented as drastically increasing temperature and decreasing baryon chemical potential. Immediately after the collision, extreme thermal energy, thanks to the high temperature, breaks the nuclei and nucleons passing over the first-order phase transition to the state of QGP. The volume of QGP expands owing to the pressure gradient inside QGP, starting to cool down with decreasing temperature. With further decreasing temperature, quarks and gluons can not exist as free particles, forming hadrons, and hadrons are allowed to only interact with

(pseudo-)elastic scattering, passing over the crossover transition to the hadron gas. Such edges of stages are called freeze-out.

The wide green line represents a possible scenario for the formation of QGP droplets in small collision systems, where the formation is not conclusive yet. The detailed description of the measurements evidencing the formation of QGP droplets will be addressed in chapter 2.

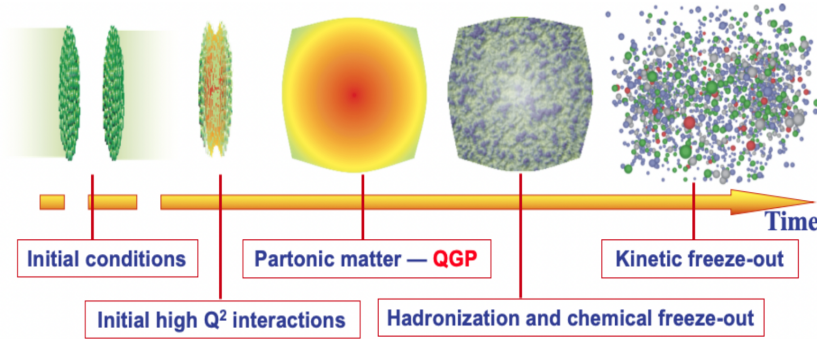


Figure 1.3: History of heavy ion collisions [3]

A history of heavy ion collisions is naively shown in figure 1.2, and the alternative representation can be found in figure 1.3 in terms of the evolution time. The history of heavy ion collisions can be divided into five stages. Due to the Lorentz contraction, accelerated ultrarelativistic heavy ions are contracted like a pancake. The impact parameter of heavy ion collisions determines which nucleons participate in collisions, which is called participant, and the other does not, which is called spectator. At the early stage of collisions, partonic hard scattering produces hard probes such as high transverse momentum particles, jets, and heavy-flavored particles. The produced QGP expands with the pressure gradient, where partons can traverse freely inside the plasma. As the temperature cools down, deconfined partons start to compose hadrons, and finally, hadrons also cease to interact with each other.

1.2.1 Initial stage of collisions

The initial collision geometry can be determined with the impact parameter of the collision, which can not be experimentally measured.

However, information on the initial geometry, such as impact parameter, can be modeled. One of the successful models is the MC-Glauber model [46] that is widely utilized for not only heavy ion collisions but also smaller collision systems [18, 21, 30, 31]. The model considers the distribution of nucleons in nuclei, the geometry of each participating nucleon, and the nuclear thickness.

The spatial distribution of the nucleons in the nucleus can be constructed by sampling a nuclear density distribution. Woods-Saxon potential [47], which is one of the common potentials of the nucleus and proportional to the nuclear density distribution, can be expressed as

$$V(r)/V_0 = -\frac{1}{1 + \exp[(r - R)/\alpha]}, \quad (1.3)$$

where α denotes the thickness of the nucleus surface. The position of nucleons can be determined with the probability distribution, which is proportional to $r^2 V(r)$, considering that the position of the nucleon is independent and the distances between nucleons are ≈ 1 fm.

The implementation of the geometry of each nucleon is important in the simulation because the probability of the interaction between two nucleons depends on how they are overlapped in the transverse plane. The nucleon density profile of each nucleon is given by the Fermi-distribution [48] and can be expressed as

$$\rho(r)_{\text{nuc}}/\rho_{\text{nuc},0} = \frac{1 + w(r/R)^2}{1 + \exp[(r - R)/\alpha]}, \quad (1.4)$$

where w is the width of the nucleon and can affect the overlapping area between nucleons.

Based on distributed nucleons and their internal density profiles, the nuclear thickness [46] is defined as the overlapping region of colliding two nuclear density profiles, where one profile can be expressed as

$$T = \int dz \rho = \sum_{\text{nuc}} \int dz \rho_{\text{nuc}}. \quad (1.5)$$

With this expression, the nuclear thickness using T_A and T_B can be represented as $T_{AB}(b) = \int T_A(s)T_B(s - b)d^2s$, where b denotes impact parameter. The T_{AB} is finally converted into the energy density with the scale factor that helps simulation to reproduce experimental data, such as charged-particle multiplicity.

1.2.2 Evolution of QGP

The pressure gradient of QGP lets it expand its volume and cool down. The observation of the near-perfect fluidity of a deconfined matter strongly favors the formation of QGP. The fluidity is nicely reproduced by the relativistic hydrodynamics [49]. The conservation of energy-momentum different charges can be expressed as

$$\delta_\mu T^{\mu\nu}(x) = 0 \quad (1.6)$$

$$\delta_\mu J_i^\mu(x) = 0 \quad (1.7)$$

where i in the continuity equation denotes index for each charge of baryon number, strangeness, electric charge, and so on to be conserved.

With the projector transverse to u defined as $\Delta^{\mu\nu} = g^{\mu\nu} - u^\mu u^\nu$, and without the heat current for simplicity, the energy-momentum tensor $T^{\mu\nu}$ can be expressed as

$$T^{\mu\nu} = \epsilon u^\mu u^\nu - (p + \Pi) \Delta^{\mu\nu} + \pi^{\mu\nu} \quad (1.8)$$

where ϵ and P denote the energy density and the equation of state, respectively. The shear and bulk pressures are denoted as $\pi^{\mu\nu}$ and Π , respectively, and key parameters for transport properties of the near-perfect relativistic fluid. The shear viscosity, η is a value of the resistance of the fluid against shear stress. Lower η induces stronger interactions between the constituents. The bulk viscosity, ζ is a measure of the resistance of the fluid to expansion and compression. Both parameters are essential to describe the evolution of QGP and extensively studied [50] by measuring differential spectra of final-state hadrons to determine their values to understand the characteristics of QGP better.

1.2.3 Freeze-out

There are two temperatures and corresponding timescales during the expansion and cooling down of QGP. One is the chemical freeze-out, representing a limit when all inelastic interactions, such as hadronization, among the constituents are expected to cease. The other is the kinetic freeze-out, representing a limit when all elastic interactions stop [51]. Understanding of two stages of freeze-outs is essential to study

particle production and its modification from small to large collision systems and is also important to interpret the results of the present thesis. Hence, they deserve to be in a separate section.

1.3 Chemical freeze-out

At the end of the hadronization of QGP, all inelastic interactions stop, admitting no existence of deconfined constituents. Therefore, the relative particle composition is expected to freeze at the chemical freeze-out.

The relative particle composition at the chemical freeze-out can be studied using the statistical thermal model [52] in nuclear–nuclear collisions. The statistical thermal model assumes that the hadronized QGP is in chemical and thermal equilibrium. The thermal particle composition is calculated by constructing the partition function. In the Grand Canonical Ensemble (GCE), the partition function for particle species i is expressed as

$$\ln Z^{\text{GCE}}(T, V, \mu_Q)_i = \frac{V g_i}{2\pi^2} \int \pm dp p^2 \ln [1 \pm \exp(-(E_i - \mu_i)/T)], \quad (1.9)$$

where T , V , Q , μ_Q , g_i , and E_i denote the temperature, the extractable volume, conserved charges, the chemical potential conserving the corresponding charge Q , the spin degeneracy of $(2J_i + 1)$, and the total energy, respectively. The plus (minus) sign is for fermions (bosons). For each particle species i , (net) baryon number (B), isospin (I_3), strangeness (S), electric charge (Q), and even charmness (C) are conserved at the chemical freeze-out and implemented in the total chemical potential as $\mu_i = \sum_j \mu_{Q_j} Q_j$, where j is index for different conserved charges. Based on the constructed partition function, the relative particle composition, or the particle density (n_i), can be derived from the partition function, which can be expressed as

$$n_i = N_i/V = -(T/V) \frac{\delta \ln Z_i}{\delta \mu} \quad (1.10)$$

The baryon chemical potential, the temperature, and the volume at the chemical freeze-out can be extracted by reproducing the particle composition measured by the experiments.

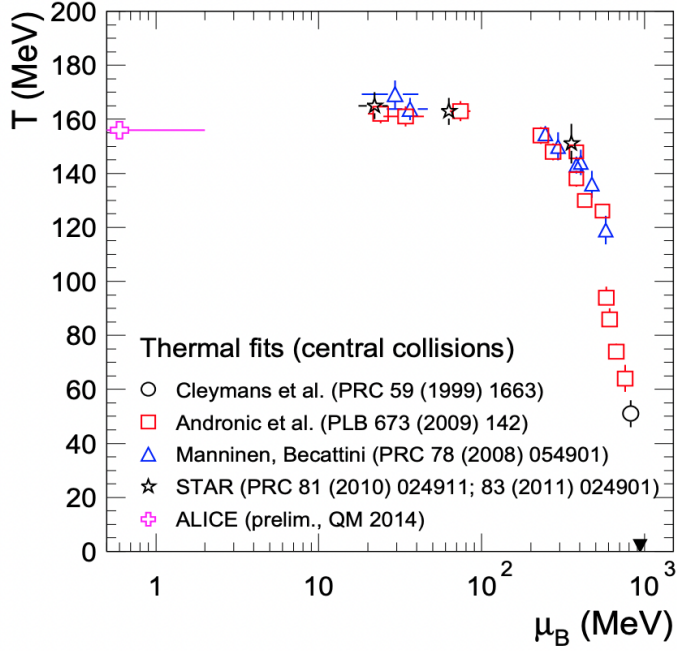


Figure 1.4: The phase diagram from the thermal fit [4]

The thermal fit can be used to examine the QCD phase diagram [4] by extracting the temperature and the baryon chemical potential as shown in figure 1.4. The chemical freeze-out temperature is saturated for $\mu_B < 300$ MeV, implying the hadronization temperature is the same as the chemical freeze-out temperature.

Figure 1.5 shows different hadron yields in central (0–10%) Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. The experimental data are reproduced with the canonical statistical model, extracting the temperature and the baryon chemical potential at the chemical freeze-out [5]. The setup of different temperature and baryon chemical potential is tested to see the sensitivity of the model and provides a worse expectation for the particle yields. The canonical statistical model with the extracted parameters nicely describes the particle yields except for K^{*0} and (anti-)proton. The K^{*0} yield is significantly modified after the chemical freeze-out through elastic scatterings, which will be further discussed in section 1.4. Protons coalesce with each other after the chemical freeze-out, resulting in the

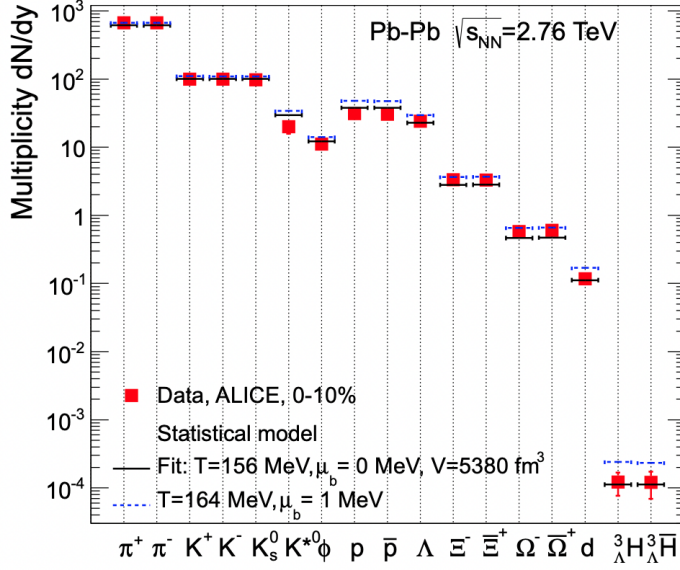


Figure 1.5: Hadron yields in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV and calculated particle yields using the statistical canonical model together with the extracted temperature and the chemical potential. Different temperatures and chemical potentials are used to test the sensitivity of the model [5]

overestimation of the yield by the model, and deuterons and Helium-3 are underestimated, in turn.

1.4 Kinetic freeze-out

After the chemical freeze-out, the density of hadrons is still high enough for them to interact with each other elastically, mostly modifying the momentum of hadrons. Due to its high density of hadrons, the stage between chemical and kinetic freeze-out is called hadron gas or hadronic phase. Modifications on p_T during the hadron gas stage are modeled with the blast wave model [53], which assumes that p_T spectra are by-products originating from the specific temperature. The model performs a global fit on p_T spectra for different particle species, extracting the kinetic freeze-out temperature (T_{kin}) and the average velocity of

the hadron gas expansion ($\langle\beta_T\rangle$). The fit range for p_T is limited to reject feed-down effects, which are dominant at low p_T , and products from hard scatterings, which are dominant at high p_T .

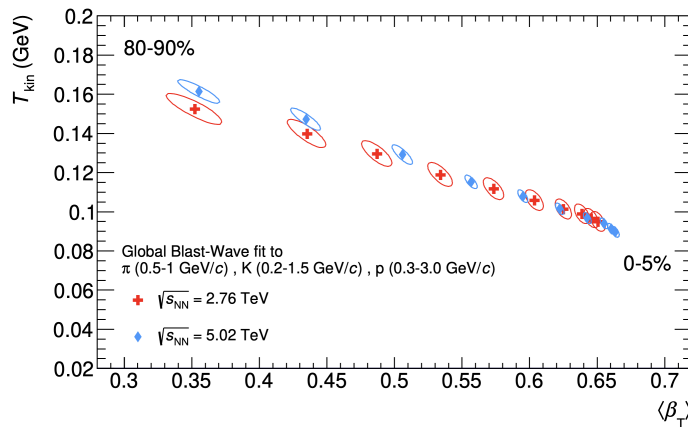


Figure 1.6: Correlations of expansion velocity and the kinetic freeze-out temperature extracted with the blast wave fits in peripheral to central Pb–Pb collisions [6]

Figure 1.6 shows the correlations of the expansion velocity and the kinetic freeze-out temperature in Pb–Pb collisions. The parameters are extracted from the blast wave model. The correlated uncertainties from the global fit are shown as ellipses. With the clear correlations between the temperature and the velocity for each data point, the temperature decreases with the increasing expansion velocity, indicating that the expansion is more rapid in central than peripheral collisions. The lower kinetic freeze-out temperature in central collisions is owing to its larger volume, allowing longer expansions and, correspondingly, a decrease in the temperature.

Moreover, the properties of the hadron gas can be further explored using resonance particles, which have the same constituent quarks as their ground state and a different quantum number, resulting in heavier mass than the mass of the ground state. Resonances are likely to have a very short lifetime, comparable with the duration time of the hadron gas as $\tau \approx 10 \text{ fm}/c$ [54]. Due to its comparability, resonances can be uniquely useful probes to study the properties of the hadron gas.

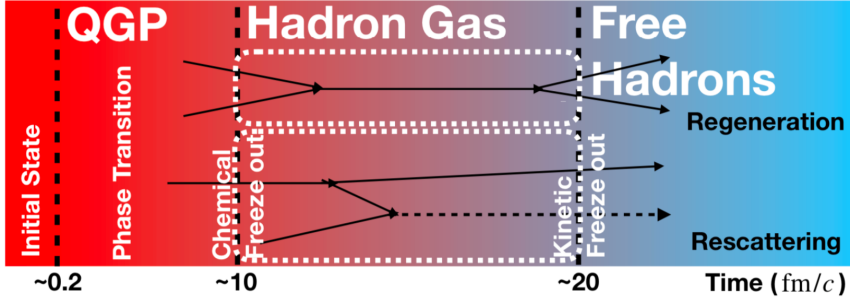


Figure 1.7: Resonances in heavy ion collisions

As hadrons can (pseudo-)elastically interact with the hadron gas, the decay products of resonances can also interact with them, modifying their momenta. And such modifications can significantly affect the measurement of the resonance yield via two processes of rescattering effects and regeneration, as shown in figure 1.7. The rescattering effects reduce the resonance yield by modifying the momenta of the decay products of the resonance, resulting in failure of the reconstruction of the invariant mass of the resonance. The regeneration effects enhance the resonance yield by combining the hadrons into a resonance. The cross sections for both processes are dependent on the scattering cross section between the decay product and the hadron gas [55]. The study of the properties of the hadron gas can be achieved by comparing the particle yields for different particles in different collision systems, where the modification can be observed by measuring the particle yield ratio or the nuclear modification factor.

Figure 1.8 shows the particle yield ratio of $K^*(892)^0$ and $\phi(1020)$ to charged kaons as a function of charged-particle multiplicity elevated to $1/3$ in different collision systems. The lifetime of $K^*(892)^0$ is measured to be ~ 4.2 fm/c [1], which is short enough for decay products to interact with the hadron gas. It is also worth mentioning that K^{*0} and the charged kaon include one strange quark, resulting in no significant modification of the ratio from the strangeness enhancement expected (see chapter 2). The ratio of K^{*0}/K decreases with increasing multiplicity throughout the entire system, indicating the sequential suppression of K^{*0} yields. On the other hand, the ratio of ϕ/K increases with increasing multiplicity.

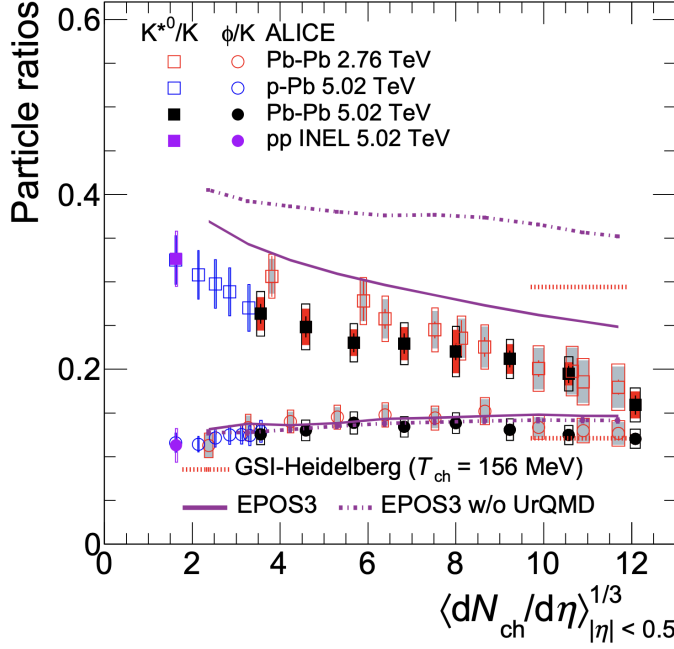


Figure 1.8: Particle yield ratio of $K^*(892)^0$ and $\phi(1020)$ to charged kaons as a function of charged-particle multiplicity elevated to 1/3 in different collision systems and predictions from EPOS LHC with and without UrQMD afterburner [7]

Because the lifetime of ϕ is 46.2 fm/c [1], the interaction with the hadron gas is weakly expected. On the other hand, ϕ mesons include one strange quark and one anti-strange quark, exhibiting stronger sensitivity on the strangeness enhancement.

Such trends are reproduced with the Ultra-relativistic-Quantum-Molecular-Dynamics approach (UrQMD) model [56], which simulates rescatterings and regenerations processes after the chemical freeze-out while initial particle composition is calculated with the EPOS LHC [57]. The UrQMD nicely describes the decreasing trend of K^{*0}/K and mildly increasing trend of ϕ/K at the same time, implying the dominant rescattering effects from small to large collision systems for K^{*0} resonances. Reversely, such trends are not reproduced without UrQMD.

Figure 1.9 shows the ratio of K^{*0}/K as a function of p_T in pp and cen-

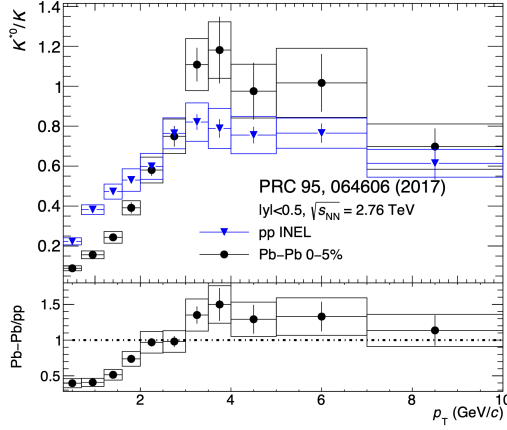


Figure 1.9: p_T -differential particle yield ratio of K^{*0}/K in pp and Pb-Pb collisions [8].

tral Pb-Pb collisions. The ratios increase at low p_T due to radial effects and saturate at high p_T . The ratios measured in pp and Pb-Pb collisions are consistent at high p_T region. However, they are getting discrepant with decreasing p_T , exhibiting the suppression of the ratio in central Pb-Pb collisions as shown in the double ratio of K^{*0}/K in Pb-Pb to pp collisions in the lower panel. From the suppression observed in figure 1.8, the rescattering effects are dominant at low p_T , and no suppression can be found at high p_T .

Figure 1.10 shows the ratio of K^{*0}/K as a function of p_T in high- and low-multiplicity pp collisions. The same trend measured in pp and Pb-Pb collisions can be observed in figure 1.9 when one assumes the correspondence of central Pb-Pb to high-multiplicity pp and minimum-bias pp to low-multiplicity pp events, suggesting the existence of rescattering effects in small collision systems.

Figure 1.11 shows the particle yield ratio of different hadrons to charged pions as a function of charged-particle multiplicity. Rather than other ratios, which are already discussed before, the ratio of K^{*0} to charged pions is worth discussing to understand the interplay of competing two different effects. The ratio of K^{*0}/π is consistent in pp and p-Pb collisions and starts to decrease in Pb-Pb collisions. Because K^{*0} includes one strange quark, the K^{*0} yield is enhanced due to the strangeness en-

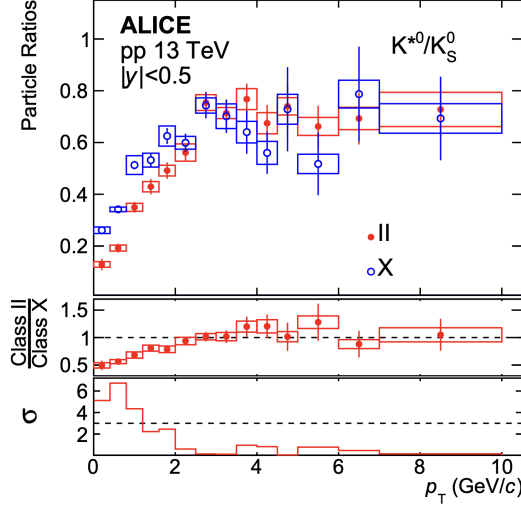


Figure 1.10: p_T -differential particle yield ratio of K^{*0}/K in high- and low-multiplicity pp collisions [9].

hancement. At the same time, K^{*0} has very short lifetime yielding the suppression of the K^{*0} yield. As a result of an interplay of two competing effects, the ratio is consistent in small collision systems. However, the ratio starts to decrease in heavy ion collisions because the strangeness enhancement is saturated (see section 2.2), and the remaining effect, rescattering effects only decrease the K^{*0} yield with increasing charged-particle multiplicity. Therefore, the ratio of K^{*0}/π decreases in large collision systems.

Such approaches can also be applied to different particles to better understand the properties of the hadron gas, and to explore if a given particle experiences the strangeness enhancement to confirm the inclusion of strange quarks for the particle ($f_0(980)$ resonance in the present thesis).

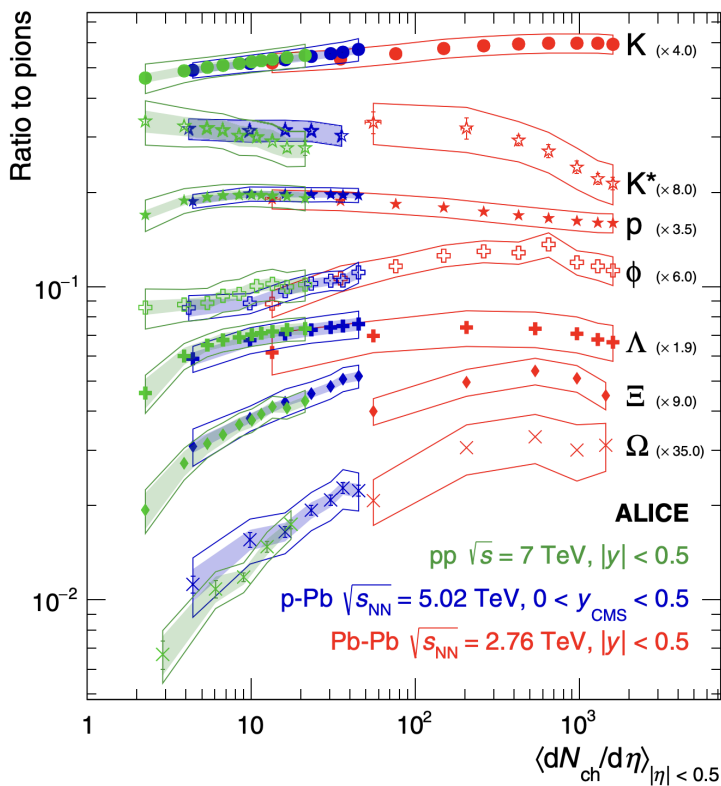


Figure 1.11: Particle yield ratio of different hadrons to charged pions as a function of charged-particle multiplicity [10]

1.5 Collectivity

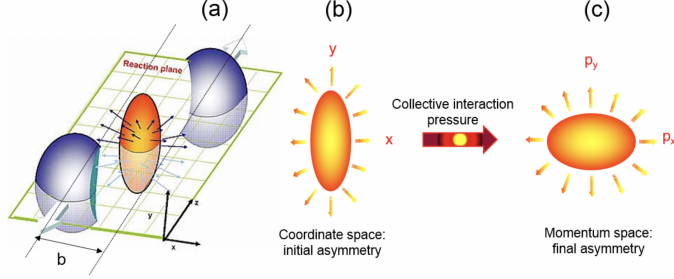


Figure 1.12: Initial collision geometry of semi-central Pb–Pb collisions and its evolution [11]

One of the most unambiguous evidence indicating the formation of the QGP in heavy ion collisions is the anisotropic particle distribution. Collectivity represents the behavior of produced particles exhibiting common properties of their source. In the expanding fireball, produced particles interact with the medium, exhibiting properties of the fireball until the kinetic freeze-out reaches. Due to the anisotropic initial collision geometry, produced particles experience interactions with the medium depending on their relative azimuths to the reaction plane, which determines how long the particle traverses the medium. The interplaying effects of the initial collision geometry and its expansion can be expressed as

$$v_n = k\varepsilon_n, \quad (1.11)$$

where ε_n , k , and v_n denote the initial collision geometry of being n -th figure (i.e., ε_2 is eccentricity), the evolution power of the fireball, and how many particles are collectively flowing in the final state, well known as flow coefficients with n -th mode, respectively. Figure 1.12 shows the initial anisotropic collision geometry (a) and its expansion (b), leading to anisotropic final-state particle distribution (c).

With the success of the reproduction of the initial state thanks to the Glauber model [46], the hydrodynamic calculation for describing the expansion of the fireball reproduces the measured v_2 in Pb–Pb collisions. Figure 1.13 shows the v_2 of charged hadrons as a function of centrality

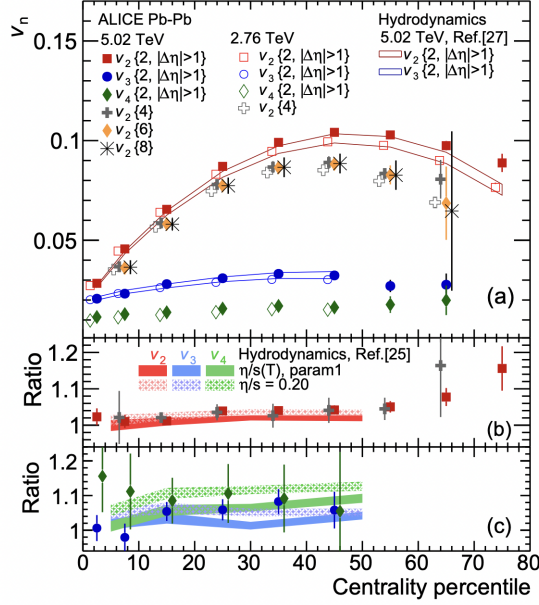


Figure 1.13: Flow coefficients as a function of centrality in Pb–Pb collisions. The measured data are reproduced using the AMPT model and hydrodynamic calculations. [12]

in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ and 5.02 TeV, which is reproduced by the hydrodynamic calculation [12]. The sizable v_2 can be seen in the most central collisions, where the ε_2 is expected to be very small, and v_2 is largest in semi-central collisions. Finally, the v_2 is again decreasing with decreasing centrality.

The above trend can be understood with the evolution power, $k = v_n/\varepsilon_n$. Figure 1.14 shows the v_2/ε_2 as a function of the charged-particle multiplicity inversely elevated to 1/3 [13]. The evolution power is expected to increase with increasing multiplicity. Because the evolution power is strongest in the most central collision, v_2 can be measurable with the very small ε_2 .

The measurement of flow coefficients can be performed by decomposing the azimuthal particle distribution using Fourier series with the reconstruction of the event plane (Ψ_n) [58], and the procedure can be

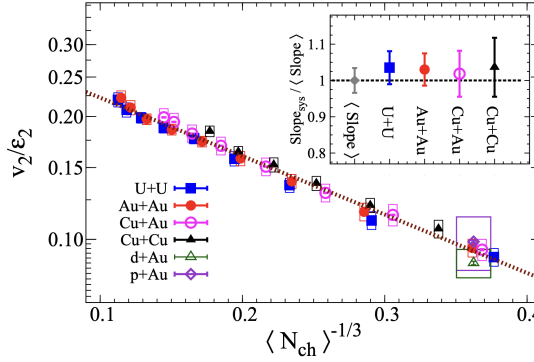


Figure 1.14: Evolution power (k) as a function of charged-particle multiplicity elevated to $-1/3$ [13]

expressed as

$$\frac{dN}{d\varphi} = G \left(1 + \sum_{n=1}^{\infty} 2v_n \cos[n(\varphi - \Psi_n)] \right) \quad (1.12)$$

The extracted v_n with the above equation is corrected for the resolution of the event plane and can be expressed as $v_n = v_n^{\text{obs}}/\mathcal{R}_{\Psi_n}$ [59]. On the other hand, correlating two particles of a “trigger” particle and an “associated” particle also allows one to measure the flow coefficients without reconstructing the event plane. Consider the convolution of equation 1.12 for the trigger and associated particles, which can be expressed as

$$\begin{aligned} \frac{dN^{\text{trig}}}{d\varphi} &= G^{\text{trig}} \left(1 + \sum_{n=1}^{\infty} 2v_n^{\text{trig}} \cos[n(\varphi^{\text{trig}} - \Psi_n)] \right) \\ \frac{dN^{\text{assoc}}}{d\varphi} &= G^{\text{assoc}} \left(1 + \sum_{n=1}^{\infty} 2v_n^{\text{assoc}} \cos[n(\varphi^{\text{assoc}} - \Psi_n)] \right) \\ \frac{1}{G^{\text{trig}}} \frac{dN^{\text{pair}}}{d\Delta\varphi} &= G^{\text{assoc}} \left(1 + \sum_{n=1}^{\infty} 2v_n^{\text{trig}} v_n^{\text{assoc}} \cos(n\Delta\varphi) \right), \end{aligned} \quad (1.13)$$

where the $G^{\text{trig(assoc)}}$ is the number of trigger (associated) particles and $\Delta\varphi = \varphi^{\text{trig}} - \varphi^{\text{assoc}}$ is relative azimuthal angle between trigger and associated particles. The correlations of trigger particles with associated

particles in different modulations are averaged out because of no correlations of Ψ_n and Ψ_m . The two-particle correlation method naturally constructs two-particle pairs originating from jets (see section 1.6), contaminating the measurement of v_n . Such a contamination can be reduced by taking large $\Delta\eta$ when one observes the flow coefficients because jet contributions are mostly concentrated on $(\Delta\eta, \Delta\varphi) \sim (0, 0)$.

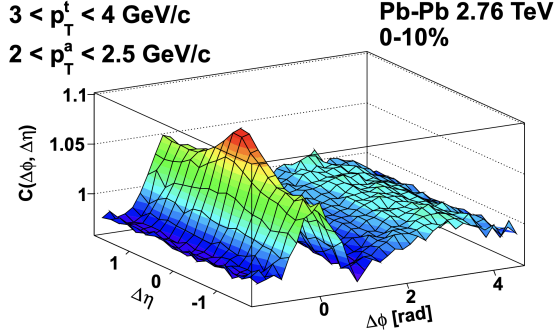


Figure 1.15: Two-particle angular correlations in Pb–Pb collisions [14]

Figure 1.15 shows the scatter plot of relative azimuthal angle and relative pseudorapidity between the trigger and associated particles in Pb–Pb collisions at $\sqrt{s_N} = 2.76$ TeV. The region of small and large $\Delta\varphi$ ($\Delta\eta$) is called near-side and away-side (short-range and long-range), respectively. From the correlation function, clear modulation can be found along $\Delta\varphi$ direction. On the short-range near-side region, there is a sharp peak, and this is coming from the correlations of particles in the same jet as mentioned above. Long-range near-side regions can be selected to reduce jet contributions to the anisotropy measurement. It is worth mentioning that the flat distribution at the away-side along $\Delta\eta$ is mostly attributed to the correlations of jets and back-to-back jets, the so-called η -swing.

1.6 Partonic energy loss

Products originating from hard scatterings of partons at the initial collision stage can be characterized by their high transverse momenta or large masses scale. Such hard probes traverse the QGP medium, fully ex-

perienicing the evolution of the medium and exhibiting the properties of the medium [60, 61]. Interactions of the hard probes with the medium lead to parton energy loss, which is transformed into fragmented products in the final state.

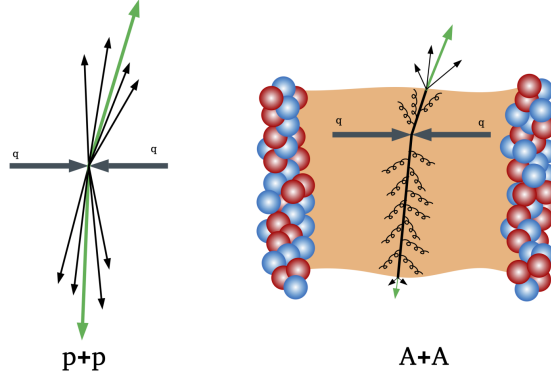


Figure 1.16: Production of jets in proton-proton (left) and heavy ion (right) collisions

One of the crucial observables to study the parton energy loss is the “jet” [62], which is a collimated spray of fragmented particles originating from the interaction between high-momentum initial parton and the medium. Jets may not be measured when such interactions are abundant, which is called jet quenching and was first proposed by Bjorken [63]. The measurement of jet quenching brings consensus on the existence of the QGP in heavy ion collisions. Considering the momentum conservation, the creation of a high-momentum parton should happen in both opposite directions (jet and back-to-back jet). Figure 1.16 shows a cartoon for the creation of two jets from initial hard scattering. The jet escaping on the upper side loses less energy due to a shorter journey in the medium, compared with the jet escaping on the lower side, while such effects are not expected in small systems.

The measurement of such an imbalance between two measured transverse momenta of two jets can indicate the jet quenching [15]. Furthermore, the relative azimuthal angle between two jets should be $\Delta\varphi \sim \pi$, which can fluctuate if there are significant interactions within the medium. Figure 1.17 shows the asymmetry of energies of two jets,

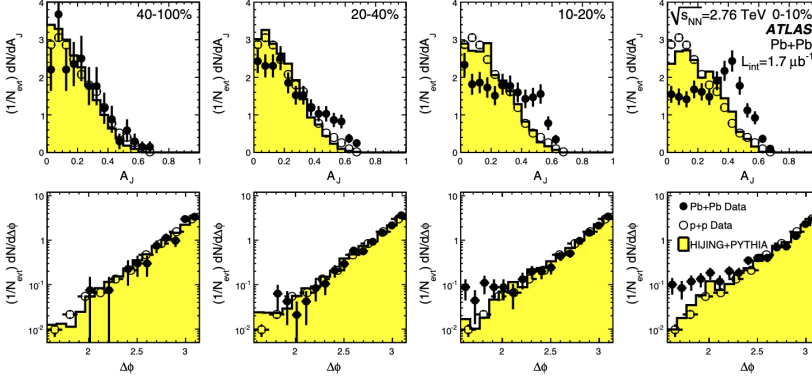


Figure 1.17: (Upper) A_J and (Lower) $\Delta\varphi$ distributions for different centrality classes in pp and Pb–Pb collisions and from the simulation. The agreement between results from pp collisions and the simulation can be seen. The distributions from Pb–Pb collisions show a different trend from pp collisions [15].

$A_J = (E_{T1} - E_{T2})/(E_{T1} + E_{T2})$ and relative azimuthal angle between two jets. The measured data exhibit significant asymmetry and smaller relative azimuthal angle, which are not expected from the model and not observed in results in pp collisions.

Another key measurement of the parton energy loss is the nuclear modification factor, which compares the particle productions in small and large collision systems per the number of binary collisions. In this context, the small (large) system corresponds to the pp (Pb–Pb) collisions. The number of binary collisions can be determined with the Glauber model in Pb–Pb collisions [46]. Because the hard probes are produced from initial hard scattering and characterized by their high transverse momenta, the particle production at high transverse momentum should be only proportional to the number of binary collisions between small and large collision systems. To quantitatively measure this, the nuclear modification factor can be defined as

$$R_{AA} = \frac{d^2N_{AA}/dp_T dy}{N_{\text{coll}} d^2N_{pp}/dp_T dy}, \quad (1.14)$$

where $d^2N/dp_T dy$ and N_{coll} denote the p_T spectrum of produced particles and the number of binary collisions, respectively. The smaller (larger)

R_{AA} than unity represents the suppression (enhancement) of the particle production in Pb–Pb collisions compared with the production in pp collisions.

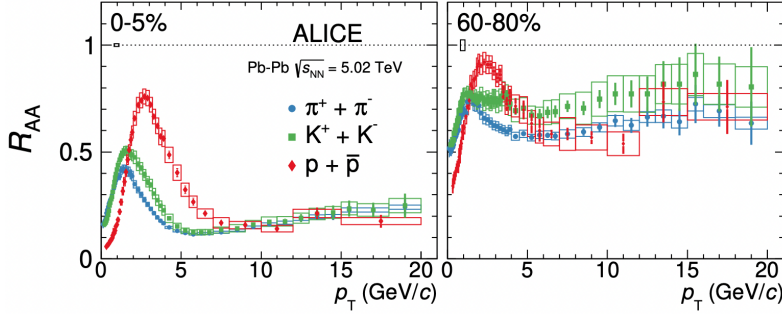


Figure 1.18: The nuclear modification factor as a function of transverse momentum in central and peripheral Pb–Pb collisions for different particles [6].

Figure 1.18 shows the R_{AA} of π , K , and p in central and peripheral Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV as a function of p_T . All particles shown in the figure exhibit a suppression at high p_T region in peripheral collisions, and the suppression gets significant in central collisions. The system-size dependent suppression of high- p_T particles can be explained by the parton energy loss due to interactions with the medium, QGP.

1.7 Suppression of heavy quarkonium

The production of heavy quarkonium, which is a meson composed of a heavy quark and its antiquark ($q\bar{q}$), is predicted to be suppressed due to the color screening effect [64]. The color screening is the Debye screening for strong interaction. The sequential suppression of quarkonium states for charmonium and bottomonium is one of the apparent evidence proving the formation of a deconfined matter in which the constituents of the medium exhibit color charges. The competition between the binding energy of the quarkonium and the potentials of the strong interaction associating with surrounding light quarks and gluons in a given temperature can be represented as the size of hadron and the Debye screening length [65, 66], respectively. For instance, the temperatures for the dis-

sociations of Υ states [67] are estimated to be $2T_c$, $1.2T_c$, and T_c for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively, where T_c denotes the deconfinement temperature.

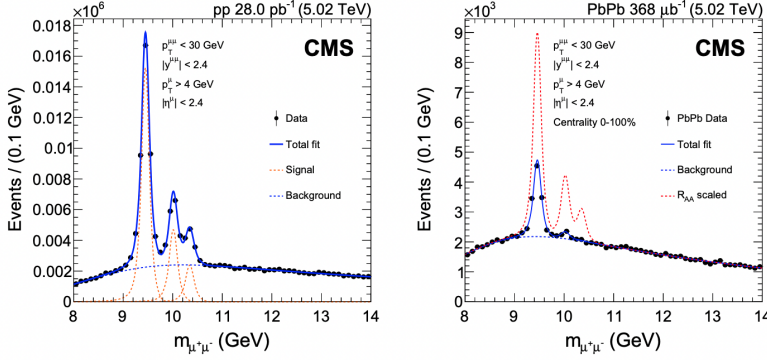


Figure 1.19: Invariant mass distribution of $\mu^+\mu^-$ in pp (left) and Pb–Pb (right) collisions at $\sqrt{s_{NN}} = 5.02$ TeV [16].

Figure 1.19 shows the invariant mass distribution of $\mu^+\mu^-$ in pp (left) and Pb–Pb (right) collisions at $\sqrt{s_{NN}} = 5.02$ TeV [16]. The red dashed line in the right panel shows extracted signals scaled up by the nuclear modification factor, which is the same as extracted signals in pp collisions. Especially, the peak of $\Upsilon(3S)$ is measurable in pp collisions and rarely seen in Pb–Pb collisions. This suggests the sequential melting of bottomonium states depending on the dissociation temperature that can be effectively converted into the radius of the hadron.

Figure 1.20 shows the nuclear modification factor as a function of the number of participating nucleons for different Υ states [16, 17]. The nuclear modification factor decreases with the increasing number of participants for the measured three states. The suppression of $\Upsilon(2S)$ and $\Upsilon(3S)$ is more significant than that of $\Upsilon(1S)$, which is expected, considering the binding energy of each state. The prediction is provided from Simulation for Heavy Ion Collision with Heavy-quark and ONia (SHINCHON). The SHINCHON calculates the dissociation of Υ in the simulated hydrodynamic evolution. The calculation from the SHINCHON agrees with the experimental data for $\Upsilon(1S)$. However, the SHINCHON overestimates the nuclear modification factor for excited states, $\Upsilon(2S)$ and $\Upsilon(3S)$, especially in central collisions. This discrepancy may be attributed

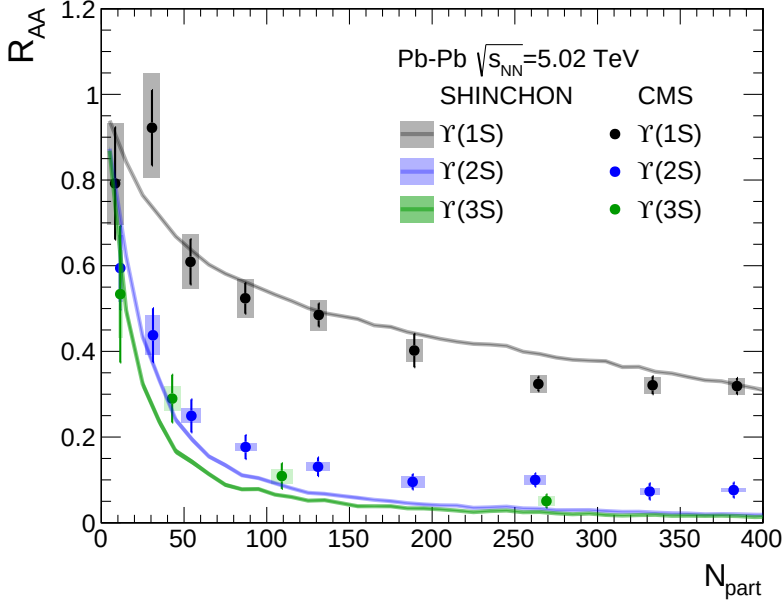


Figure 1.20: The nuclear modification factor of $\Upsilon(nS)$ in Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV measured by CMS collaboration [16, 17]. The model description from SHINCHON (see text) is represented with solid lines [18].

to the no recombination processes in the simulation. Further details on SHINCHON can be found in Ref. [18] and appendix C.

Chapter 2

QGP droplets in small systems

In the previous chapter, experimental evidence for the existence of QGP in heavy ion collisions is reviewed. Surprisingly, various recent measurements provide signatures of QGP in small collision systems at LHC and RHIC energies, where the evolution of the QGP in small systems is expected to be shorter than that of heavy ion collisions. Reflecting shorter evolution time and smaller volume size, the QGP created in small systems is called QGP droplets. Despite a vast of studies to explore the origin of such phenomena favoring the QGP formation in heavy ion collisions, no strong consensus has been reached so far. In this chapter, different measurements that may or may not indicate the formation of the QGP droplets will be discussed.

2.1 Anisotropic flow

The anisotropic particle distribution is one of most evident measurements indicating the formation of the QGP in heavy ion collisions. Surprisingly, many experiments have observed such phenomena in small collision systems [20, 21, 68, 69], where the baryon density had been

expected to be not enough to create the QGP. The observations of anisotropic flow in small systems lead to vast experimental and theoretical strategies and approaches for the origin of the collectivity.

There are two main distinct arguments (still developing) for understanding the origin of collectivity, considering which component contributes more to it. Referring to equation 1.11, the one argument considers constructing the initial collision geometry of small collision systems, not admitting the formation of the QGP droplets. One of the most famous models is the IP-Glasma model [70]. In the past, the geometry of each nucleon had been not considered, and the point-like geometry is accepted. However, the IP-Glasma model constructs the parton distribution inside the nucleon and fluctuation of the distribution [71] to derive the initial geometry (ϵ). In heavy ion collisions, the treatment of the point-like nucleon geometry is still acceptable because effects from the individual parton distribution for each nucleon are averaged out in many-body systems.

The event classification with the initial collision geometry, which can be represented as the impact parameter or the number of participants, is highly connected with the final charged-particle multiplicity in heavy ion collisions. However, it is not the case in small collision systems, especially in pp collisions. The selection of high-multiplicity events does not guarantee the smaller impact parameter at all. However, the anisotropy can still be measured as a function of charged-particle multiplicity [72] to explore the origin of the collectivity in small systems and help constrain the parameters that rule hydrodynamic calculations [73].

Figure 2.1 shows v_2 and v_3 and the model comparisons of them. The configurations of T_RENTo and IP-Glasma are T_RENTo+iEBE-VISHNU [74, 75] and IP-Glasma+MUSIC+UrQMD [76, 77], respectively. The T_RENTo also conducts a hadronic cascade using UrQMD. Measured v_2 is decreasing with decreasing multiplicity in pp collisions and is consistent with multiplicity in p–Pb collisions. Nowadays, v_2 can be observed in any small system, and one interest can arise, such as what the smallest size of the system exhibiting the flow is. From the decreasing trend of v_2 , an experimental hint can be provided. The T_RENTo largely overestimates the experimental data at high-multiplicity events, and the multiplicity dependence is oppositely reproduced of experimental data for

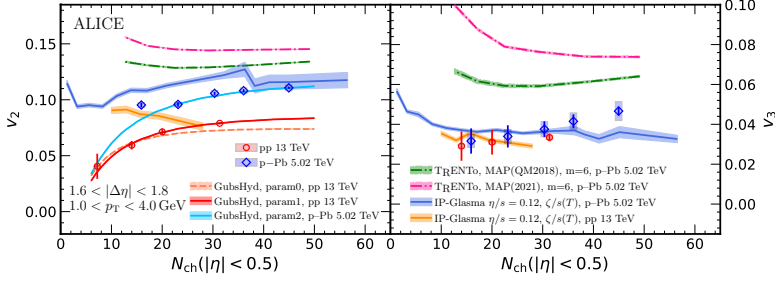


Figure 2.1: v_2 and v_3 as a function of charged-particle multiplicity at midrapidity. The data are compared to different configurations of models.

pp and p-Pb systems. The IP-Glasma model nicely describes the v_2 and v_3 at high-multiplicity events. However, The IP-Glasma model also oppositely estimates the multiplicity dependence of experimental data. On the other hand, the Gubshrd model [78], which is based on an analytical Gubser solution, nicely describes measured data in both magnitude and multiplicity dependence. The Gubshyd model suggests that the current smallest system (most left point in the figure) is still large enough to be hydrodynamized and agrees with the existence of the flow modulation in the system.

On the other hand, the event can be classified with the loosely constrained impact parameter of pp collisions [20, 79]. A theory calculation [19] suggests the impact parameter of pp collisions can be shrunk with the requirement of the number of produced jets. Figure 2.2 shows the impact parameter distribution when two jets or four jets are produced. The distribution is getting narrower with the increasing number of required jets. Because the production of four jets is rare, current experiments start with a single jet [20] or hard probes such as Z-boson [79].

Figure 2.3 shows the associated ridge yield as a function of event-scale selection using the leading particle (a particle having the highest transverse momentum) and reconstructed jets [20]. The event-scale selection is the requirement of the transverse momentum of the particle having the highest transverse momentum or the reconstructed jets. The measured ridge yield is found to be consistent (or mildly increasing) with increasing event-scale selection, which is veiled by its poor statistical un-

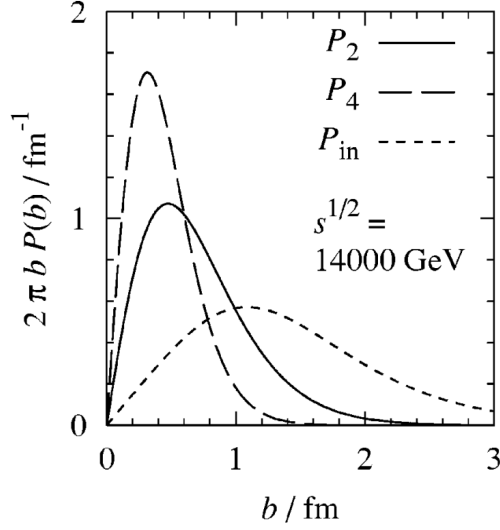


Figure 2.2: Impact parameter distribution of inelastic pp collisions and with the requirement of produced jets [19]

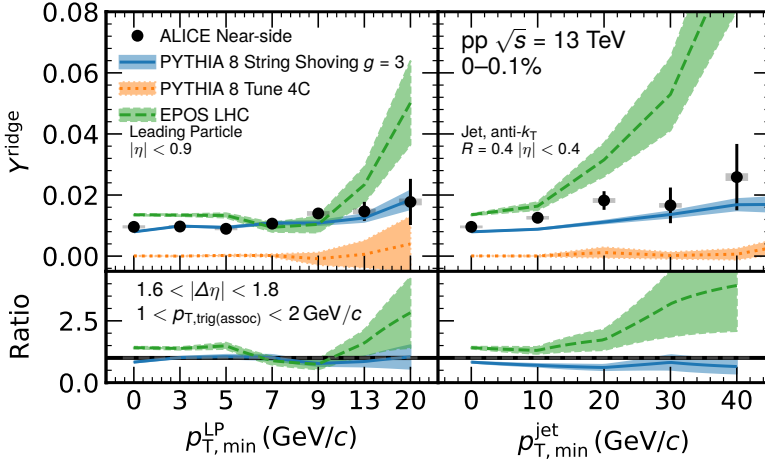


Figure 2.3: Ridge yield as a function of the minimum requirement of the transverse momentum of leading particles and reconstructed jets [20]

certainties.

The comparison of experimental data with model descriptions is also prepared. PYTHIA 8 event generator with the parameter setup of

Tune 4C [80] does not exhibit any collectivity in small systems, and PYTHIA 8 Tune 4C expects the ridge yield to be zero, which is nice as the zero ridge yield guarantees that the measurement is not contaminated by the non-flow particles, mainly attributed to jets. On the other hand, PYTHIA 8 String shoving model [81], which generates transverse pressures by overlapping the expanding strings, nicely describes the experimental data, exhibiting the mild increasing trend of the ridge yield. PYTHIA 8 String shoving model will be further discussed in this section. Finally, EPOS LHC [57] also produces the ridge yield, overestimating the experimental data. The measurement is expected to further constrain the model parameters through the comparison.

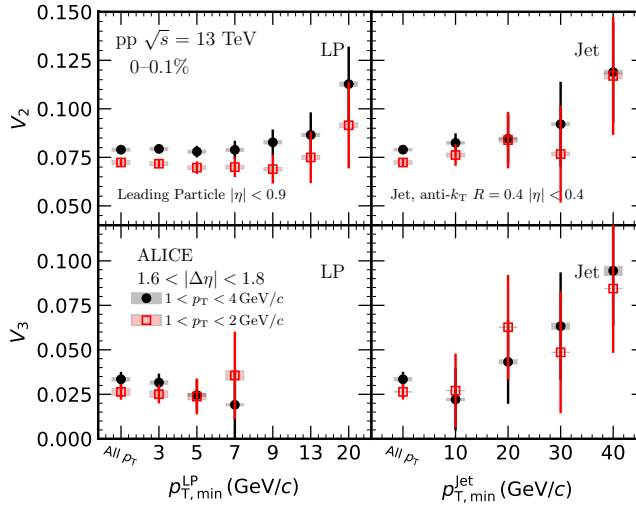


Figure 2.4: v_2 and v_3 as a function of event-scale selection using the leading particle and reconstructed jets

Figure 2.4 shows the v_2 and v_3 as a function of event-scale selection using the leading particle and reconstructed jets for two different transverse momentum ranges. The extracted flow is not increasing with increasing event-scale selection, which is also vailed by poor statistical uncertainties. Together with 2.3, such analyses can be further conducted with the event-scale selection using the di-jet mass or four-jet event when statistics allows.

The other argument for the origin of the collectivity in small sys-

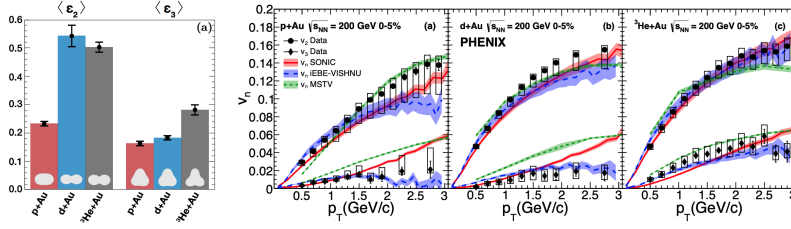


Figure 2.5: (Left) the initial collision geometries of three systems. (right) v_2 and v_3 as a function of transverse momentum from high-multiplicity (0–5%) events [21].

tems is the hydrodynamic evolution, admitting the formation of the QGP droplets and calculating the evolution of it. This picture is nicely tested by PHENIX experiment [21]. To engineer different initial collision geometries, three different collision systems, such as p–Au, d–Au, and ^3He –Au, are constructed. The left panel of figure 2.5 shows the initial collision geometry (ϵ) of three different systems. As expected, ϵ_2 of d–Au and ^3He –Au is much larger than that of p–Au, and ϵ_3 of ^3He –Au is much larger than that of p–Au and d–Au. Such initial states are directly transferred to the final particle distributions, v_2 and v_3 , as shown in the right panel of figure 2.5. v_2 measured in d–Au and ^3He –Au is larger than that in p–Au especially at low p_T , and v_3 measured in ^3He –Au is much larger than that in p–Au and d–Au. implying the identical ordering between the initial geometry and final particle distribution, indicating the hydrodynamic evolution of QGP droplets.

Finally, both arguments of initial geometry modeled by the IP-Glasma and its hydrodynamic evolution can be used at the same time, reproducing many experimental results [22]. There is also an alternative theoretical approach to understand the origin of the collectivity. One of them is the string shoving model in PYTHIA 8. The transversely expanding QDC strings, which are illustrated in figure 2.6, are overlapped with each other, producing the pressure in the transverse direction. Such a pressure results in anisotropic particle distribution without the conventional QGP formation.

The left panel of figure 2.7 shows the associated ridge yield as a

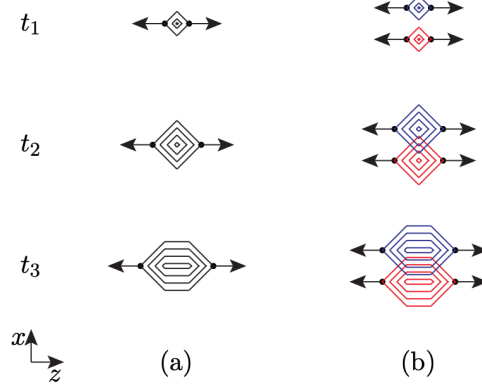


Figure 2.6: Expanding QDC strings and transverse pressure in the overlapping region [22]

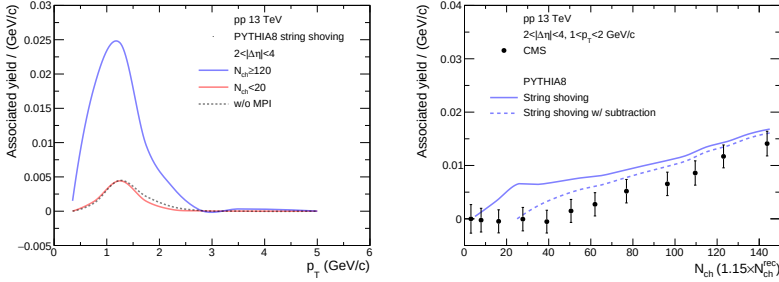


Figure 2.7: (Left) the associated ridge yield as a function of transverse momentum for the event class of high-multiplicity, low-multiplicity, and full range without multi-parton interaction in simulation. (Right) the associated ridge yield as a function of final multiplicity with and without the subtraction of jet contributions [23].

function of transverse momentum [23] for different event classes. The blue (red) solid line denotes the ridge yield in high- (low-) multiplicity pp collisions at $\sqrt{s} = 13$ TeV. The ridge yield can be clearly expected in high-multiplicity events. Unexpectedly, the string shoving model also expects non-zero ridge yield in low-multiplicity events, which is not observed in experimental data. To understand the origin of the ridge yield in low-multiplicity events, the ridge yield is simulated without any multi-parton

interactions (MPI), which is represented as the black dotted line. The calculated ridge yield without MPI is just the same as the yield in low-multiplicity events (with MPI), which implies that the jet component also contributes to the ridge yield in the string shoving model. The right panel of figure 2.7 shows the associated ridge yield as a function of charged-particle multiplicity. The simulated ridge yield is higher than the experimental data because jets contribute to the ridge yield. The ridge yield from jet contributions can be expected by turning off the MPI during the simulation. After the subtraction, better reproduction of experimental data is achieved.

2.2 Strangeness enhancement

Enhanced production of hadrons composed with strange quarks is originally suggested as a key signature of the formation of QGP [82]. However, such an enhancement is not visible with respect to the centrality and energy of the heavy ion collision from RHIC to LHC energies, only exhibiting the enhancement in peripheral collisions, which is later understood as the saturation of the enhancement in large systems. The saturation of the strangeness enhancement is also expected in the grand-canonical statistical model [52]. The extrapolation of the collision system from peripheral heavy ion collisions to smaller systems may help understand the formation of the QGP droplets and the origin of the strangeness enhancement.

Figure 2.8 shows the particle yield ratios to charged pion as a function of charged-particle multiplicity at midrapidity for different particle species. Baryons selected in the plot have a different number of strange quarks. All ratios increase with increasing multiplicity in pp and p-Pb collisions, indicating the strangeness enhancement, and cease to increase in Pb-Pb collisions as the enhancement reaches the saturation except for one point in the most peripheral collisions. The more strange quarks the particle has, the more rapid enhancement is observed. The origin of the enhancement in small systems still remains unknown, motivating the development of different models [52, 57].

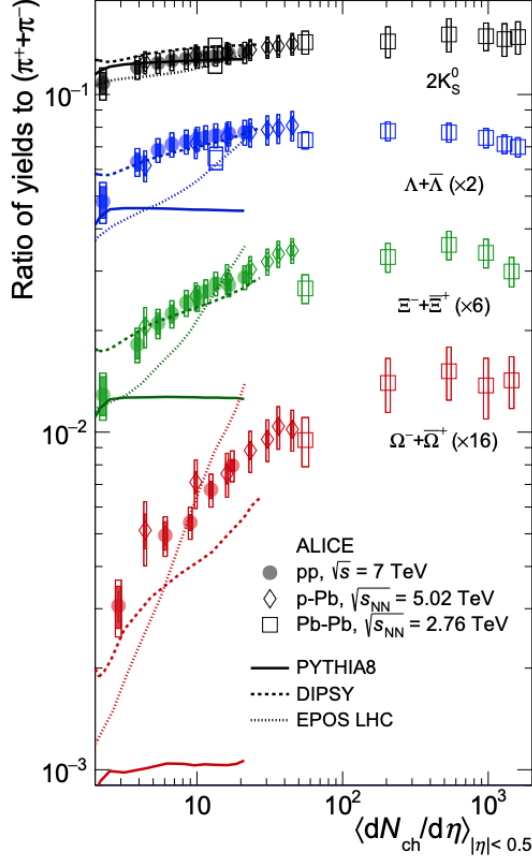


Figure 2.8: p_T -integrated particle yield ratios to charged pions as a function of charged-particle multiplicity [24].

2.3 Inconclusive measurements

In addition to the measurements addressed in previous sections, the measurements of the nuclear modification factor [25] and the energy loss of jets [26] are conducted by many experiments. In contrast with the flow measurement and the strangeness enhancement, the nuclear modification factor in p-Pb collisions does not show any suppression for the particle productions, and the jet quenching is not observed in high-multiplicity pp collisions. Conclusions for these results are not well constructed regarding if signals are not so significant to be observed or if the

QGP droplets are not formed.

2.3.1 Nuclear modification factor

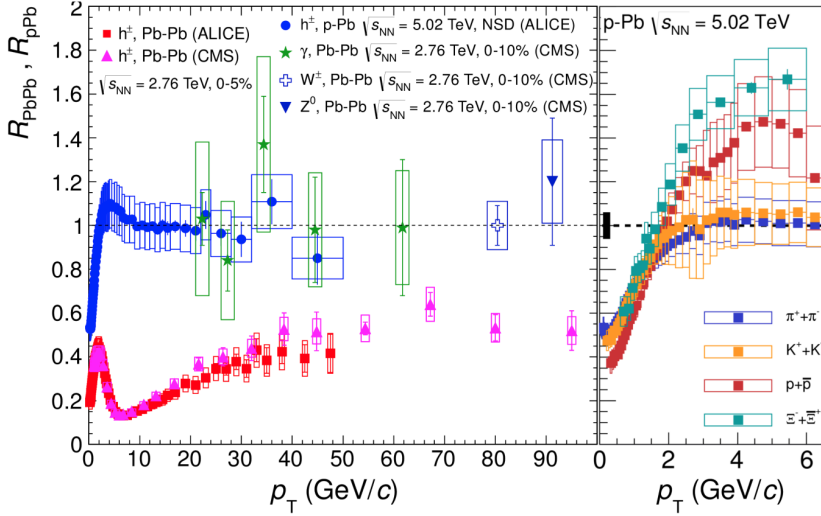


Figure 2.9: (Left) the nuclear modification factor of unidentified charged hadrons in p–Pb and Pb–Pb collisions together with γ , W , and Z bosons, which do not exhibit color charge. (Right) the nuclear modification factor of identified particles in p–Pb collisions [25]

The left panel of figure 2.9 shows the nuclear modification factor of unidentified charged hadrons in p–Pb and Pb–Pb collisions. Significant suppression of the particle production is observed in the central Pb–Pb collisions. However, such suppression is not visible in p–Pb collisions. The nuclear modification factors for γ , W , and Z bosons indicate no suppression. The electroweak bosons do not have color charges and do not interact with QGP. Therefore, the unity of the nuclear modification factors at high p_T guarantees closure of the measurement. The right panel of figure 2.9 shows the nuclear modification factor of identified particles in p–Pb collisions. In addition to the no suppression at high p_T , only baryons exhibit the Cronin-like enhancement at intermediate p_T , speculating the number of constituent quarks.

2.3.2 Jet quenching

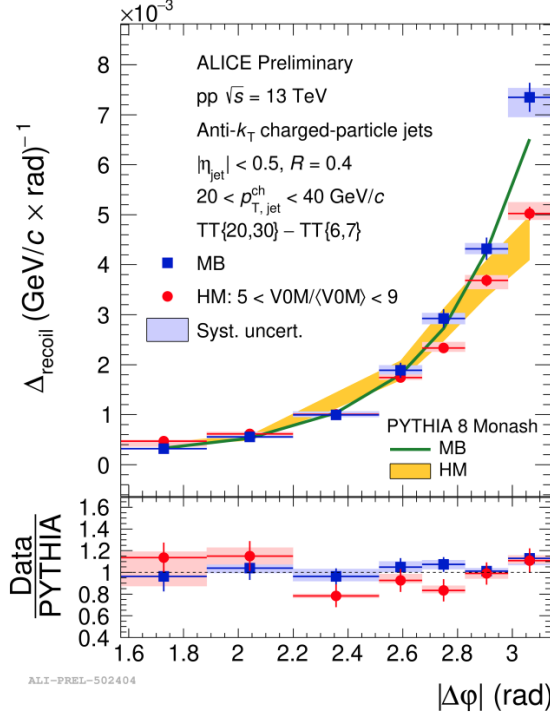


Figure 2.10: The hadron-jet acoplanarity distribution [26]

Figure 2.10 shows the hadron-jet acoplanarity in minimum bias and high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV, where the acoplanarity can be defined as

$$\Delta_{\text{recoil}}(p_{T,\text{jet}}^{\text{ch,corr}}, |\Delta\varphi|) = \frac{1}{N_{\text{trig}}} \frac{d^2 N_{\text{jets}}}{dp_{T,\text{jet}}^{\text{ch,corr}} d|\Delta\varphi|} \bigg|_{\text{TT}\{20,30\}-\text{TT}\{6,7\}}. \quad (2.1)$$

In equation 2.1, $\text{TT}\{20(6), 30(7)\}$ denotes the exclusive $p_{T,\text{trig}}$ interval of $20(6) < p_{T,\text{trig}} < 30(7)$ GeV/c, N_{trig} denotes the number of trigger particles in a given $p_{T,\text{trig}}$ interval, N_{jets} denotes the correlations of the jet and hadrons, $p_{T,\text{jet}}^{\text{ch,corr}}$ denote the jet transverse momentum which is corrected for underlying events with k_T algorithm with $R = 0.2$ [83], and $|\Delta\varphi|$ denotes the relative azimuthal angle between the jet and hadrons.

The acoplanarity of recoil jets in high-multiplicity events shows a substantial suppression or broadening compared to the corresponding distribution obtained from minimum-bias events. These distinctive characteristics are also observed in PYTHIA 8. Additional examinations of the PYTHIA 8 data indicate that the suppression and broadening are due to a bias introduced by the ALICE high-multiplicity trigger. It is expected that further significance can be achieved with enhanced statistics.

Chapter 3

$f_0(980)$ resonances

The $f_0(980)$ resonance has been the subject of intense studies in the field of particle physics for several decades. It was first observed in the late 1970s in $\pi\pi$ scattering experiments [84], and since then, numerous experiments have been conducted to better understand its properties and behavior [85–87]. The $f_0(980)$ resonance is of particular interest because it is one of the lightest scalar mesons and provides a unique window into the nature of the strong interaction. One of the key characteristics of the $f_0(980)$ resonance is its mass that is measured to be near $990 \text{ MeV}/c^2$ [1].

On the other hand, the width of $f_0(980)$ is also another disputing topic, where the width is broadly defined in Ref. [1]. The large width indicates a very short lifetime, which enables the decay products of $f_0(980)$ resonances to interact with the hadron gas before the kinetic freeze-out, where the width is estimated to be from 10 to $100 \text{ MeV}/c^2$.

The $f_0(980)$ resonance is also interesting regarding its quantum number of scalar mesons as $J^{PC} = 0^{++}$ [1]. The anomalous quantum number leads to many suggestions for its internal structure, which also does not reach a strong consensus. The candidate for the internal structure of $f_0(980)$ can be tetraquark state [88], molecular state [89], conventional meson with the orbital excitation for even parity [90], or su-

perpositions of them.

3.1 Internal structure of $f_0(980)$

The internal structure of $f_0(980)$ is still controversial despite a long history of studies. Along with the calculation of $f_0(980)$ mass as $q\bar{q}$ assumption using the QCD sum rule [91], the $f_0(980)$ can be configured as a mixture of strange and light components with a mixing angle of φ and can be expressed as

$$\begin{aligned} |f_0(980)\rangle &= \cos \varphi |s\bar{s}\rangle + \sin \varphi |n\bar{n}\rangle \\ |f_0(500)\rangle &= -\sin \varphi |s\bar{s}\rangle + \cos \varphi |n\bar{n}\rangle, \end{aligned} \quad (3.1)$$

where φ is also very naively measured, and $|n\bar{n}\rangle$ is $\frac{1}{\sqrt{2}}|u\bar{u} + d\bar{d}\rangle$. However, there are too many other $J^{PC} = 0^{++}$ states that are allowed to be explained with the $q\bar{q}$ state, such as glueball states [92]. And these concerns can be resolved with the compact tetraquark model that agrees with the dominant couplings to pairs of pseudoscalar mesons such as $\pi\pi$, KK , and so on [27].

The quantum state of $J^{PC} = 0^{++}$ has been suggested as the chiral partner of the state of $J^{PC} = 0^{-+}$, which is well known as the pseudoscalar meson [93]. While the tetraquark picture can reproduce the inverse mass hierarchy of light scalar mesons ($m_\sigma < m_\kappa < m_{f_0}, m_{a_0}$) compared to the mass hierarchy of pseudoscalar mesons ($m_\pi < m_K < m_\eta$), the tetraquark picture can not support that light scalar mesons are not allowed to be the chiral partner of light pseudoscalar mesons. In this respect, identifying the nature of $f_0(980)$ can significantly contribute to the further understanding the chiral symmetry.

The bag model calculations [27] nicely estimate the invariant masses of light scalar mesons by combining q^2 and $\bar{q}^2$, suggesting the calculated mass of 980 MeV/ c^2 is clearly indicating $f_0(980)$ and $a_0(980)$. Figure 3.1 shows the states of 0^{++} -mesons nonet. According to the model, the $f_0(980)$ resonance is anticipated to have one strange and anti-strange quark to be one of light 0^{++} -mesons nonet. However, The tetraquark structure, helping construct an exotic quantum number configuration of scalar mesons, raises a new question for the radial excitations of diquark

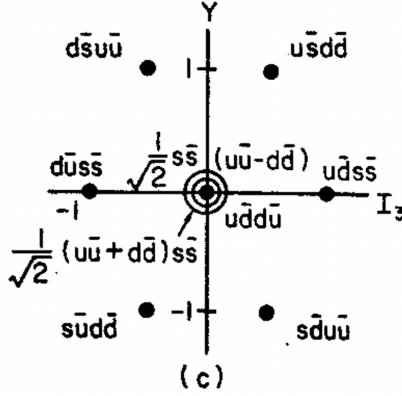


Figure 3.1: The light scalar meson nonet [27]

state [94], increasing the calculated mass and, as a result, overestimating it. Another suggestion raises the possibility of molecular states of two kaons [89]. In addition to the difficulty of the separation between molecular and tetraquark structures such as $X(3872)$ and T_{cc} [95], the measurement of two kaonic femtoscopic correlations hardly concludes its possibility for $f_0(980)$ [96].

3.2 $f_0(980)$ production at LHC energy

The measurement of $f_0(980)$ production at LHC energy can significantly contribute to the understanding of the nature of $f_0(980)$ based on the previous observations. The first measurement is related to rescattering effects in the hadronic phase. The relative suppression of the particle production for short-lived resonances with increasing multiplicity indicates the rescattering effects. Based on this, the observation of rescattering effects for $f_0(980)$ resonances can suggest the lifetime of $f_0(980)$ to be very short. This can be confirmed by measuring the multiplicity-dependent suppression of the $f_0(980)$ production only at low p_T .

The second item is the strangeness enhancement of $f_0(980)$ in small systems. If $f_0(980)$ includes (anti-)strange quarks, the $f_0(980)$ yield will increase with increasing multiplicity with the assumption of no other effects, such as rescattering effects. Meanwhile, rescattering effects can be reduced with the comparison of particle productions between short-lived

resonances. In such a case, one and the other particle would be $f_0(980)$ and $K^{*0}(892)$, respectively, because the $K^{*0}(892)$ includes one strange quark and $f_0(980)$ includes two strange quark or no strange quark while their lifetimes are short.

The other measurements are related to the number of constituent quarks (NCQ). Based on the quark recombination model [97], produced particles can exhibit their NCQ from experimental observables. One observable is the Cronin-like enhancement of the nuclear modification factor in p–Pb collisions for baryons, while such an enhancement is not observed for all measured mesons [25]. The other observable is the NCQ scaling of flow coefficients [98] (see figure 7.14). The NCQ scaling behavior of elliptic flow indicates that the flow of the particles originates from the flow of deconfined quarks. However, the unity of scaled flows between mesons and baryons starts breaking at high kinetic energy due to the radial flow. The Cronin-like enhancement and the splitting of scaled flow for $f_0(980)$ can significantly contribute to the determination of the NCQ for $f_0(980)$.

Chapter 4

Experimental setup

4.1 Detection of charged-particles

Experimental physicists can detect the existence, kinematic variables (i.e., momentum and energy), and species of particles by measuring the interaction between particles and detectors, which are developed for their purpose. In this section, a few methodologies used in many experiments including ALICE will be reviewed. In addition to them, an application of the machine learning method will be briefly introduced.

4.1.1 Transverse momentum

The transverse momentum of charged particles can be measured using the curvature of the particle trajectory in an electromagnetic field that drives the charged particle. As a result of the Lorentz force, the momentum vector can be expressed as $\vec{p} = q(\vec{v} \times \vec{B})$ with the absence of the electric field. When the magnetic field is further assumed to be homogeneous and only one component in the z direction, $B_x = B_y = 0$ and $B_z = B$, the trajectory is calculated to be a helix, which is a coaxial cylinder surface with the cyclotron frequency of ω and the radius of the helix of R . By measuring the trajectory, the transverse momentum of the charged particle can be reconstructed as $p_T = |q|B/|\kappa|$, where the

$|\kappa| = 1/R$ is the curvature.

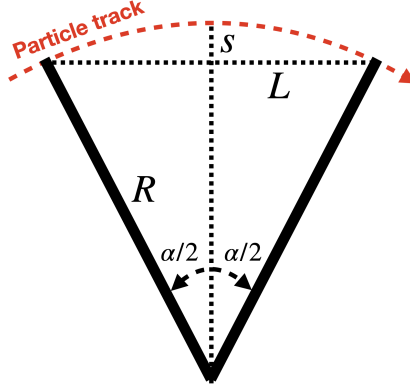


Figure 4.1: Determination of curvature.

The above method is partially applicable because the full helix can hardly be measured with practical detector acceptance. Instead, the measurement of the curvature can still be enough to reconstruct the transverse momentum. From the measurement of a part of the curve's trajectory, the curvature can be estimated to be $\kappa = 2s/L^2$ as shown in figure 4.1 with the assumption of small α , which corresponds to high p_T .

4.1.2 Ionization energy loss

Charged particles with a relativistic velocity lose their energy when the particle crosses the medium primarily by ionization of atoms constituting the medium. The charged particle ionizes the atom by depositing the kinetic energy of the particle to the electrons of the atoms, making the electrons in excited states. Such processes are researched by Bethe [99] and Bloch [100] with the consideration of quantum mechanics in the 1930s. Their research on ionization energy loss is widely utilized for the development of various detectors.

The average energy loss, dE/dx of a charged particle suggested by Bethe and Bloch can be calculated with the Bethe-Bloch formula, which can be expressed as

$$\left\langle -\frac{dE}{dx} \right\rangle = K z^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\frac{1}{2} \ln \frac{2m_e c^2 \beta^2 \gamma^2 W_{\max}}{I^2} - \beta^2 \right] \quad (4.1)$$

The critical parameters in the above equation are summarized in the table 4.1, and other parameters that are not present in the table can be found in the Ref. [1]. It is worth mentioning that so-called density correction and shell correction are not included in the equation.

Table 4.1: Parameters used in equation 4.1, which can also be found in Ref. [1]

Symbol	Definition
K	$0.307 \text{ MeV mol}^{-1} \text{ cm}^2$, coefficient for dE/dx
z	charge number of incident particle
$Z (A)$	atomic number (mass) of absorber
β	ratio of velocity to c
m_e	electron mass
$W_{\max}$	maximum energy transfer in a single collision
I	mean excitation energy

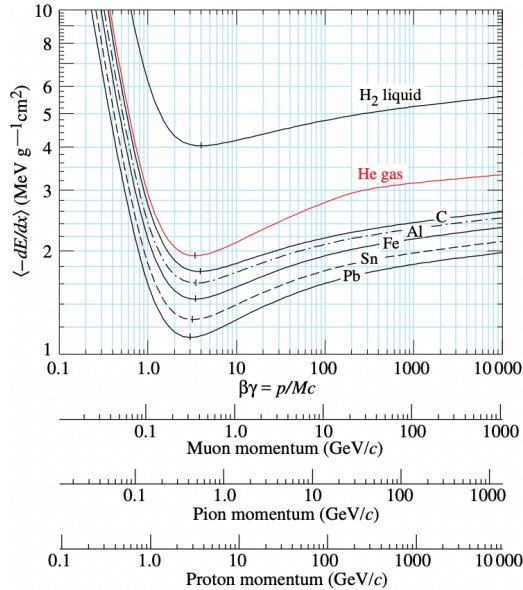


Figure 4.2: Average energy loss in different media for different incident particles as a function of momentum.

Figure 4.2 shows the dE/dx as a function of $\beta\gamma$ ($= p/mc$) and mo-

momentum of μ , π , and proton for different media from H_2 liquid to solid Pb. In low $\beta\gamma$ region ($\beta\gamma < 3$), dE/dx drastically increases with decreasing $\beta\gamma$ because the interaction time is long enough to excite many electrons. In high $\beta\gamma$ region ($\beta\gamma > 10$), dE/dx mildly increases with increasing $\beta\gamma$ because the time dilation induces the incoming particle to excite more electrons. In intermediate $\beta\gamma$ region ($3 < \beta\gamma < 4$), dE/dx is minimized, and the minimum ionization can be used for calibrating and diagnosing the detector. The particles in such a kinematic region are called “minimum-ionizing particles” (MIP). Furthermore, there is a significant difference of dE/dx at the same momentum range between particle species, especially protons and pion and muon, which have similar invariant mass. Based on this, the measurement of dE/dx can be used for identifying particle species.

4.1.3 Time-of-Flight

The flight time of particles from one subsystem to the other subsystem can be used for particle identification. With a given flight length (L) and measured time difference between two subsystems (Δt), the velocity can be calculated as

$$\beta = \frac{L}{\Delta tc}. \quad (4.2)$$

From the equation, the invariant mass of the particle can be deduced as

$$m^2 = p^2 \left[\left(\frac{\Delta tc}{L} \right)^2 - 1 \right]. \quad (4.3)$$

By comparing the calculated mass from the Time-of-Flight method with the species one wants to look for, the particle identification can be achieved. However, the coincident two signals should be proven to originate from the same particle.

4.1.4 Machine learning

There have been many applications of the machine learning methodologies to a wide range of nuclear physics, from the detection of particle species [101] and kinematic variables [102] to the reconstruction of physical observables such as jets [103] and so on. The machine learning

approach can be primarily divided into two different methods, multi-layered neural networks, and decision trees based on gradient boosting. Especially, ALICE reported on the ongoing progress for particle identification using the neural network method [101], which convolutes different information in each layer.

The author of this thesis recently reported on the feasibility study of angular reconstruction of incoming gamma particles using a sampling calorimeter and the gradient-boosting-based machine learning method. The paper shows that the machine learning method provides better resolution when the reconstruction of stochastic processes is conducted. The paper also discusses the reconstruction quality can be effectively controlled by tuning the hyperparameter, which limits and regulate each training process for constructing decision trees and enhancing the training samples. Further details can be found in appendix B and Ref [104].

4.2 Large Hadron Collider in CERN

The European Organization for Nuclear Research (Conseil Européen Pour La Recherche Nucléaire in French), CERN is the laboratory for particle and nuclear physics and the largest hadron accelerator in the world. The CERN is located at the border between France and Switzerland near Geneva and is operated, supported, and utilized by scientists more than 10,000 and countries of 23 member states.

The Large Hadron Collider (LHC) [105] in the CERN is currently the most powerful and the most giant main accelerator ring in the world. Having a circumference of 26.7 km, the LHC can accelerate protons in rotating two identical beam pipes up to the energy of 7 TeV, achieving the center of energy of 14 TeV for proton-proton collisions and of 5.02 TeV for nucleon-nucleon collisions with the luminosity of $\mathcal{O}(10^{34}) \text{ cm}^{-2}\text{s}^{-1}$. The main parameters [37] for the LHC are listed in table 4.2.

The LHC can not fully accelerate particles without the help of pre-accelerations. Rather, particle beams experience several pre-acceleration stages before the beams are injected into the LHC. The first stage for the pre-acceleration is performed with the Linear accelerators while protons and ions are accelerated by LINAC2 and LINAC3, respectively, to the energy of 50 MeV (per nucleon for ions). The accelerated protons

Table 4.2: Key parameters of the LHC design [37]

Parameter	Value
Circumference	26.7 km
Number of dipole magnets	1232
Number of quadrupole magnets	858
Number of higher-order magnets	7210
RF cavities	400.8 MHz
Superconductor temperature	1.9 k
Injection energy	450 GeV
Accelerated energy (proton)	7 TeV
Distance between bunches	7 m
Bunch spacing	25 ns
Number of bunches per each proton beam	2808
β^*	0.55 m

and ions are injected into the Proton Synchrotron (PS) Booster and Low Energy Ion Ring, respectively, to be accelerated to the energy of 1.4 GeV. PS takes accelerated protons and ions to sequentially accelerate them to the energy of 25 GeV, transferring them to the Super Proton Synchrotron (SPS), which, in turn, accelerates particles to 450 GeV. The full system of the accelerator is illustrated in figure 4.3.

The LHC is equipped with metallic chambers containing the electromagnetic field called radiofrequency (RF) cavities [105] along the accelerator ring. The particles in the ring receive an electrical impulse when the oscillation of the electromagnetic field is synchronized with the circulation of particles inside the ring. The synchronization regulates the energy of particles by increasing (decreasing) the acceleration force when a particle is under (over-)accelerated. The frequency for the electromagnetic field is set to 400.8 MHz, which corresponds to the bunch separation of 2.5 ns, while actual bunch separation is conservatively set to 25 ns.

The guidance of the beam in the ring is performed with the superconducting electromagnets. The dipole magnet is used to maintain the circular orbit, generating a magnetic field of 8.33 T for accelerating the proton beam to 7 TeV [106]. The quadrupole magnet is used to preserve the beam to be focused during beam traveling inside the ring. The higher-

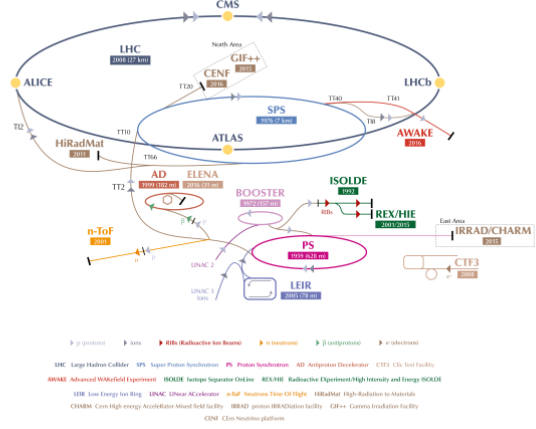
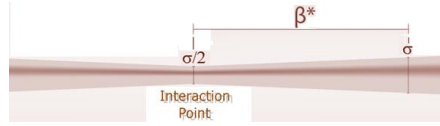


Figure 4.3: Accelerator complex in CERN [28]

order multipole magnets are used for additional beam adjustment. The numbers of the dipole, quadrupole, and higher-order multipole magnets are 1232, 858, and 7210, respectively.

Figure 4.4: Definition of β^*

In accelerator physics, the beam emittance (ϵ) [107, 108] represents how many particles in a bunch are deviating spatially and energetically, implying that smaller ϵ suggests better beam transportation and more probability of collisions between bunches. The amplitude function, β can be expressed as $\beta = \pi\sigma^2/\epsilon$, where the σ is the size of the bunch. The schematic of β can be found in figure 4.4. The amplitude function, β is designed to be minimized to achieve higher luminosity in the experiment considering the expression of the luminosity as

$$\mathcal{L} \sim \frac{fN^2}{4\epsilon\beta^*}, \quad (4.4)$$

where the N , f , and β^* are the population of the bunch, the revolution

frequency, and the amplitude function at the interaction point, respectively [109]. The improvement on β^* is essential to enhance the luminosity.

In the LHC, there are four interaction points where bunches collide, and huge detector systems are installed by each experiment along the accelerator ring. Out of these, A Large Ion Collider Experiment (ALICE) [34, 110] is only the experiment mainly dedicated to focusing the heavy ion physics. The vast measurement from ALICE is to research the properties of QGP. On the other hand, A Toroidal LHC Apparatus (ATLAS) [111] and Compact Muon Solenoid (CMS) are the two largest experiments capable of particle and nuclear physics together. The most famous contribution from them is the observation of Higgs boson [112, 113]. In addition to them, LHC beauty (LHCb) [114] conducts the study of bottom quark physics. In the weak interaction inducing the changes of flavor, the measurement for the transition of bottom quarks to up-type quarks provides the rates of CP violation, which explain the matter-dominant universe. The measurement of $B_s^0 \rightarrow \mu\mu$, which is sensitive to New Physics, successes in verifying the Standard Model expectation by LHCb [115].

4.3 ALICE detector

ALICE, which is the experiment dedicated to primarily studying the heavy ion physics in Pb–Pb collisions, is optimized for the environment of Pb–Pb collisions, where the produced number of charged particles is expected to be the order of 2,000 for the central collisions [116]. Because the dominant particles are thermally produced, the high momentum resolution at low p_T region is essentially required. Based on this, the requirement of the ALICE detector is to separate and reconstruct each charged particle and to identify the particle species. The analysis in the present thesis uses subsystems of VZERO (V0), Zero Degree Calorimeter (ZDC), Inner Tracking System (ITS), Time Projection Chamber (TPC), and Time Of Flight (TOF), and they will be intensively described while descriptions for other systems can be found in Ref [34, 117]. Table 4.3 summarizes the usage of each subsystem.

Figure 4.5 shows a schematic of the ALICE detector. The overall di-

Table 4.3: Main purpose of each subsystem in the present thesis

Detector	Tasks
V0	Triggering, Centrality, Event plane
ZDC	Centrality
ITS	Tracking, Vertexing
TPC	Tracking, PID
TOF	PID

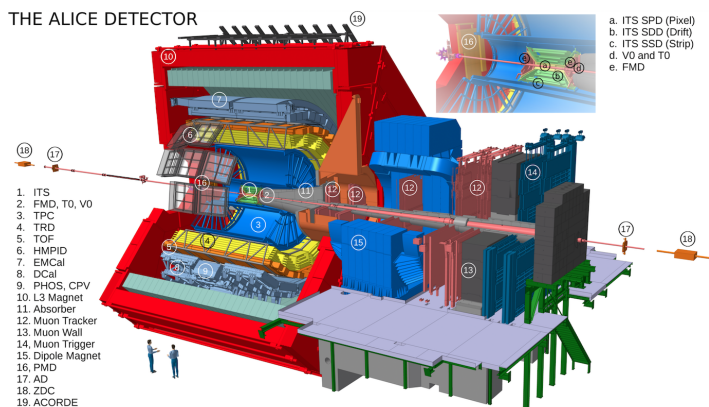


Figure 4.5: Schematic of the ALICE detector in Run 2.

mension and weights of the apparatus are $16 \times 16 \times 26 \text{ m}^3$ and 10,000 t, respectively. The multilayered configuration ensures high acceptance by capturing produced particles in each event. The tracking system is located close to the center of the interaction point in order to accurately reconstruct the primary vertex and charged-particle tracks with high resolution. The tracking chamber is embedded in the solenoid magnet, which provides a magnetic field of 0.5 T. The calorimeters on the outer layer measure particles by absorbing the energy of the particle.

4.3.1 The V0 Detector

The V0 is the hodoscope of two arrays of the scintillator counter, called V0A and V0C, each made of 32 plastic scintillators [30]. The V0A is located at 340 cm from the average position ($z = 0$) of the interaction point. The V0C is located at 90 cm from the interaction position in the

opposite direction to the VOA. The V0 measures the passage of charged particles, covering the full azimuthal angle in the pseudorapidity range of $-3.7 < \eta < -1.7$ and $2.8 < \eta < 5.1$ by V0C and V0A, respectively. The sum of the V0A and V0C signals is widely used for the event classification and triggering and is named the VOM.

The coincident signal of the V0A and V0C is used as an event trigger signal of a beam-beam collision during data taking online. ALICE additionally applies an offline event selection to reject possibly machine-induced backgrounds, measuring the time of the passages.

One of the main usages of the V0 in an offline analysis is the determination of the centrality (or the multiplicity) class. The centrality class can be either defined using the ITS, which uses the midrapidity particles, while the V0 uses forward rapidity particles. The analyses, which use the midrapidity particles, can be biased when the ITS is used for the determination of the centrality class owing to so-called auto-correlations. The auto-correlations effect can be observed in figure 4.6, showing the yields of $K^\pm$ and K^0 as a function of the multiplicity measured in midrapidity (forward rapidity) on the left (right) panel. The midrapidity definition shows biased enhancement for the charged kaon yield, while the unity of two kaon yields can be preserved with the forward rapidity definition, implying that no auto-correlations can be expected with the VOM.

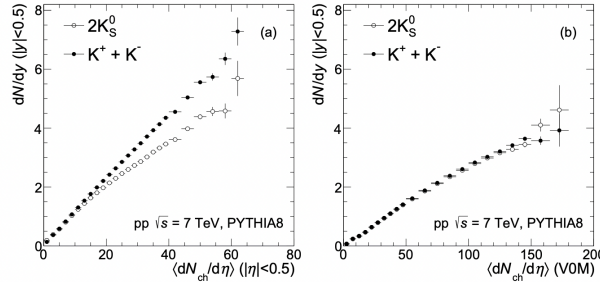


Figure 4.6: Neutral and charged kaon yields as a function of multiplicity class defined by the ITS (left) and V0 (right) in PYTHIA8 simulations of pp collisions at $\sqrt{s} = 7$ TeV [10].

The centrality class can be represented as a percentage that expresses how short the impact parameter is or how many nucleons are

participating in the collision, where the impact parameter is closely related to the number of participating nucleons and the number of produced charged particles from the collisions. Based on such correlations, the multiplicity class in small collision systems can also be represented as the percentage that expresses how many charged particle is produced, even though such correlations are weak in pp or pA collisions. The lower percentage represents a larger number of produced charged particles, called central collisions, and can be expressed as

$$\frac{\int_{V0_C}^{\infty} (dN_{\text{evt}}/dV0)dV0}{\int_{V0_{\text{AP}}}^{\infty} (dN_{\text{evt}}/dV0)dV0} \times \frac{1}{\text{AP}}, \quad (4.5)$$

where the AP represents the anchor point that background contamination is negligible.

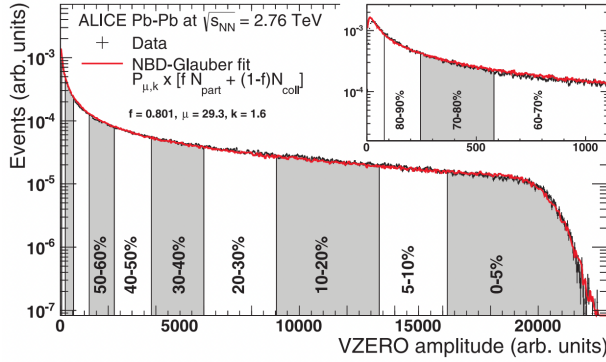


Figure 4.7: VOM amplitude distribution in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. The VOM amplitude is reproduced with the NBD-Glauber model [29].

The centrality distribution can be modeled with the Glauber model [46] convoluted with a model of particle production represented as negative binomial distribution (NBD) using the measurement of charged particle multiplicity in a wide range of rapidity in pp collisions [118]. The Glauber model calculates the number of participating nucleons for a given impact parameter by reproducing the multiplicity distributions assuming two different sources of particle production, soft and hard interactions. Figure 4.7 shows the VOM distribution in Pb-Pb collisions at

The ZDC is sampling calorimeters made of tungsten alloy for passive material and quartz fibers for active detector [38]. The ZDC is located 114 m from the interaction point. The ZDC consists of two hadronic calorimeters for measuring spectator nucleons, which represent the nucleons not participating in the collision, and an additional electromagnetic calorimeter (ZEM). The hadronic calorimeters consist of two distinct detectors for spectator neutrons (ZN) and spectator protons (ZP). A Schematic of ZDC location can be found in figure 4.9.

The ZDC operates in a very harsh radiation environment up to 10^4 Gy/day, and quartz fibers are selected because of their intrinsic radiation hardness. The fibers generate signals via Cherenkov light, which is very fast thanks to the emission process and is also hardly sensitive to radioactivity. The main parameters for the ZDC can be found in table 4.4.

Table 4.4: Main parameters for the ZN and ZP [38]

Detector	ZN	ZP
Dimension	$7 \times 7 \times 100 \text{ cm}^3$	$23 \times 12 \times 150 \text{ cm}^3$
Absorber	W-alloy	brass
Density	17.6 g/cm^3	8.5 g/cm^3
The number of slabs	44	30
Thickness of slabs	1.6 mm	4 mm
Number of fibers	1936	1680
Fiber spacing	1.6 mm	4 mm
Diameter of Fiber	$365 \mu\text{m}$	$550 \mu\text{m}$

The primary purpose of the ZDC is the determination of the multiplicity class by measuring the spectator nucleons. The nuclear modification factor of charged particles in p-Pb collisions for different VOA multiplicity classes shows dynamical biases due to the large multiplicity fluctuations for the small number of participating nucleons, which can be seen in figure 4.10. On the other hand, the measurement of spectator nucleons, where most of the nucleons in Pb ions are spectator, can provide the multiplicity definition without such a dynamical bias.

In the present thesis, the hybrid method reported in Ref. [31] is used to measure the nuclear modification factor using the ZDC. The ZN is used for this method, measuring nucleons knocked out from wounded nucle-

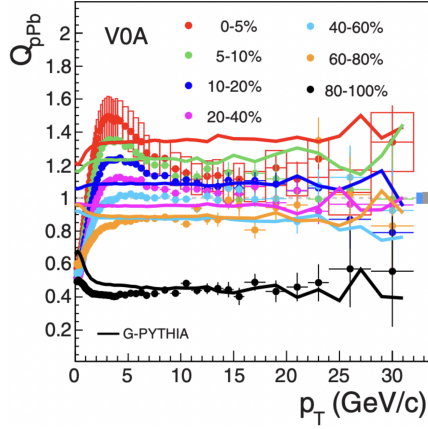


Figure 4.10: Nuclear modification factor of primary charged particle for different multiplicity classes provided by the V0A in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Uncertainties from normalization and scaling of the number of binary collisions are shown as gray and light blue boxes, respectively [31].

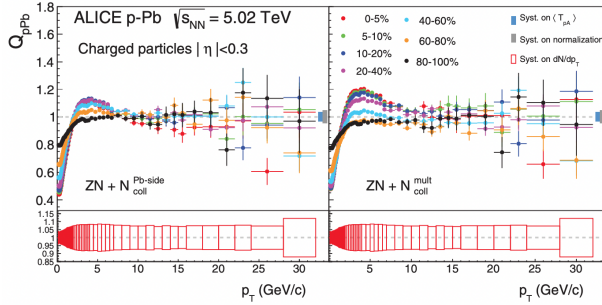


Figure 4.11: Nuclear modification factor of a primary charged particle with the hybrid method for scaling [31].

ons or emitted from the colliding nucleus by nuclear de-excitation processes. Two assumptions for the hybrid method are the followings. The first main assumption is that the event selection based on the ZN does not induce any bias on high- p_T production and on the bulk production at midrapidity. The second assumption is that the charged particle produced at mid-rapidity is well scaling with the number of participating nucleons. Figure 4.11 shows the nuclear modification factor of charged

particles using the hybrid method. The hybrid method provides the least biased multiplicity definition using the ZN signals.

4.3.3 The ITS Detector

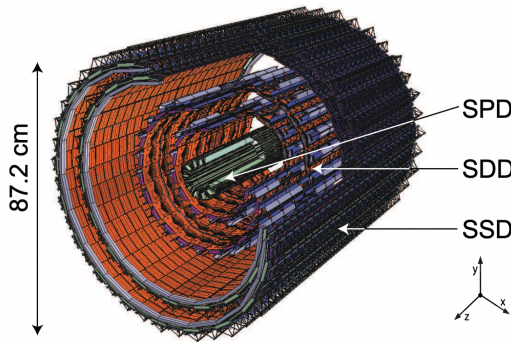


Figure 4.12: Schematic of the ITS [32]

The ITS is the closest detector to the interaction point surrounding the beam pipe [32]. The ITS is composed of six layers of cylindrical silicon detectors. Two innermost layers are Silicon Pixel Detectors (SPD), and then Silicon Drift Detector (SDD) is placed for two middle layers, and the two outmost layers are Silicon Strip Detectors (SSD). Figure 4.12 shows the layout of the ITS detector. Table 4.5 shows the dimension of each layer.

Table 4.5: Dimension of the ITS

Layer	Name	Type	Acceptance	Radius (cm)	Length (cm)
1	SPD	pixel	$ \eta < 2$	3.9	14.1
2	SPD	pixel	$ \eta < 1.4$	7.6	14.1
3	SDD	drift	$ \eta < 0.9$	15	22.2
4	SDD	drift	$ \eta < 0.9$	23.9	29.7
5	SSD	strip	$ \eta < 1$	38	43.1
6	SSD	strip	$ \eta < 1$	43	48.9

The main purposes of the ITS are the reconstruction of the primary vertex with the resolution of $100 \mu\text{m}$ and the improvement of the momentum resolution of charged tracks reconstructed by the TPC. More-

over, the ITS can identify particle species at low transverse momentum down to $p_T \approx 0.15$ GeV/c by measuring the energy loss of the charged particle (see section 4.1.2) in the SDD and SSD.

4.3.4 The TPC Detector

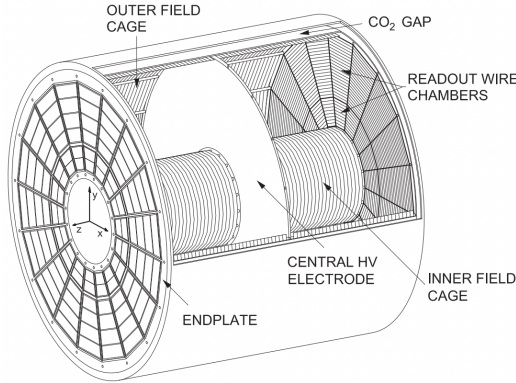


Figure 4.13: Schematic of the TPC [33]

The TPC is a 5-m-long large cylindrical gas detector with an inner (outer) radius of 85 (250) cm, where two different readout chambers cover the radial range of 85 to 132 cm (135 to 250 cm) for the inner (outer) readout chamber [33]. The 90 m³ gas in the TPC is a mixture of Ar/CO₂/N₂ (90/10/5) from 2016, while Ne had been used instead of Ar. The TPC is the main tracking detector in ALICE, together with the ITS. The TPC accurately reconstructs 3-dimensional charged tracks for as many particles as possible. The TPC measures the curvature and energy loss to determine the momentum and species of the charged particle, respectively.

The operation of the TPC is performed by the ionization of gas particles owing to the charged particle that traverses the gas volume. Electrons emitted due to the ionization of the gas molecules are accelerated by a longitudinal electric field of 400 V/cm, which is created by the position of 100 kV on both sides of the TPC cage. The accelerated electrons drift toward the end plates. The anode wires of the Multi-Wire Proportional Chamber (MWPC) [119] detector measures the number of electrons that is amplified in anode wires. The drift velocity of electrons in

the gas volume is constant because of sequential collisions between gas molecules and electrons, enabling one to reconstruct z coordinate with the measured drift time, while x and y coordinates can be reconstructed by measuring which wires are fired. The maximal drift time of electrons in the potential of 100 kV is about $100 \mu\text{s}$, which is long enough to affect the data acquisition rate negatively.

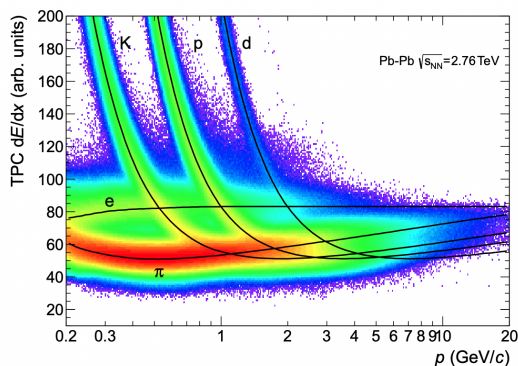


Figure 4.14: Particle identification using the TPC [34]

The energy loss of charged particles in the gas volume can be measured from the number of amplified electrons together with the flight length from the track reconstruction. Based on the Bethe-Bloch formula, particles having different invariant masses deposit different energies at the same momentum. Figure 4.14 shows the energy loss as a function of momentum for different particle species, and particle species are clearly separated in the low momentum range.

4.3.5 The TOF Detector

The TOF is a hodoscope surrounding the TPC (thus, covering the full azimuthal angle) with the pseudorapidity coverage of $|\eta| < 0.9$. The schematic of the TOF can be found in the left panel of figure 4.15. The TOF consists of 1593 Multi-gap Resistive Plate Chambers (MRPC). Each MRPC has a cross-sectional surface of $122 \times 1.3 \text{ cm}^2$ with an active area of $120 \times 7.4 \text{ cm}^2$ and provides an intrinsic timing resolution of 40 ps. The main purpose of the TOF is to identify the particle species by measuring the Time-of-Flight, where the method can be found in section 4.1.3. The

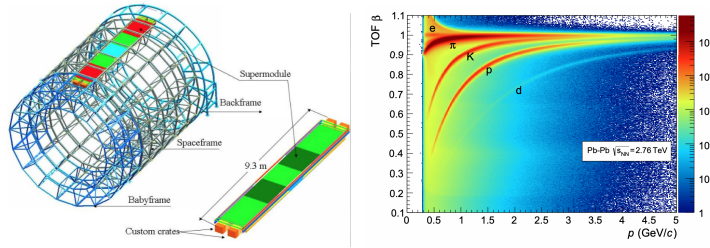


Figure 4.15: Schematic of the TOF (left) [35] and Particle identification using the TOF (right) [34]

measured flight time is constructed by calculating the relative time between the time of the primary collision and the time measured by the TOF. The right panel of figure 4.15 shows the reconstructed velocity using the Time-of-Flight as a function of momentum, and the separation of the particle species is clearly seen.

4.3.6 Track reconstruction

The track reconstruction in ALICE is mainly performed with the central barrel detector of the ITS and TPC and the TRD if a track reaches. First of all, the reconstruction starts with clustering the raw signals of each subsystem separately. Each cluster includes the position, signal amplitude, and signal time.

Based on the clustered information, the reconstruction of charged tracks is performed with the method called inward-outward-inward scheme [120]. The “inward” (“outward”) denotes that the track reconstruction starts from the outer (inner) layer, and the track is getting connected with signals in the inner (outer) layers. The track finding and fitting processes are based on the Kalman-filtering approach [121]. The first inward step starts with the TPC signal, which is represented as the red line in figure 4.16. A track seed, which is an initial direction and position for the reconstruction algorithm, is constructed with cluster information in the initial layer. The track seed keeps updating for each reconstruction step when the seed satisfies standard proximity selection. One of the selections is the number of TPC clusters participating in the track reconstruction. If a track has clusters contributor larger than 159,

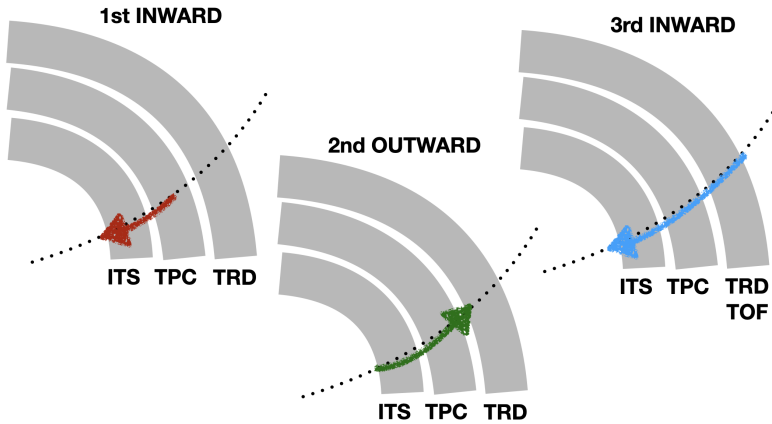


Figure 4.16: Schematic of the track reconstruction

such a track can be rejected because the total number of pad rows from inner to outer readout is 159 [34]. The other selection would be the requirement for the minimum number of participating clusters concerning the resolution of the reconstructed track.

The reconstructed track in the first-inward step becomes the track seed for the second-outward step, which starts from the ITS (the green line in figure 4.16). The update of the seed continues with the proximity selection. If there is no compatible cluster expected, A penalty factor is added to the χ^2 . As a result of the second step, each TPC track (seed) yields track candidates in the ITS. The ITS tracks for each TPC track are ranked according to the reduced χ^2 [120].

After the reconstruction of the ITS track, the third-inward process begins, represented with the blue line in figure 4.16. In this step, matches of the clusters in the TRD and TOF are performed when clusters are there. For instance, the integrated track length and the flight time can be expected for particle identification. The purity of reconstructed primary particles can be further enhanced using the kinematic information. One of the main parameters is the distance of the closest approach (DCA) to the primary vertex [122], which is also applied in the present thesis as the analysis uses primary charged particles.

Chapter 5

Analysis method

5.1 Event Selection

The data samples of p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV for the present analysis are collected during 2016 in Run 2 using the minimum-bias trigger fired by the V0. Additional offline selections are applied to the collected data to reduce backgrounds and to ensure a better quality of data used in the present analysis. The total luminosity of analyzed data passing full selections is 0.24 nb^{-1} [123]. The standard offline selections implemented in the AliPhysics package and applied in the present analysis are listed below.

- Standard physics selection using AddPhysicsSelectionTask
- Minimum Bias Trigger
 - (inputHandler → IsEventSelected()) & (AliVEvent::kINT7)
- Rejection of pileup events at the SPD.
 - AliMultSelection → IsEventSelected(), which includes
 - * inputHandler → IsPileupFromSPD()
 - * inputHandler → IsSPDClusterVsTrackletBG()

Table 5.1: Charged particle multiplicity in different multiplicity classes at $\sqrt{s_{NN}} = 5.02$ TeV. The 60–80 (100)% class is defined with the ZDC (VOA).

Estimator	0–20%	20–40%	40–60%	60–80 (100) %
ZDC	28.12	23.70	17.99	11.68
VOA	35.42	23.12	15.89	6.88

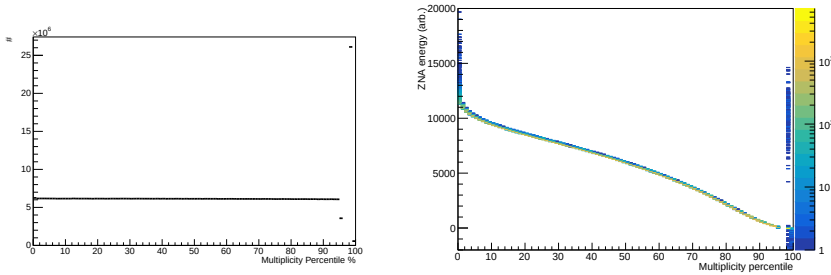


Figure 5.2: (left) multiplicity percentile distribution provided by the ZDC. (right) scatter plot of multiplicity percentile and ZNA tower energy.

not biased from such an effect, the ZDC is not used when defining the multiplicity class lower than 80%.

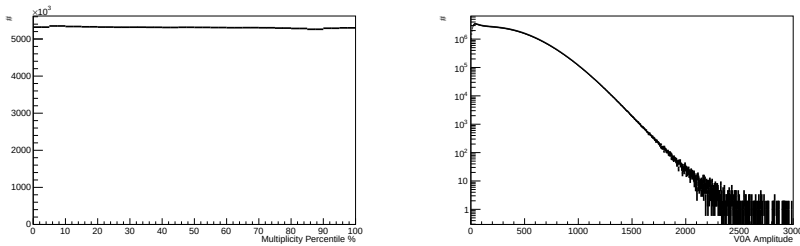


Figure 5.3: (left) multiplicity percentile distribution provided by the VOA. (right) the VOA amplitude.

The left panel of figure 5.3 shows the multiplicity percentile distribution obtained using the VOA. The distribution is uniform for the entire

range of 0–100%. The right panel of figure 5.3 shows the VOA amplitude distribution.

5.2 Track Selection

The $f_0(980)$ resonance dominantly decays to $\pi^+\pi^-$ with the branching ratio (B.R.) of $46 \pm 6\%$ [125]. Therefore, primary charged pions are selected with high quality to reconstruct the $f_0(980)$. The charged pions reconstructed in the TPC are required to satisfy $p_T > 0.15$ GeV/ c and $|\eta| < 0.8$ to ensure uniform detection acceptance. The reconstructed tracks are also required to pass the standard selection criteria, which are reported in [122], to ensure the high quality of the reconstruction. One of the important selections is the number of participating clusters to be larger than 70 in the TPC (out of 159) for good track momentum resolution as well as the existence of two hits on the ITS (one hit in SPD) to have a track seed in the second OUTWARD procedure (see section 4.3.6). The purity of the primary tracks can be achieved with the requirement on the distance of closest approach (DCA) to the primary vertex [122]. The transverse (d_z) and longitudinal (d_{xy}) DCA are required to be $d_z < 2$ cm and $d_{xy} < (0.0105 + 0.0350 \times p_T^{-1.1})$ cm, respectively. Such treatments can be conducted with

- track $\rightarrow$ TestFilterMask(0x20)

The particle identification (PID) is performed using the TPC and TOF with the method of ionization energy loss described in section 4.1.2 and time-of-flight described in section 4.1.3, respectively. The $\sigma_{\text{TPC(TOF)}}$ is defined as the difference between measured ionization energy loss (time-of-flight) and prediction with pion mass assumption divided by its resolution. The $\sigma_{\text{TPC(TOF)}}$ is required to satisfy

- $|\sigma_{\text{TOF}}| < 3$ and $|\sigma_{\text{TPC}}| < 5$ if TOF is available
- $|\sigma_{\text{TPC}}| < 2$ if TOF is not available

where the probability that signals in the TPC and TOF are coming from different tracks can be considered as

- `fPIDResponse → GetTOFMismatchProbability( track ) < 0.01`

If the probability that the signals are not originating from the same track is larger than 1%, The PID is performed with the TPC as a standalone system.

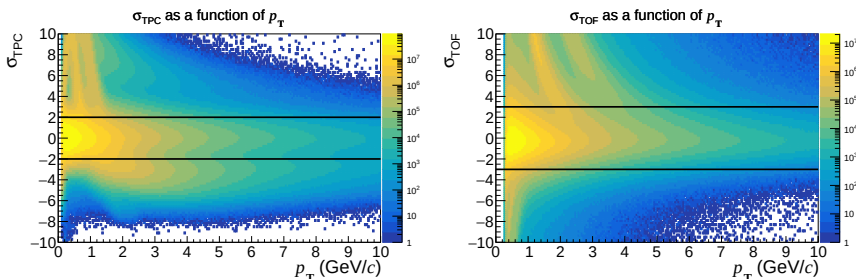


Figure 5.4: (left) σ_{TPC} as a function of p_T for charged pions. (right) σ_{TOF} as a function of p_T for charged pions. Black lines represent the limit of the selection.

Figure 5.4 shows the $\sigma_{\text{TPC(TOF)}}$ as a function of pion p_T evaluated using the TPC and TOF with the assumption of the charged pion mass.

5.3 Signal Extraction

Two oppositely-charged pions, which fully satisfy the full requirements described in section 5.2, are paired to reconstruct $f_0(980) \rightarrow \pi^+ \pi^-$ events at $-0.5 < y < 0$ [126] in p-Pb collisions. There are many pairs of two pions, which are uncorrelated with each other, and the $f_0(980)$ signals would be contaminated by such uncorrelated pairs. The Like-sign background subtraction method [127] is implemented to subtract combinatorial backgrounds coming from uncorrelated pairs. Invariant mass distribution of uncorrelated Like-sign and Unlike-sign pairs would be the same as each other. Therefore, by pairing pions that have the same charge, backgrounds from the uncorrelated pairs can be estimated. Pairs of $(\pi^+$ and $\pi^-)$ are called as Unlike-sign, and pairs of $(\pi^+$ and $\pi^+)$ or $(\pi^-$ and $\pi^-)$ are called as Like-sign. By subtracting Like-sign from the Unlike-sign, The signal of resonances decaying to $\pi^+ \pi^-$ is extracted. Like-sign

distribution is constructed by combining invariant mass distributions of $\pi^+\pi^+$ and $\pi^-\pi^-$ pairs with the geometric average.

$$\mathcal{Y}_{\text{Like-sign}} = 2\sqrt{y_{++}y_{--}} \quad (5.1)$$

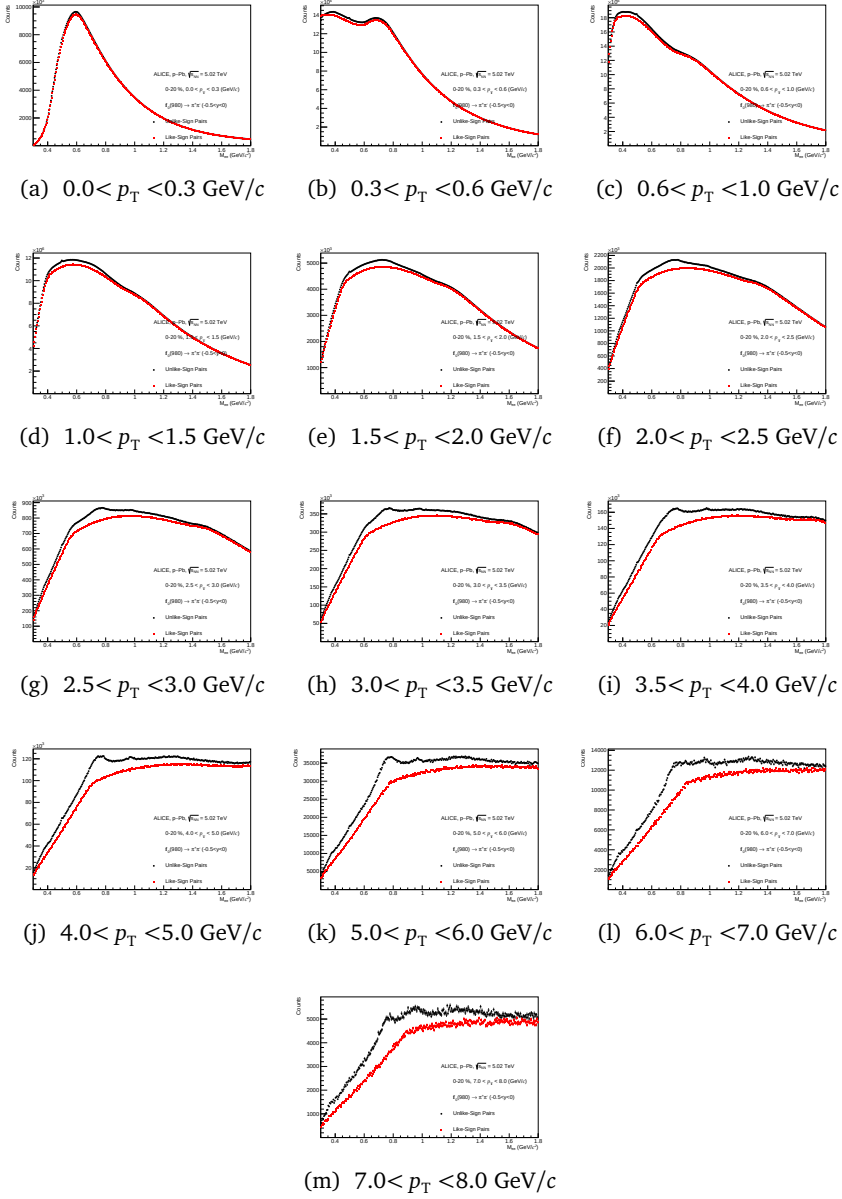


Figure 5.5: Unlike-sign (black) and Like-sign (red) distributions in 0–20 % p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

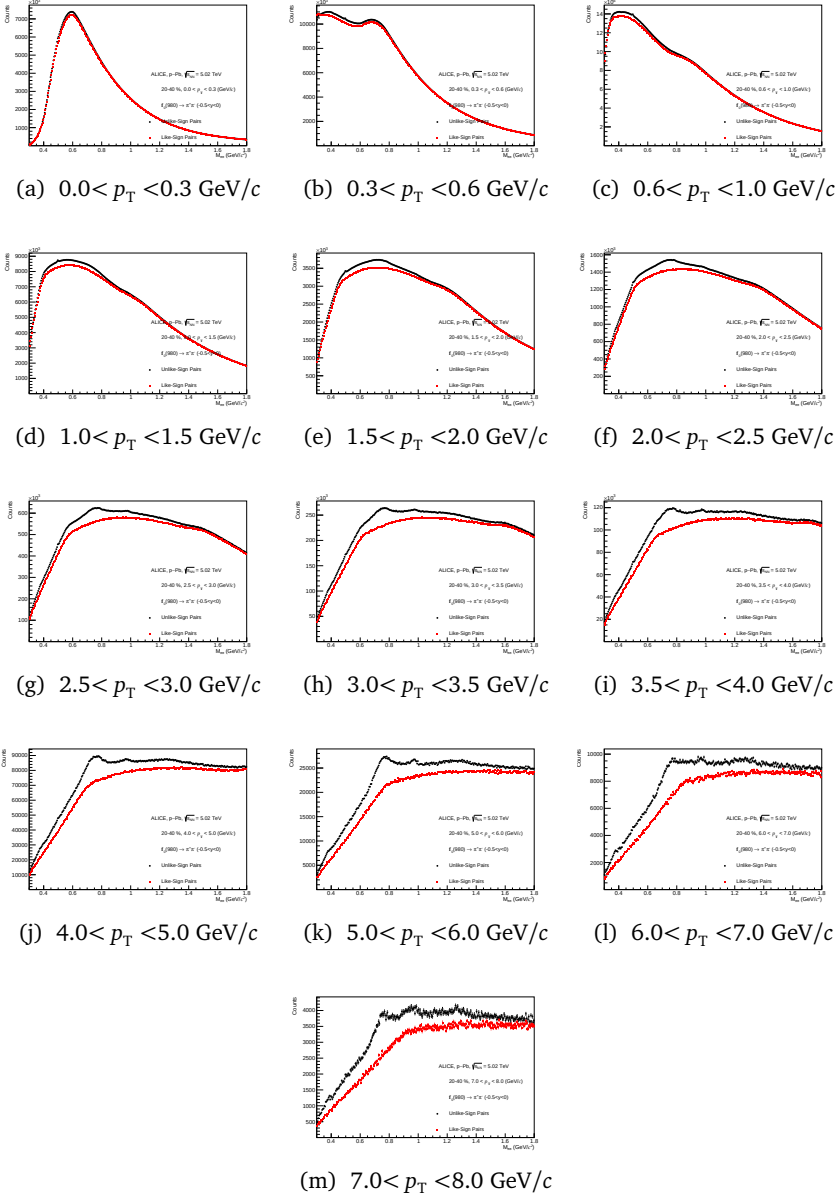


Figure 5.6: Unlike-sign (black) and Like-sign (red) distributions in 20–40 % p–Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$

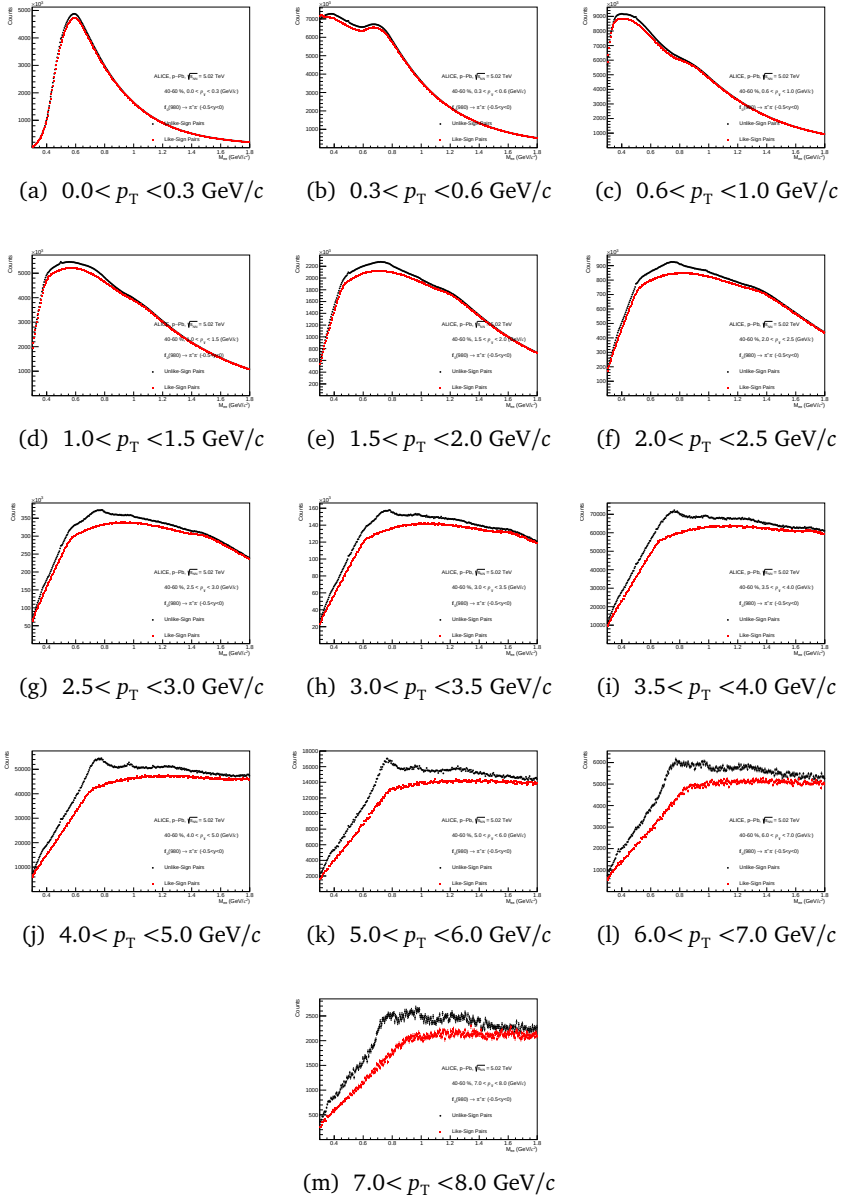


Figure 5.7: Unlike-sign (black) and Like-sign (red) distributions in 40–60 % p–Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$

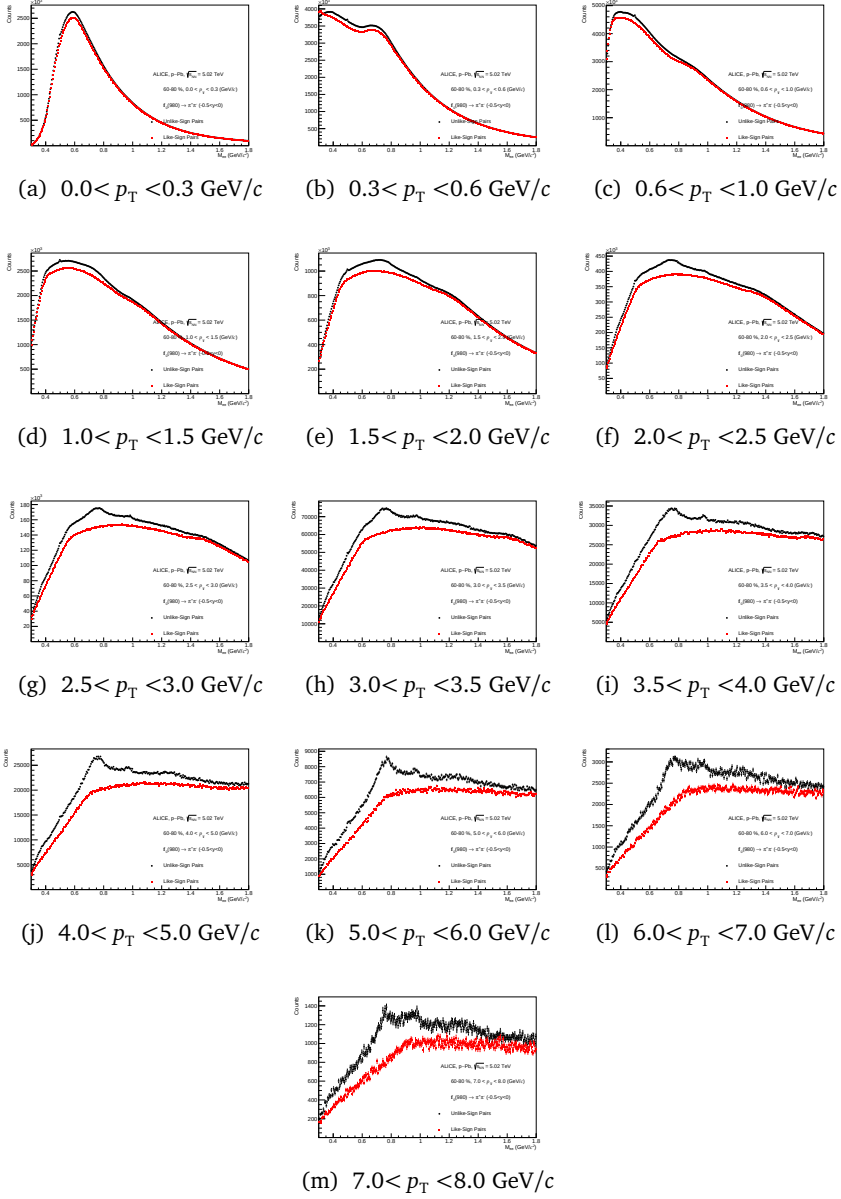


Figure 5.8: Unlike-sign (black) and Like-sign (red) distributions in 60–80 % p–Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$

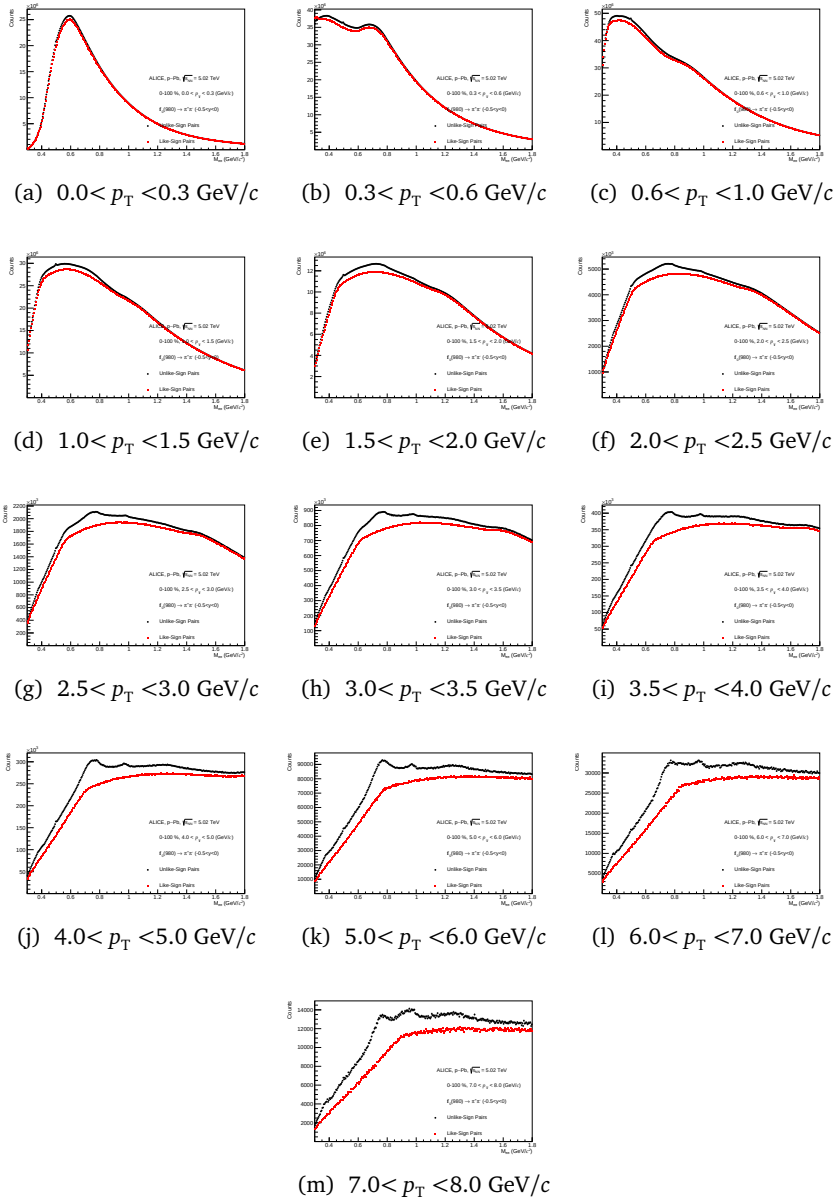


Figure 5.9: Unlike-sign (black) and Like-sign (red) distributions in 0–100 % p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

Figures from 5.5 to 5.9 show Unlike-sign and Like-sign distributions for different multiplicity classes in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV.

5.4 Fit procedure

Many resonances, which have comparable mass with the $f_0(980)$, also decay to $\pi^+\pi^-$, and these resonances can contaminate the signal of $f_0(980)$. Therefore, contributions of other resonances should be considered to exclusively extract the signal of $f_0(980)$. The present analysis mainly considers that contributions of $\rho^0(770)$ and $f_2(1270)$ are dominant resonances contaminating the signal. The contribution of each resonance is represented as relativistic Breit-Wigner function (rBW) [122], where the width (Γ) of rBW can be varying with their reconstructed masses ($M_{\pi\pi}$) and spin numbers (J). The present analysis considers three resonances near the signal region so that three rBW's participate in the construction of the fit function. Each rBW can be expressed as

$$\text{rBW}(M_{\pi\pi}) = \frac{AM_{\pi\pi}\Gamma(M_{\pi\pi})M_0}{(M_{\pi\pi}^2 - M_0^2)^2 + M_0^2\Gamma^2(M_{\pi\pi})}, \quad (5.2)$$

where $\Gamma(M_{\pi\pi})$ is defined as

$$\Gamma(M_{\pi\pi}) = \left[\frac{(M_{\pi\pi}^2 - 4m_\pi^2)}{(M_0^2 - 4m_\pi^2)} \right]^{(2J+1)/2} \times \frac{\Gamma_0 M_0}{M_{\pi\pi}}. \quad (5.3)$$

In equation 5.2 and 5.3, A , m_π , and Γ_0 correspond to the amplitude of rBW, the rest mass of charged pions, and the rest width of the resonance, respectively.

Anticipated residual backgrounds (f_{bkg}) mostly attributed to misidentified particles and mini-jets are fitted with a Maxwell-Boltzmann-like distribution, which can be expressed as [128]

$$f_{\text{bkg}}(M_{\pi\pi}) = B(M_{\pi\pi} - 2m_\pi)^n \exp(c_1 M_{\pi\pi} + c_2 M_{\pi\pi}^2), \quad (5.4)$$

where, B , n , c_1 , and c_2 are free parameters. Each rBW of the resonance is corrected for the $\pi\pi$ interference [129] due to the large phase space for $\pi\pi$ pairs (processes such as $h_1 \rightarrow \pi\pi \rightarrow h_2$), which can be expressed as

$$\text{PS}(M_{\pi\pi}) = \frac{M_{\pi\pi}}{\sqrt{M_{\pi\pi}^2 + p_{\text{T}}^2}} \times \exp(-\sqrt{M_{\pi\pi}^2 + p_{\text{T}}^2}/T_{\text{kin}}), \quad (5.5)$$

where p_T in the above equation denotes the transverse momentum of the $\pi\pi$ pair, and T_{kin} is the kinetic freeze-out temperature and set to be 160 MeV [122] for different multiplicity classes.

The fit procedure is divided into three procedures considering roughly defined width of $f_0(980)$ as $10 < \Gamma_0^{f_0} < 100 \text{ MeV}/c^2$. [1]. The fit function, three rBW's and f_{bkg} , is defined with nine free fit parameters out of fifteen. The T_{kin} and p_T used in equation 5.5 are fixed to 160 MeV and median value in a given p_T range during all three fit procedures, respectively. The masses and widths of $\rho(770)$ and $f_2(1270)$ are well defined in Ref. [1], and those are fixed to $m_\rho = 775.3 \text{ MeV}/c^2$, $\Gamma_\rho = 149.1 \text{ MeV}/c^2$, $m_{f_2} = 1,275.5 \text{ MeV}/c^2$, and $\Gamma_{f_2} = 186.7 \text{ MeV}/c^2$ during all fit procedures. However, nine free parameters are still too many for the fit result to be in local minima. In this respect, additional constraints for fit parameters are considered in each step to minimize the failure of the fit.

The purpose of the first step is to obtain unbiased and initial $f_0(980)$ width. The first step is only performed using minimum-bias events for the merged p_T intervals to be less affected by statistical fluctuations. During the first step, all nine parameters are left free. The initial $f_0(980)$ width is shown in figure 5.10. The second step aims at constraining the f_{bkg} . In the second step, the $f_0(980)$ width is fixed with the value determined in the first step. The last fit procedure is processed with the fixed f_{bkg} , while the $f_0(980)$ width is set free again in the range of $10 < \Gamma_0^{f_0} < 100 \text{ MeV}/c^2$. In each step, parameters for amplitudes of each resonance and the mass of $f_0(980)$ are left free, and the fit range is set to $0.8 < M_{\pi\pi} < 1.76 \text{ GeV}/c^2$.

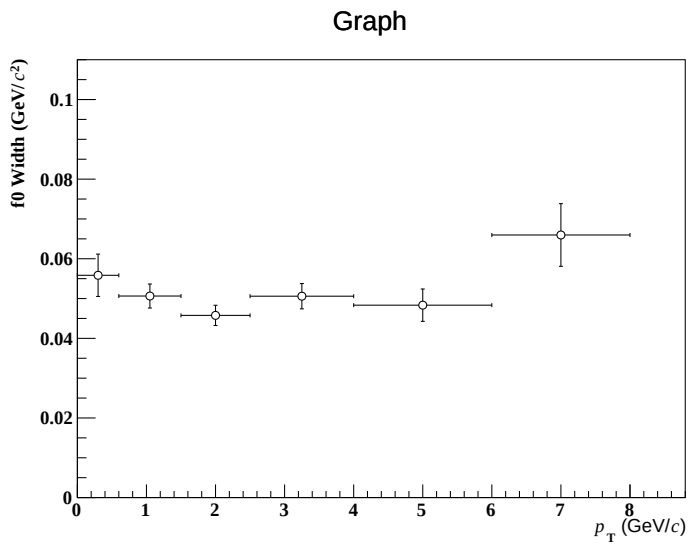


Figure 5.10: Initial $f_0(980)$ width obtained from the first step

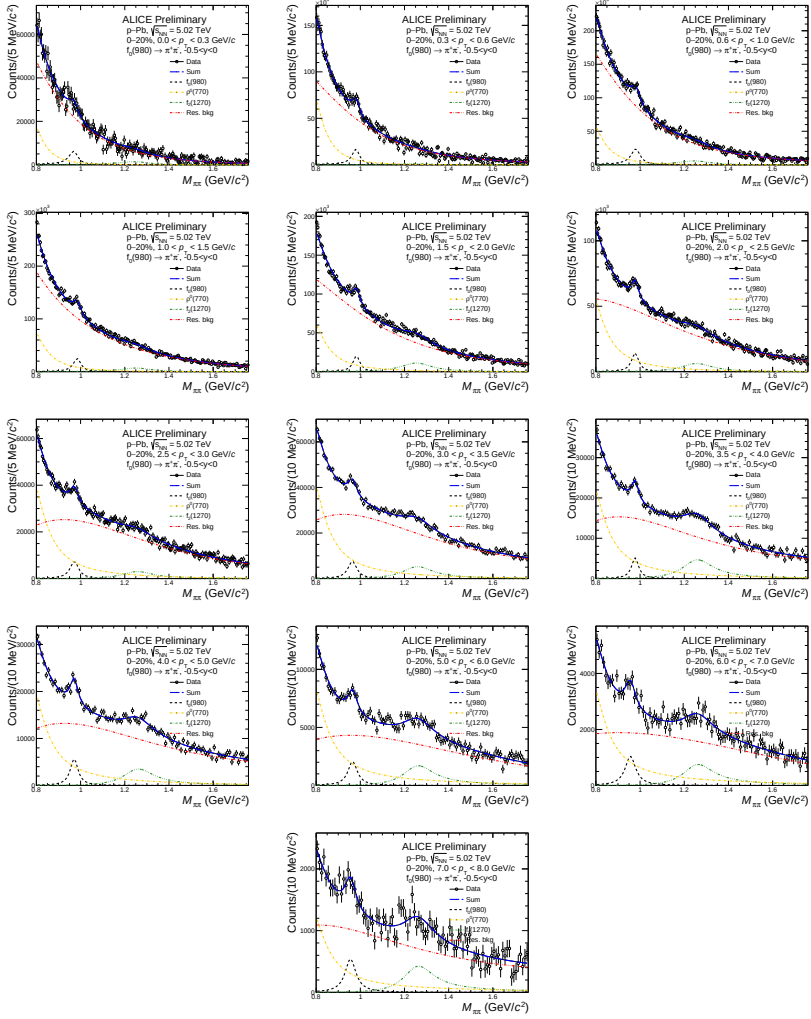


Figure 5.11: Like-sign-subtracted invariant mass distribution in 0–20% p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

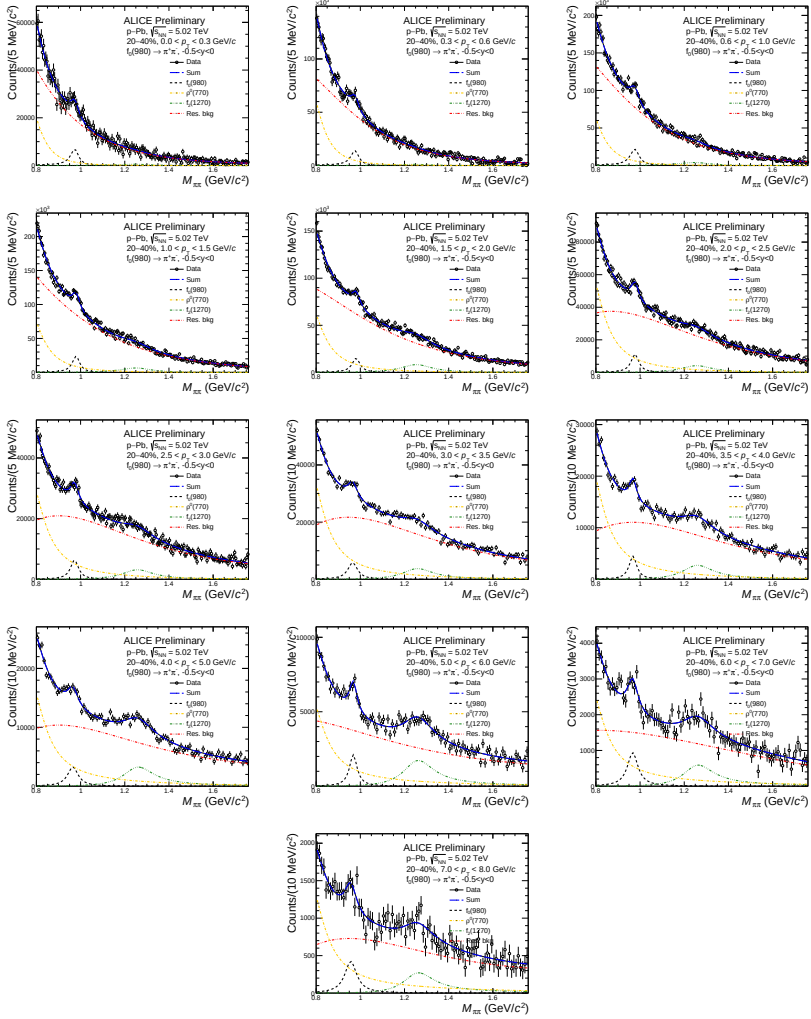


Figure 5.12: Like-sign-subtracted invariant mass distribution in 20–40% p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

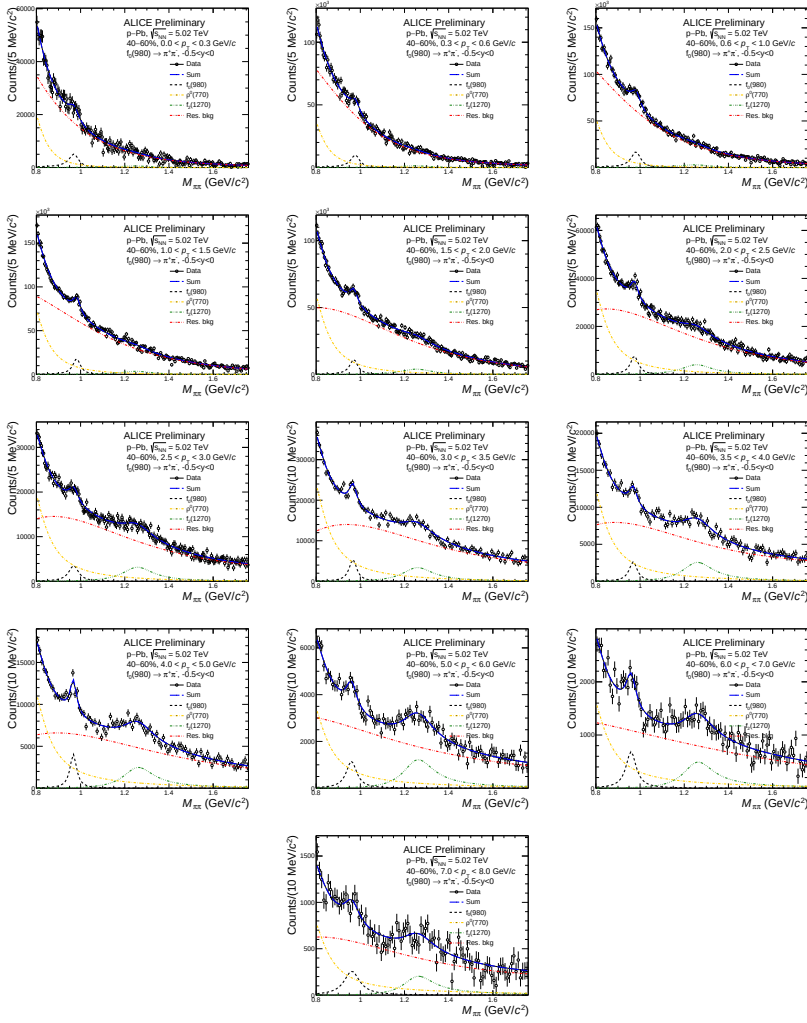


Figure 5.13: Like-sign-subtracted invariant mass distribution in 40–60% p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

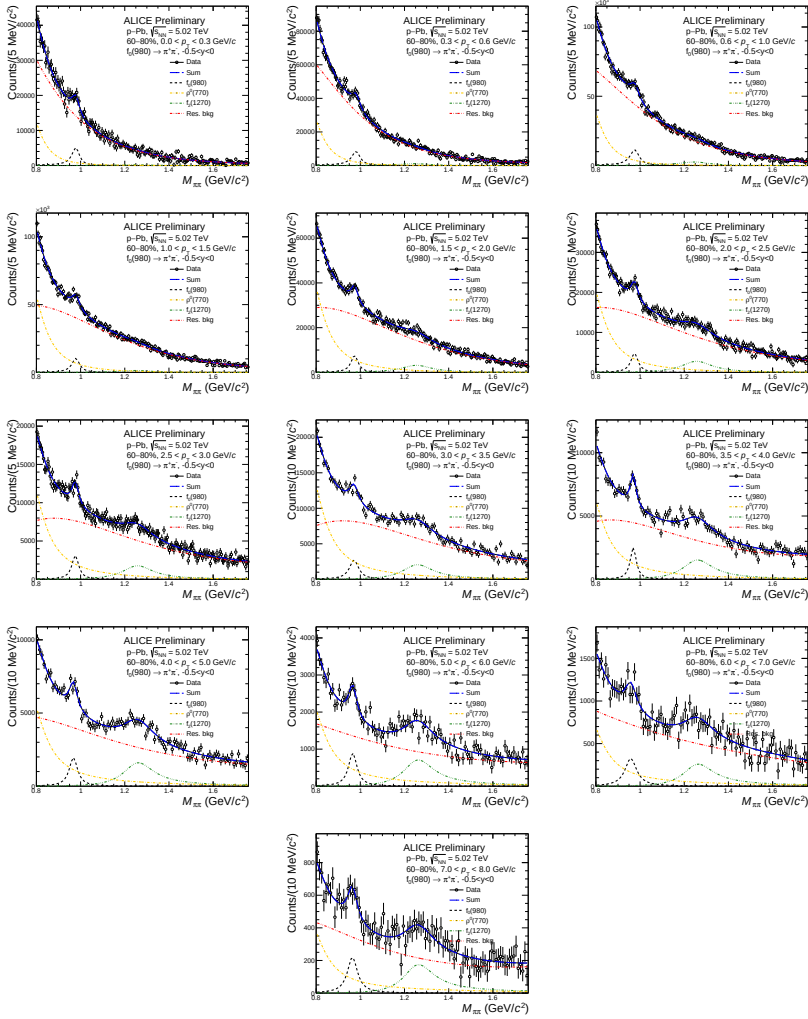


Figure 5.14: Like-sign-subtracted invariant mass distribution in 60–100% p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

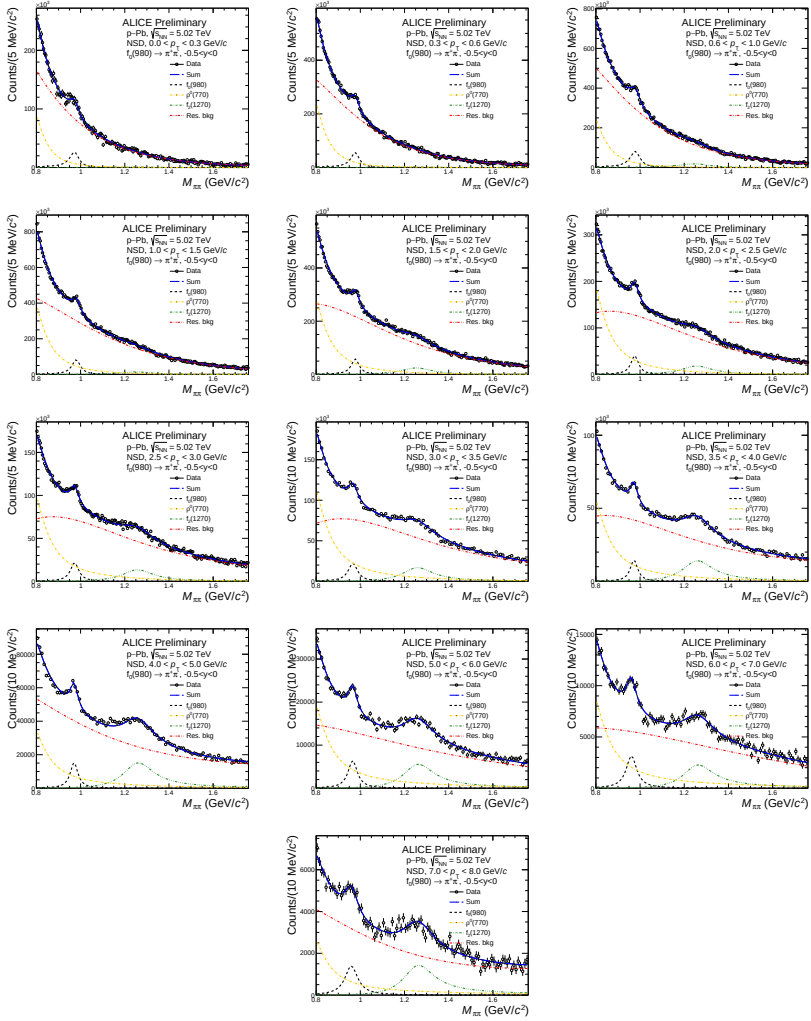


Figure 5.15: Like-sign-subtracted invariant mass distribution in 0–100% p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

Figures from 5.11 to 5.15 show Like-sign-subtracted invariant mass distribution for different multiplicity classes in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The contributions of $f_0(980)$, $\rho^0(770)$, $f_2(1270)$, and f_{bkg} are represented as black lines, yellow lines, green lines, and red lines, respectively. Blue lines show the total fit function to the data, which are shown as open circles.

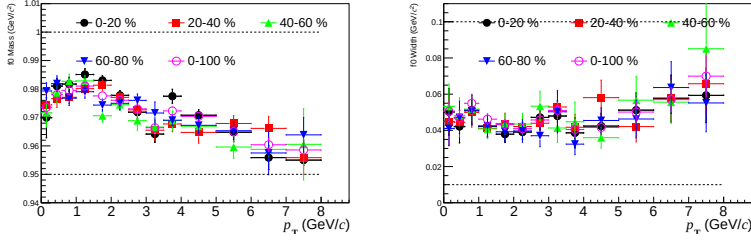


Figure 5.16: Distribution of estimated mass (left) and width (right) of $f_0(980)$ in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, respectively. Dotted lines in the figures denote the upper and lower limit of the fit parameter.

Figure 5.16 shows the finally estimated mass and width of $f_0(980)$ as a function of p_T in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The estimated mass is mildly decreasing with increasing p_T . The estimated width is independent on p_T , exhibiting similar values in figure 5.10.

5.5 Correction

Correction factors are evaluated by analyzing MC data samples anchored to experimental data used in the present analysis. Two types of MC data samples are utilized. General-purposed and Resonance-injected MC data samples are used to estimate normalization factors and the acceptance and efficiency of resonances, respectively. Resonances are injected with a uniform p_T distribution, which will be corrected afterward.

5.5.1 Normalization to NSD

Several corrections are addressed to normalize the number of events and the number of signals to NSD class. The procedure of normalization can be expressed as

$$\frac{Y_{\text{NSD}}}{N_{\text{NSD}}} = \frac{Y_{\text{kINT7 \& vtx}}}{N_{\text{kINT7 \& vtx}}} \times \frac{N_{\text{kINT7}}}{N_{\text{NSD}}} \times \frac{N_{\text{kINT7 \& vtx}}}{N_{\text{kINT7}}} \times \frac{Y_{\text{NSD}}}{Y_{\text{kINT7 \& vtx}}}, \quad (5.6)$$

where Y denotes the number of signals, and N denotes the number of events. The first term on the right side denotes a measured term without normalization. The second term corresponds to trigger efficiency, which corrects for the trigger efficiency. The third term is evaluated by comparing the number of events before and after applying good vertex quality selection as a function of the multiplicity class. The event loss due to vertex quality selection cuts is estimated to be almost unity in the following Ref. [130]. The second and third terms are evaluated in each multiplicity interval defined in the present analysis and are shown in figure 5.17.

5.5.2 Signal-Loss Correction

The last term can be evaluated by comparing the number of generated $f_0(980)$ before and after applying the event selection as a function of transverse momentum and multiplicity class. Due to the insufficient samples for $f_0(980)$ resonances in general-purposed MC data, the correction factor is hardly evaluated. Instead, the coefficients are estimated using other abundant particles, $\rho^0(770)$ in the present analysis, with the assumption that the slopes of $m_T = \sqrt{p_T^2 + m_0^2}$ spectra for entire hadrons are universal [131]. The correction factors are estimated with different

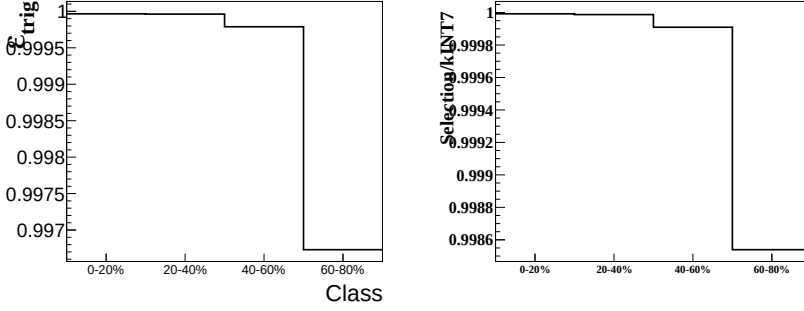


Figure 5.17: Trigger efficiency (left) and event loss due to the vertex quality selection (right) as a function of multiplicity class in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV, respectively.

particles, π , K , and ϕ mesons, to confirm the universality of the correction factor. The estimated correction factors from different particles are identical to one from $\rho^0(770)$ so that the systematic uncertainty from the signal-loss correction is concluded to be negligible. The correction factor of the signal-loss correction, CF can be expressed as

$$CF = \frac{\frac{1}{p_T} \frac{d^2 N_{f_0}^{\text{NSD}}}{dy dp_T}}{\frac{1}{p_T} \frac{d^2 N_{f_0}^{\text{E.S.}}}{dy dp_T}} = \frac{\frac{1}{m_T} \frac{d^2 N_{f_0}^{\text{NSD}}}{dy dm_T}}{\frac{1}{m_T} \frac{d^2 N_{f_0}^{\text{E.S.}}}{dy dm_T}} \sim \frac{\frac{1}{m_T} \frac{d^2 N_h^{\text{NSD}}}{dy dm_T}}{\frac{1}{m_T} \frac{d^2 N_h^{\text{E.S.}}}{dy dm_T}} = \frac{\frac{1}{p_T} \frac{d^2 N_h^{\text{NSD}}}{dy dp_T}}{\frac{1}{p_T} \frac{d^2 N_h^{\text{E.S.}}}{dy dp_T}}, \quad (5.7)$$

where h denotes (any) hadrons tested in the present analysis. Figure 5.18 shows the CF as a function of m_T for different particle species in different multiplicity classes. The correction factor is estimated to be 1% (3%) in low-multiplicity events with the ZDC (VOA) at m_T and becomes unity with increasing multiplicity class and m_T .

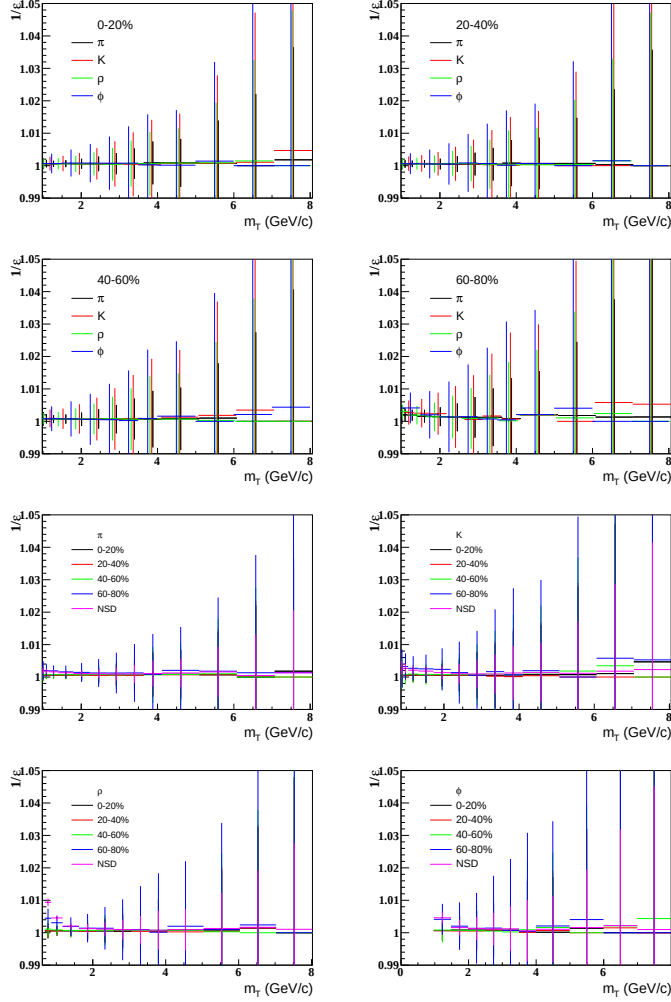


Figure 5.18: Signal loss distribution in multiplicity class for charged pion, kaon, rho meson, and phi meson in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$.

5.5.3 Acceptance $\times$ efficiency

Correction factors for acceptance $\times$ efficiency are estimated by comparing the number of primarily generated and reconstructable $f_0(980)$ resonances, which can be expressed as

$$\text{Acceptance} \times \epsilon = \frac{N_{\text{rec}}^{f_0}}{N_{\text{gen}}^{f_0}}, \quad (5.8)$$

where $N_{\text{rec}}^{f_0}$ and $N_{\text{gen}}^{f_0}$ denote the number of reconstructed and generated $f_0(980)$ resonances, respectively. The multiplicity dependence of the correction factor is tested with the VOA multiplicity instead of the ZDC because simulations or models of low-energy neutrons are not concrete [132]. Figure 5.19 shows the Acceptance $\times$ efficiency as a function of p_T for different multiplicity classes defined by the VOA. The correction factor is 30% at low p_T and gradually increases up to 60% with increasing p_T . The multiplicity dependence of the correction factor is $\pm 10\%$ at low p_T and disappears as p_T increases.

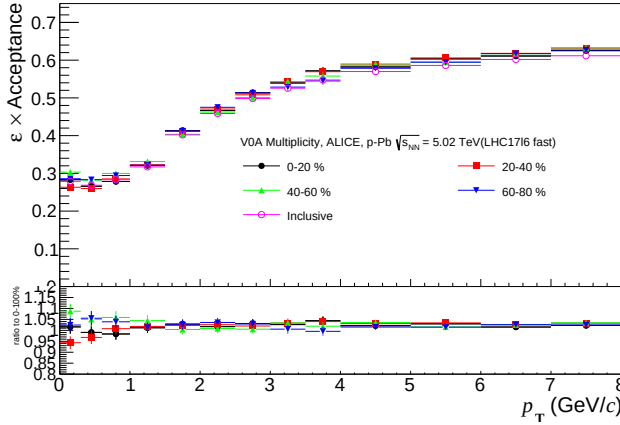


Figure 5.19: Acceptance $\times$ efficiency in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$. Multiplicity dependence was checked with the VOA.

5.5.4 Re-weighting procedure

The $f_0(980)$ resonances are artificially generated with a uniform p_T spectrum, which possibly induces bias. The re-weighting procedure is

applied to the correction factor evaluated in section 5.5.3 to correct for possible bias owing to the uniform spectrum, which re-weights the uniformly generated p_T spectrum with the p_T interval used in the experiment and realistic spectrum obtained in the experiment. The standard process can be found in;

PWGLF/RESONANCES/macros/utils/ReweightEfficiency.C.

Figure 5.20 shows the re-weighted correction factors as a function of p_T for different multiplicity classes. The re-weighting procedure corrects the acceptance $\times$ efficiency up to 2% level.

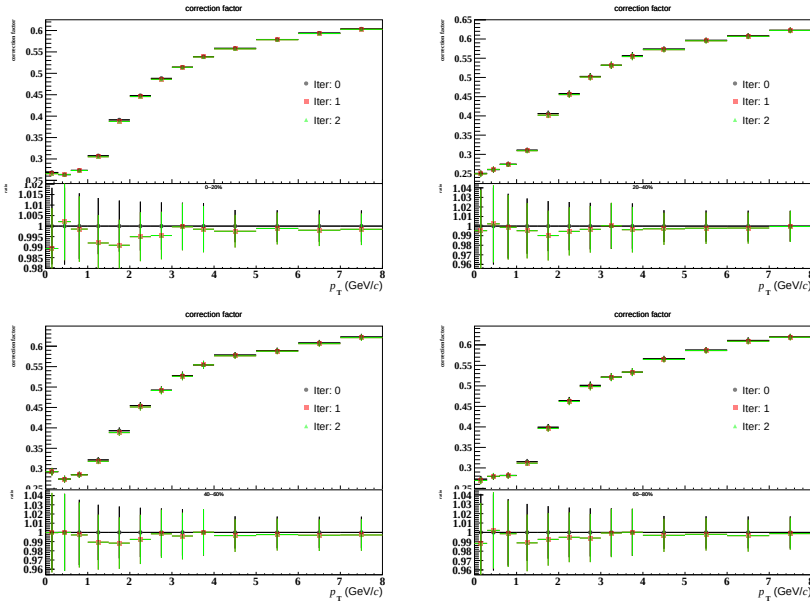


Figure 5.20: Re-weighted acceptance $\times$ efficiency

Chapter 6

Systematic uncertainties

The systematic uncertainties of invariant yields in the present analysis are evaluated by varying the analysis selection criteria and corrections, where each variation is expected to be a systematic source and is conformed with the Barlow test [133]. The Barlow test is performed with the distribution of the Barlow factor (B), which can be expressed as

$$B = \frac{y_{\text{def}} - y_{\text{var}}}{\sqrt{|\sigma_{\text{def}}^2 \pm \sigma_{\text{var}}^2|}}, \quad (6.1)$$

where the y_{def} and y_{var} are observables with the default and variation setups, respectively, and the σ_{def} and σ_{var} denote the statistical uncertainties obtained with the default and variation setups, respectively. The Barlow factor compares the systematic deviation with the statistical fluctuation. If the Barlow factor is smaller than unity, it means the systematic deviation is smaller than the statistical fluctuation, implying that the systematic deviation can be contaminated by the statistical fluctuation and that no significant systematic deviation is expected from the corresponding variation. However, the Barlow factor can also fluctuate itself, and the variation is expected to induce a global deviation. In this respect, the Barlow test is not conducted for each data point but is conducted for the entire data for each variation. The distribution of the Barlow factor is utilized to determine if the variation systematically affects the result. In the

present analysis, the following four requirements on the B distribution are used.

- $-0.8 < \text{Average} < 0.8$
- $\text{Standard Dev.} < 1.3$
- $P(|B| < 1) > 0.55$
- $P(|B| < 2) > 0.75$.

If the B distribution for a given variation meets four requirements, the variation is thought to be consistent with the default setup and not considered to induce systematic uncertainty. If multiple variations are tested for the one possible source, the variations are combined to one systematic uncertainty as the one source by calculating RMS of the variations.

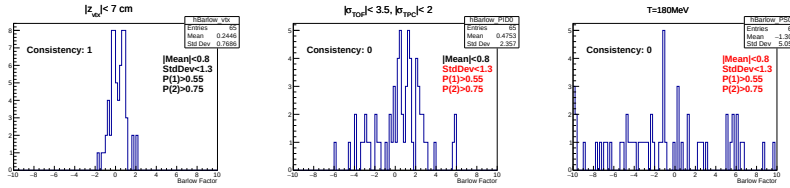


Figure 6.1: Barlow factor distributions from the variation of the primary vertex selection (left), the particle identification (middle), and the phase space correction (right)

Figure 6.1 shows the Barlow factor distributions from different variations. The variation of the primary vertex selection is considered as not a systematic source because the B distribution satisfies four requirements. On the other hand, the B distributions from the particle identification and the phase space correction do not meet four requirements, so they are considered as the systematic source.

The smoothing procedure is applied to the calculated RMS with the assumption that the systematic uncertainty is not discontinuous between neighboring p_T intervals. The smoothing procedure can be described as

$$y_i^{\text{smooth}} = \frac{y_{i-1} + y_i + y_{i+1}}{3}, \quad (6.2)$$

where the smoothing procedure is not applied to the first and last bins.

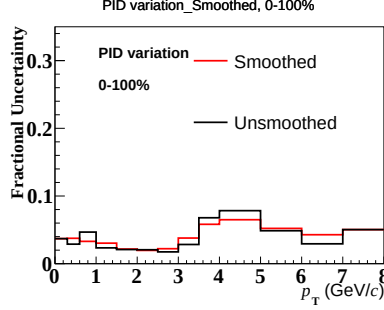


Figure 6.2: Evaluated systematic uncertainty and the uncertainty after applying the smoothing procedure.

Figure 6.2 shows the evaluated systematic uncertainty from the variation of the particle identification and the result of the smooth procedure for the uncertainty.

The R method [9] is used to estimate how much uncertainties are correlated with N_{ch} . The relative variations between the specific multiplicity class and minimum bias are calculated as

$$R_{\text{var}} = \frac{\frac{S_{\text{var}}^j}{S_{\text{def}}^j}}{\frac{S_{\text{var}}^{\text{MB}}}{S_{\text{def}}^{\text{MB}}}} \quad (6.3)$$

The R_{var} is calculated in different p_{T} , and total uncorrelated uncertainty can be expressed as

$$\left(\frac{\delta S^j}{S^j} \right)^2 = \sum_{\text{var}} (R_{\text{var}} - 1)^2 \quad (6.4)$$

The systematic uncertainties of invariant yields with the above procedures are summarized in table 6.1. The total systematic uncertainty is calculated as the sum in quadrature of the different contributing sources. Each uncertainty does not significantly depend on the chosen p_{T} range and the multiplicity class.

The systematic uncertainty from the primary vertex selection is inspected with the Barlow test by narrowing the requirement to 7 cm, and

Table 6.1: The relative systematic uncertainty of invariant p_T -differential yields. Numbers given in ranges correspond to minimum and maximum uncertainties.

Sources	Systematic uncertainty (%)
Primary vertex	negligible
Pileup rejection	negligible
Tracking	± 4 –6
Particle identification	± 4 –12
$f_2(1270)$ parameters	± 3 –9
Signal extraction $\rho(770)$ parameters	± 3 –8
Fit range	± 0 –6
Initial f_0 width	± 2 –12
Phase space correction	± 3 –8
Total (in quadrature)	± 15 –27

the uncertainty is estimated to be negligible. The systematic uncertainty from the pileup rejection is tested with the Barlow test by varying the minimal number of track segments contributing to the reconstruction of pileup event vertices from 5 to 3. The uncertainty is estimated to be also negligible.

The systematic uncertainty from tracking is assigned as the value obtained from [126], where the uncertainties are estimated by varying the requirements to select reconstructed tracks such as d_{xy} , d_z , and the number of contributing TPC readout channels. The systematic uncertainty from the particle identification is inspected with different requirements on the number of standard deviations ($\pm 0.5\sigma$) for the TPC and TOF. The uncertainties are evaluated to be 4–12%.

The systematic uncertainties from possible deviations of masses and widths of $f_2(1270)$ and $\rho(770)$ are evaluated by shifting the masses and the widths three times within their uncertainties, which are reported at Ref. [1]. The evaluated uncertainties from $f_2(1270)$ and $\rho(770)$ parameters are 3–9% and 3–8%, respectively. The systematic uncertainty from the fit range is evaluated by changing the range inward or outward as much as 40 MeV/ c^2 . The uncertainties are estimated to be less than 6%.

The systematic uncertainty from the initial value for $f_0(980)$ width,

which is obtained in the first step described in section ??, is estimated by varying the width within the statistical uncertainties in both directions as shown in figure 5.10. The changes are expected to affect the background distribution determined in the second step, and the estimated systematic uncertainties are 2–12%. The systematic uncertainty from the phase space correction is estimated by varying the kinetic freeze-out temperature (T_{kin}) in the range of $140 < T_{\text{kin}} < 180$ MeV. The evaluated uncertainties are 3–8%.

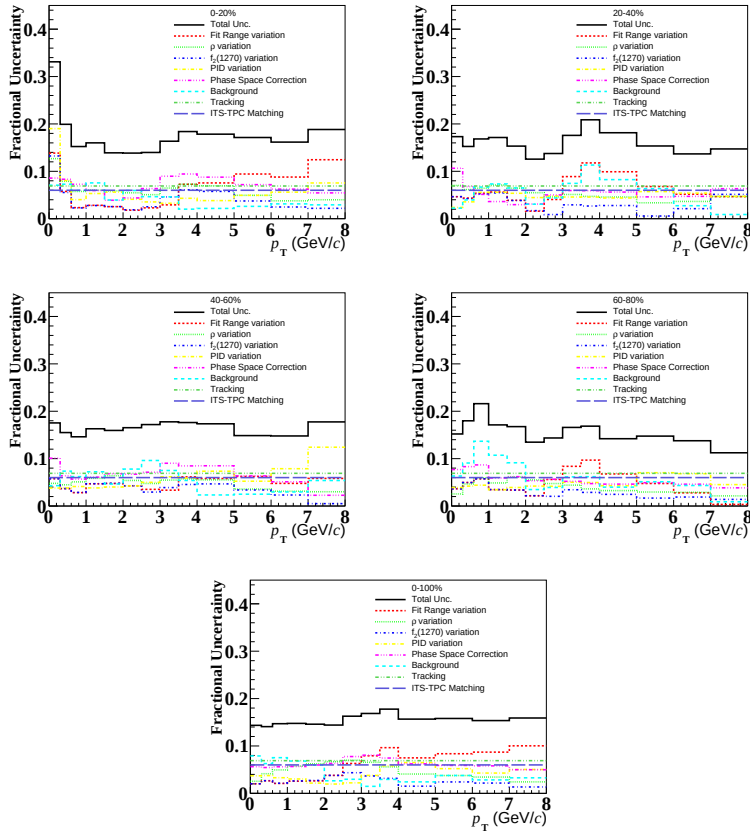


Figure 6.3: Total systematic uncertainty.

Figure 6.3 and figure 6.4 show the total systematic uncertainty and total uncorrelated systematic uncertainty, respectively. The total systematic uncertainty is evaluated to be 15–27%, and the total uncorrelated

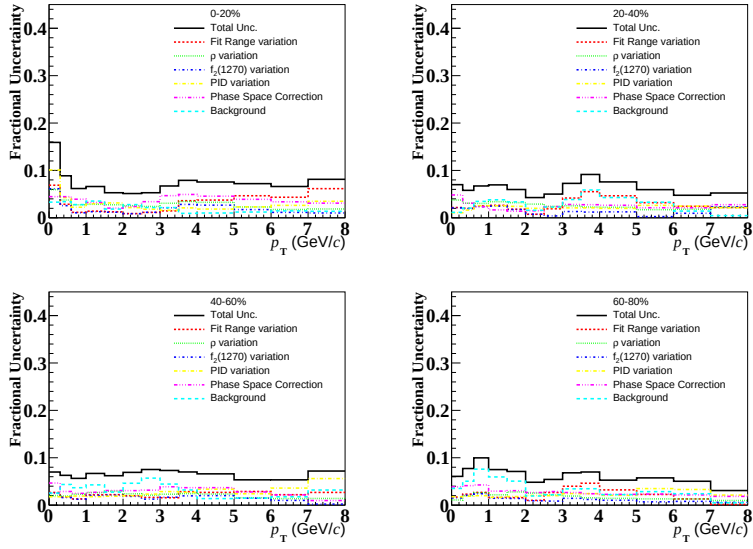


Figure 6.4: Total uncorrelated systematic uncertainty.

systematic uncertainty is evaluated to be 7-10%.

Chapter 7

Results

7.1 $f_0(980)$ in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV

The upper panel of figure 7.1 shows the p_{T} spectra of $f_0(980)$ in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV measured in the range of $0 < p_{\text{T}} < 8$ GeV/ c for different multiplicity classes, which are determined by VOA, and non-single diffractive (NSD) events at Pb-going midrapidity ($-0.5 < y < 0$). Each spectrum with a different color is scaled with the number represented on the legend in the figure for better visibility. The lower panel of figure 7.1 shows the ratios of the specific p_{T} spectrum to the NSD spectrum. The systematic uncertainties of the ratios are evaluated by propagating the uncorrelated uncertainties along the multiplicity class on each p_{T} spectrum. In low $p_{\text{T}} < 4$ GeV/ c , the hardening of the spectrum from low- to high-multiplicity events is visible, while the same shape is seen at high $p_{\text{T}} > 4$ GeV/ c . Such a trend can be generally observed in different particle species [134].

Table 7.1 shows the integrated yield (dN/dy) and mean transverse momentum ($\langle p_{\text{T}} \rangle$) of $f_0(980)$ for different multiplicity classes in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The dN/dy linearly increases with the multiplicity, while the $\langle p_{\text{T}} \rangle$ weakly increases with the multiplicity.

Three different particles are selected to be shown together with

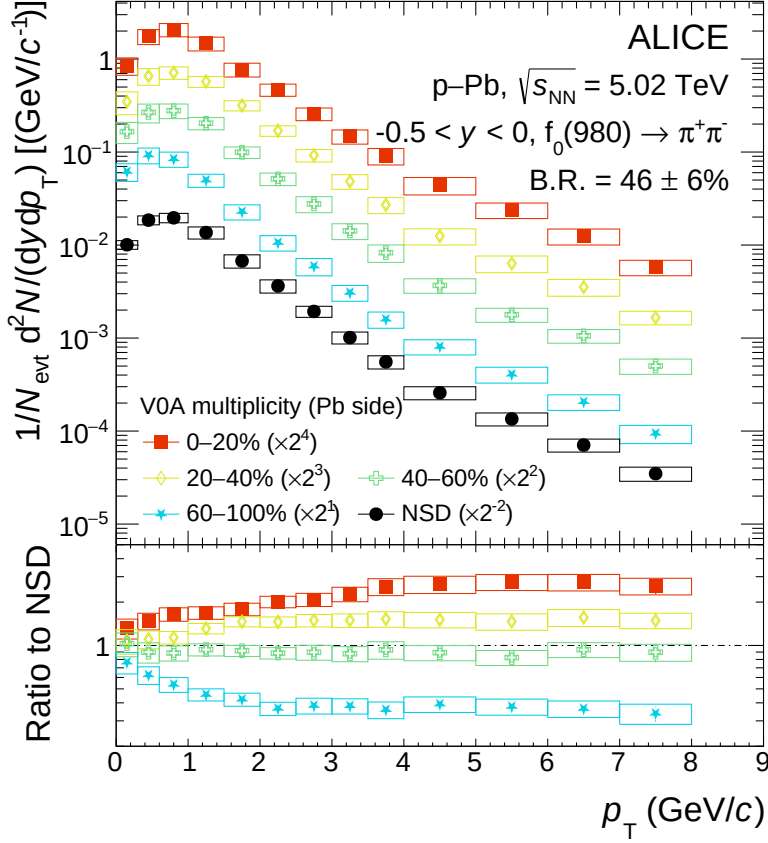


Figure 7.1: The transverse momentum spectra for $f_0(980)$ in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Each color denotes a different multiplicity class, scaled for better visibility. Systematic and statistical uncertainties are shown as boxes and error bars, respectively. The ratios of the specific spectra to the NSD spectrum are shown in the lower panel.

the measurement of $f_0(980)$ production because they can be classified according to their lifetimes and whether they contain (anti-) strange quarks. The charged pion has no strangeness, and its yield is not expected to be largely affected by interactions with the hadron gas. Therefore, measured yields of charged pions are a good reference to study mechanisms suppressing or enhancing the measured particle yields. The $K^{*0}(892)$ has one strangeness and a short lifetime of 4.2 fm/c, which

Table 7.1: The values of dN/dy and mean transverse momentum ($\langle p_T \rangle$) measured in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV for different multiplicity classes. The first and second uncertainties represent the statistical and systematic uncertainties, respectively.

Multiplicity class (V0A)	dN/dy	$\langle p_T \rangle$ (GeV/c)
0–20%	$0.206 \pm 0.005 \pm 0.014$	$1.287 \pm 0.034 \pm 0.010$
20–40%	$0.153 \pm 0.004 \pm 0.010$	$1.250 \pm 0.029 \pm 0.082$
40–60%	$0.113 \pm 0.002 \pm 0.008$	$1.142 \pm 0.025 \pm 0.088$
60–100%	$0.064 \pm 0.001 \pm 0.005$	$0.999 \pm 0.014 \pm 0.080$

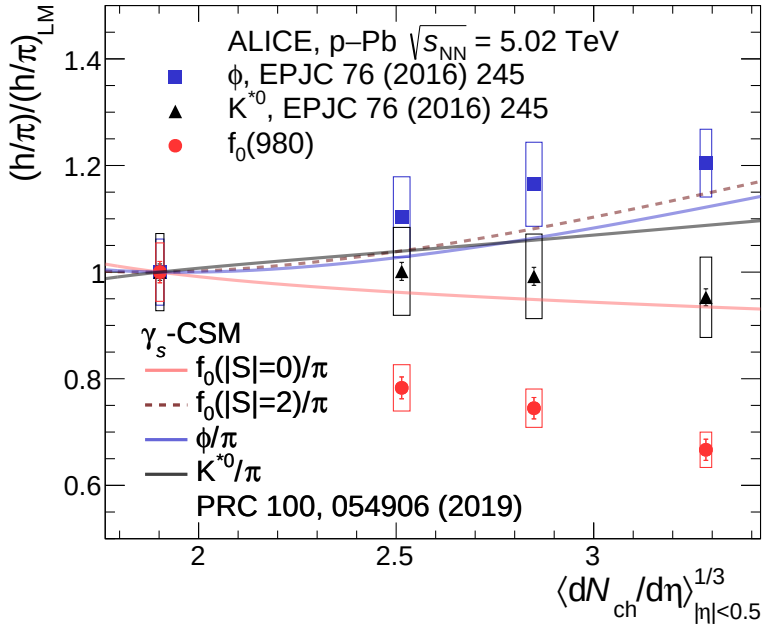


Figure 7.2: Double ratios of ϕ , $K^{*0}(892)$, and $f_0(980)$ to π as a function of charged particle multiplicity elevated to 1/3 in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The ratios are divided by the ratio in low-multiplicity (LM) events to make the first point unity and reduce systematic uncertainties by canceling out the correlated uncertainties. Predictions from the Canonical statistical model [36] are represented with lines.

is comparable with the duration time of the hadronic phase, enabling the decay products of K^{*0} to interact with the hadron gas. On the other hand, the ϕ meson has strange and anti-strange quarks and a long lifetime of 46.2 fm/c, letting the ϕ meson only experiences the strangeness enhancement.

In this respect, figure 7.2 shows the double ratios of different particles to the charged pion yields as a function of charged particle multiplicity elevated to 1/3 in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The measured charged particle multiplicity is classified using the V0A detector. The systematic uncertainties of the double ratios are evaluated using the uncorrelated uncertainties. The particle yield ratio of the ϕ meson to the charged pion is clearly increasing with increasing multiplicity due to the strangeness enhancement for the ϕ meson. On the other hand, the ratio of the K^{*0} to the charged pion is consistent with increasing multiplicity despite one strangeness for the K^{*0} . This trend is attributed to the two competing effects between the strangeness enhancement and the rescattering effects in small collision systems. Finally, the particle yield ratio of $f_0(980)$ to the charged pion is clearly decreasing with increasing multiplicity. The decreasing ratio implies that significant rescattering effects are expected for $f_0(980)$.

The particle yield ratio is calculated using the γ_s -Canonical Statistical Model (CSM) [36], which constructs system-size-dependent hadrochemistry at vanishing baryon density with local conservations of electric charge, baryon density, and strangeness while allowing for undersaturation of strangeness especially in small collision systems. Note that γ_s is the parameter [135] for the undersaturation of strangeness derived from measured particle yields. CSM nicely reproduces the increasing trend of the ϕ meson to the charged pion yields with increasing multiplicity, indicating that CSM can successfully describe the strangeness enhancement. However, CSM largely overestimates the yield ratio of the K^{*0} to charged pion. Such an overestimation is largely attributed to the no consideration of interactions between the hadron gas and the decay products of K^{*0} while the strangeness enhancement for the K^{*0} is well reflected, hence, resulting in the overestimation of the particle yield ratio.

Finally, CSM model calculates the particle yield ratios of $f_0(980)$ to the charged pion with two assumptions for the hidden strangeness

of $f_0(980)$. The hidden strangeness ($|S|$) represents the number of strange quarks and anti-strange quarks. The $\rho^0(770)$ and ϕ have zero and two hidden strangeness, respectively, while both particles have zero strangeness. The first assumption of $|S| = 2$ results in the increasing trend of the ratio, which is the opposite trend of experimental data because CSM only considers the strangeness enhancement for $f_0(980)$ and no rescattering effects with the $|S| = 2$ assumption. The second assumption of $|S| = 0$ results in the consistent trend of the ratio. With the $|S| = 0$ assumption, the strangeness enhancement is not considered together with no rescattering effects for $f_0(980)$. In conclusion, CSM calculations with both assumptions overestimate the ratio of $f_0(980)$ to the charged pion yields because of no rescattering effects.

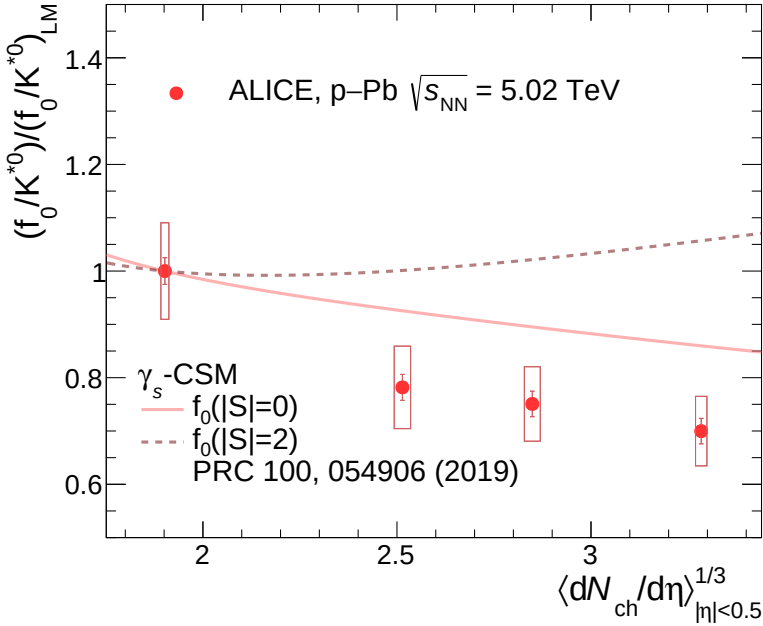


Figure 7.3: Double ratio of $f_0(980)$ to $K^{*0}(892)$ as a function of charged-particle multiplicity elevated to 1/3 in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The ratios are divided by the ratio in low-multiplicity (LM) events to make the first point unity and reduce systematic uncertainties by canceling out the correlated uncertainties. Predictions from the Canonical statistical model [36] are represented with lines.

From the previous exploration, both processes of the strangeness enhancement and the rescattering effects are important to predict particle productions. Taking a yield ratio between two particles experiencing similar processes helps emphasize the specific process that only one of the two particles experiences. Figure 7.3 achieves this by measuring the particle yield ratio of $f_0(980)$ to K^{*0} . Because both particle yields are expected to be largely affected by the rescattering effects, the ratio of $f_0(980)$ to K^{*0} is expected to be less affected by the rescattering effects (as their lifetimes are expected to be similar as well), and the ratio is clearly decreasing with increasing multiplicity. The decreasing trend is qualitatively predicted using CSM model with the $|S| = 0$ assumption. Because CSM only considers the strangeness enhancement for K^{*0} and no rescattering effects for both particles, the ratio is expected to be decreasing with multiplicity. However, CSM estimates the ratio to be increasing with the $|S| = 2$ assumption because the strangeness enhancement is more significantly there for $f_0(980)$ than K^{*0} , indicating no strange quarks in $f_0(980)$. It is worth mentioning that the absence of strange quarks in $f_0(980)$ is also contradictory with the tetraquark suggestion for the $f_0(980)$ state [27].

The upper panel of figure 7.4 shows the p_T -differential particle yield ratio of $f_0(980)$ to the charged pion in high-multiplicity and low-multiplicity p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The ratio is increasing with increasing p_T for both event classes. The ratios are consistent within one sigma at high $p_T > 4$ GeV/c. The lower panel of figure 7.4 shows p_T -differential double ratio of particle yields of $f_0(980)$ to the charged pion. The systematic uncertainties are calculated by propagating only uncorrelated uncertainties. The double ratio proves the consistency of ratios at high p_T . On the other hand, the double ratio begins to be suppressed with decreasing p_T , indicating that the suppression of $f_0(980)$ over the charged pion ratio is dominantly happening at low p_T . The multiplicity-dependent suppression of $f_0(980)$ only at low p_T suggests the rescattering effects for $f_0(980)$.

The upper panel of figure 7.5 shows the p_T -differential particle yield ratio of $f_0(980)$ to the K^{*0} in high-multiplicity and low-multiplicity p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The ratio in high-multiplicity events is systematically lower than that in low-multiplicity events for the entire p_T

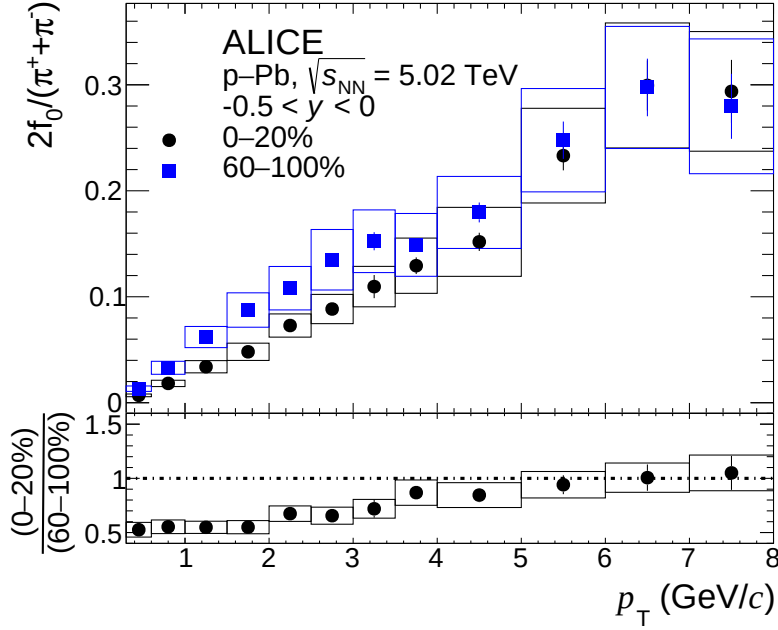


Figure 7.4: Particle yield ratios of $f_0(980)$ to π as a function of p_T in high-multiplicity (circles) and low-multiplicity (triangles) p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The lower panel shows the double ratio of $f_0(980)/\pi$ between the high-multiplicity and low-multiplicity events.

range. The lower panel of figure 7.5 shows p_T -differential double ratio of particle yields of $f_0(980)$ to the K^{*0} . The systematic uncertainties are calculated by propagating only uncorrelated uncertainties. The double ratio is systematically lower than unity in the entire p_T range. The different p_T dependence of the double ratio from figure 7.4, K^{*0}/K [9], and ρ/π [122] indicates that other effects, rather than the rescattering effects, are present. For instance, the enhancement of K^{*0} yield due to the strangeness enhancement can explain the suppression in the entire p_T range, while no strangeness enhancement for $f_0(980)$ yield. As a result, the decreasing trend of the ratio in higher multiplicity events suggests no strange quarks in $f_0(980)$.

Figure 7.6 shows the nuclear modification factor (Q_{pPb}) of $f_0(980)$ as a function of p_T in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV for different multiplicity classes, which is defined using the ZDC detector and the hy-

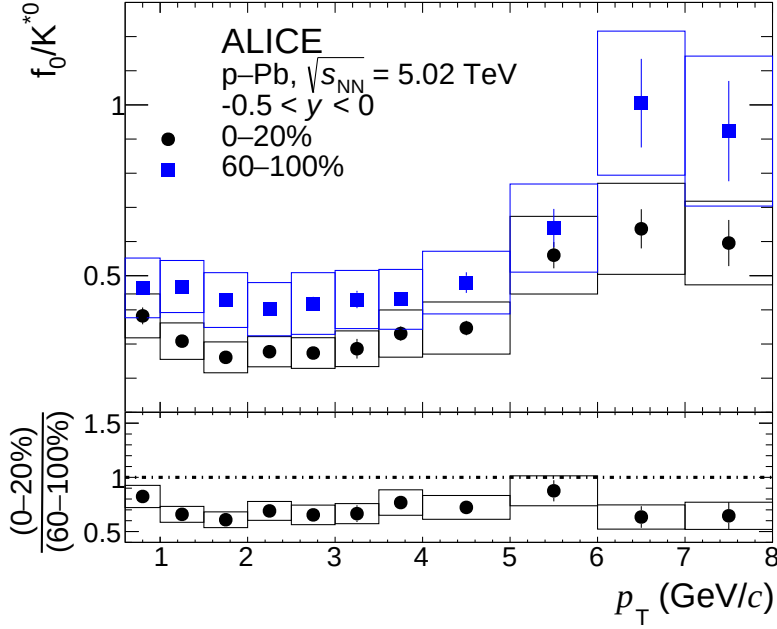


Figure 7.5: Particle yield ratios of $f_0(980)$ to K^{*0} as a function of p_T in high-multiplicity (circles) and low-multiplicity (triangles) p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The lower panel shows the double ratio of $f_0(980)/K^{*0}$ between the high-multiplicity and low-multiplicity events.

brid method for the scaling with the number of binary collisions (see section 4.3.2 for details). The branching ratio of $f_0(980) \rightarrow \pi^+ \pi^-$ is canceled out as the same value is used for two different analyses. The systematic uncertainties are calculated with the assumption that there is no correlated uncertainty between the yield in pp and p-Pb collisions. The Q_{pPb} of $f_0(980)$ exhibits stronger suppressions in low p_T compared with Q_{pPb} of charged hadrons, and the suppression increases with increasing multiplicity. Such a multiplicity-dependent suppression only at low p_T indicates the confirmation of the rescattering effects. In addition, the Q_{pPb} does not exhibit the Cronin-like enhancement [136] at the intermediate p_T in HM events. This is a consistent trend of conventional mesons, as shown in the right panel of figure 2.9. No Cronin-like enhancement for $f_0(980)$ might suggest that the number of constituent quarks of $f_0(980)$ is smaller than three.

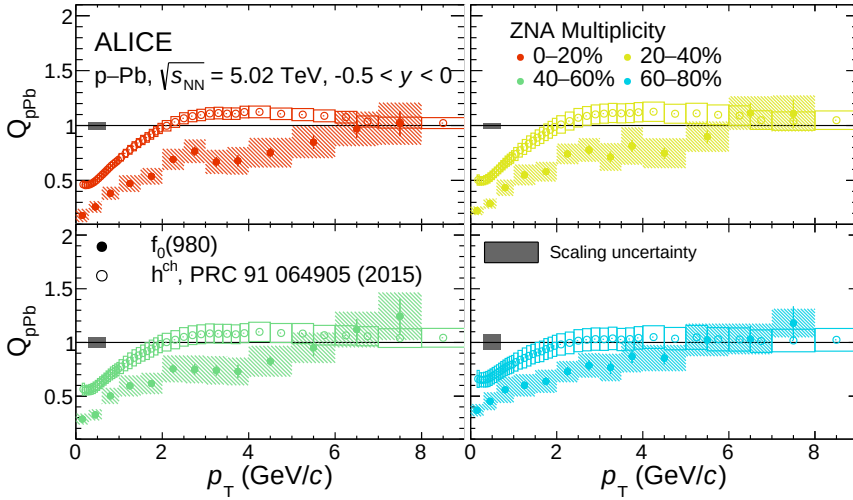


Figure 7.6: Nuclear modification factor (Q_{pPb}) of $f_0(980)$ as a function of p_T in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV for different multiplicity classes. Statistical and systematic uncertainties are shown as error bars and boxes, respectively. Open boxes around $p_T = 0.5$ GeV/ c represent the binary collision scaling uncertainties. The Q_{pPb} of $f_0(980)$ is compared with Q_{pPb} of charged hadrons.

7.2 Extension to different collision systems

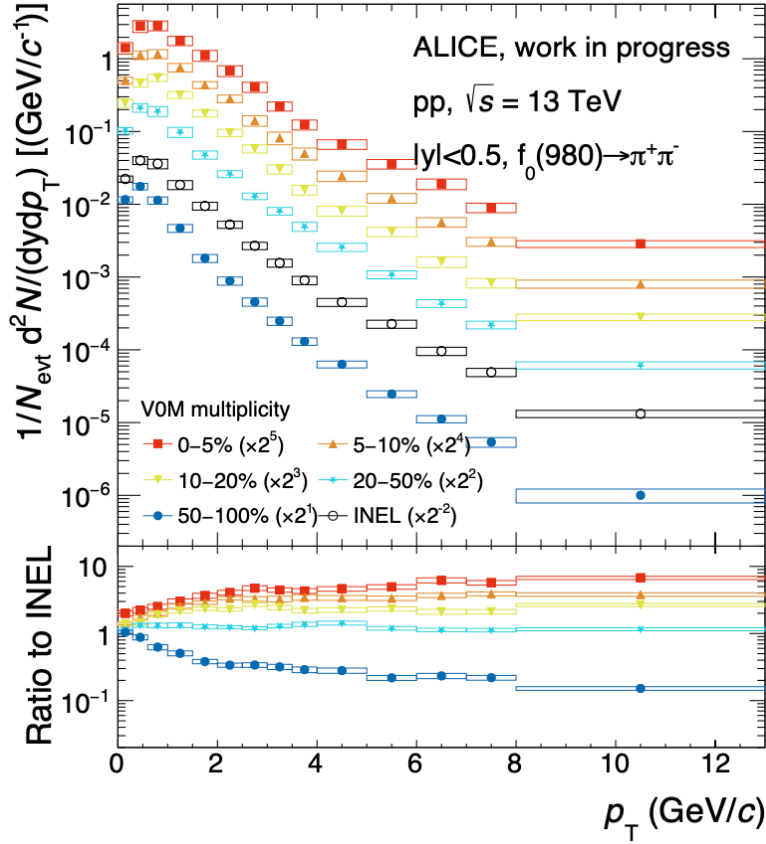


Figure 7.7: The transverse momentum spectra for $f_0(980)$ in pp collisions at $\sqrt{s} = 13$ TeV. Each color denotes a different multiplicity class, which is scaled for better visibility. Systematic and statistical uncertainties are shown as boxes and error bars, respectively. The ratios of the specific spectra to the INEL spectrum is shown in the lower panel.

Figure 7.7 shows the p_T spectra of $f_0(980)$ in pp collisions at $\sqrt{s} = 13$ TeV for different multiplicity classes measured in the range of $0 < p_T < 13$ GeV/c. Each spectrum is scaled with the number denoted in the legend for better visibility. The spectra are divided by the p_T spectrum obtained in the INEL event class. The systematic uncertainties are

calculated using the uncorrelated uncertainties between two spectra. At low p_T , the hardening of the spectrum is observed, and the hardening ceases at high p_T , exhibiting similar multiplicity and p_T dependencies with figure 7.1.

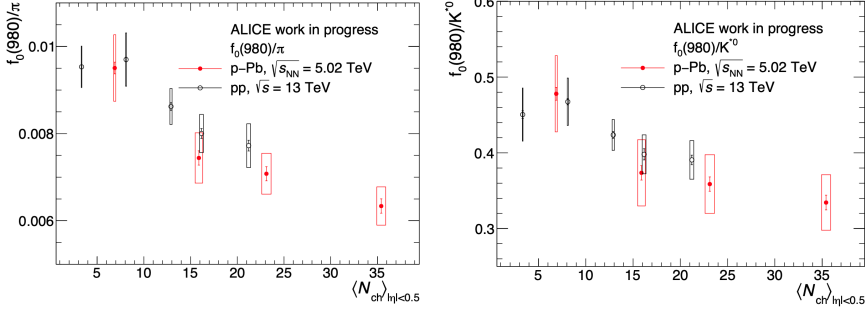


Figure 7.8: The particle yield ratio of $f_0(980)$ to the charged pion (left) and the K^{*0} (right) as a function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV and p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV.

Figure 7.8 shows the particle yield ratio of $f_0(980)$ to the charged pions (left) and the K^{*0} (right) as a function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV and p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Sequential suppressions for both ratios are observed, exhibiting a similar trend to those measured in p-Pb collisions.

Figure 7.9 shows the p_T spectra of $f_0(980)$ in different centrality Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV in the range of $0 < p_T < 8$ GeV/ c . The measurement of the p_T spectrum is conducted only with statistical uncertainties. Two different triggers are used to collect data in Pb-Pb collisions. In the measurement of the p_T spectrum, dedicated triggers for 0–10% and 30–50% classes are used together with the minimum-bias trigger.

Figure 7.10 shows the p_T -integrated yield (dN/dy , left) and mean transverse momentum ($\langle p_T \rangle$, right) as a function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV, p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The integrated yield is linearly increasing with increasing multiplicity. The mean transverse momentum is suppressed for heavier collision systems at the same multiplicity, which can be generally observed in different particle

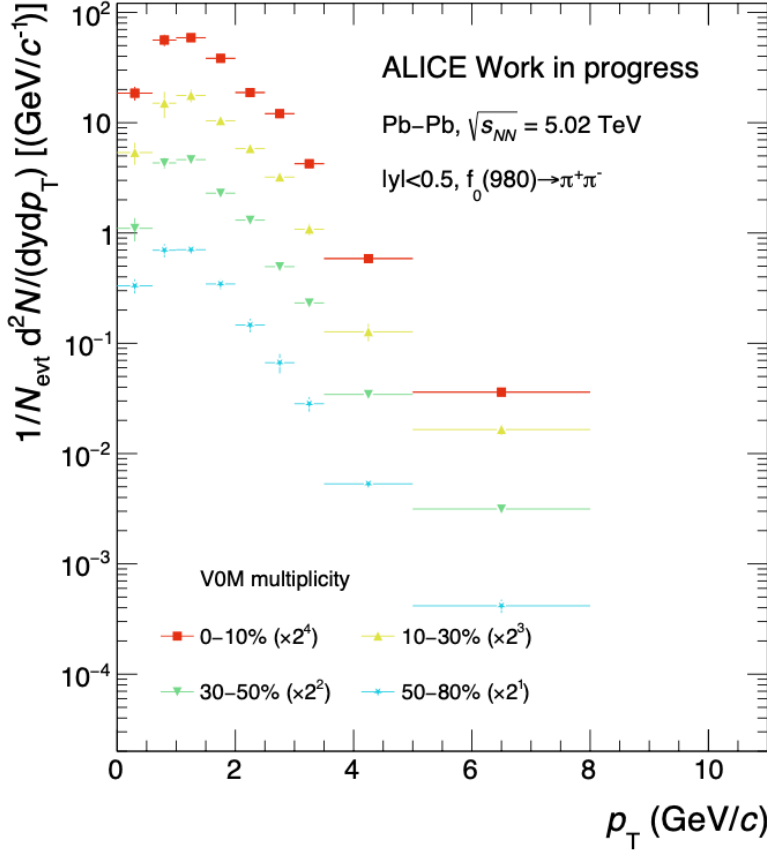


Figure 7.9: The transverse momentum spectra for $f_0(980)$ in Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Each color denotes a different multiplicity class, which is scaled for better visibility. Statistical uncertainties are shown as error bars.

species [137].

Figure 7.11 shows the particle yield ratio of $f_0(980)$ to the charged pion as a function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV, p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Sequential suppression from pp collision to Pb–Pb collisions is visible, which is visible for different resonances such as K^{*0} [9], ρ [122], and so on [138].

Figure 7.12 shows the particle yield ratio of $f_0(980)$ to the K^{*0} as a

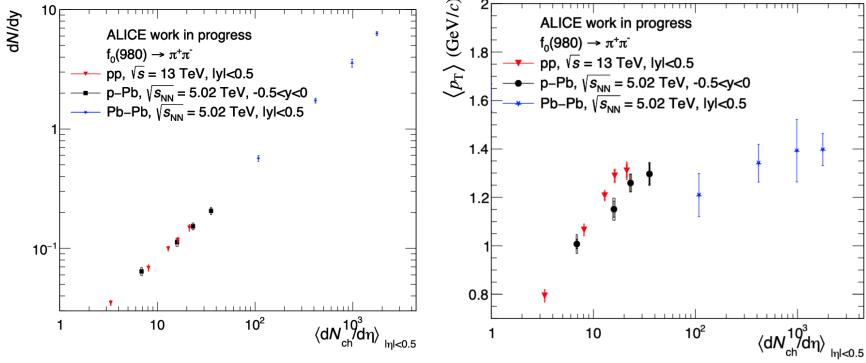


Figure 7.10: Integrated yield (left) and mean transverse momentum (right) of $f_0(980) \rightarrow \pi^+\pi^-$ from pp collisions to Pb–Pb collisions.

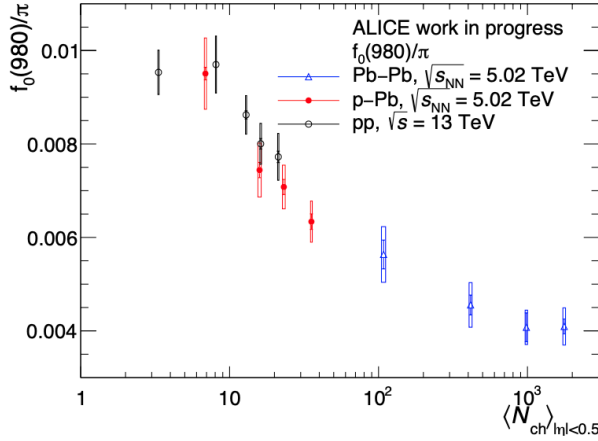


Figure 7.11: The particle yield ratio of $f_0(980)$ to the charged pion as a function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV, p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV.

function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV, p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The ratio decreases with increasing multiplicity in pp and p–Pb collisions and ceases to decrease in Pb–Pb collisions. Such a consistent ratio in Pb–Pb collisions indicates the rescattering effects between them are not much different and can also be explained with the strangeness enhancement of K^{*0} is saturated in Pb–Pb collisions, favoring no strange quarks in $f_0(980)$.

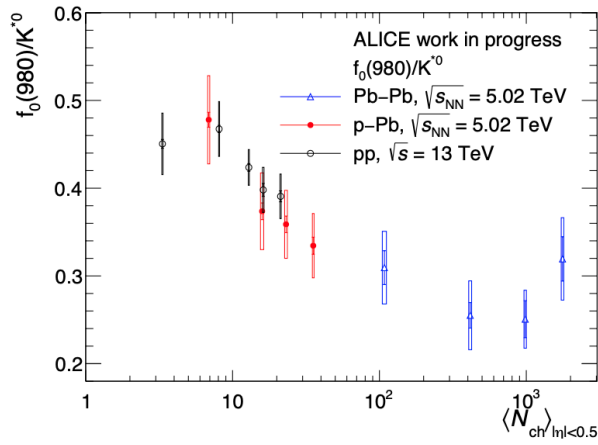


Figure 7.12: The particle yield ratio of $f_0(980)$ to the K^{*0} as a function of charged particle multiplicity in pp collisions at $\sqrt{s} = 13$ TeV, p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, and Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV.

7.3 Event plane dependent production

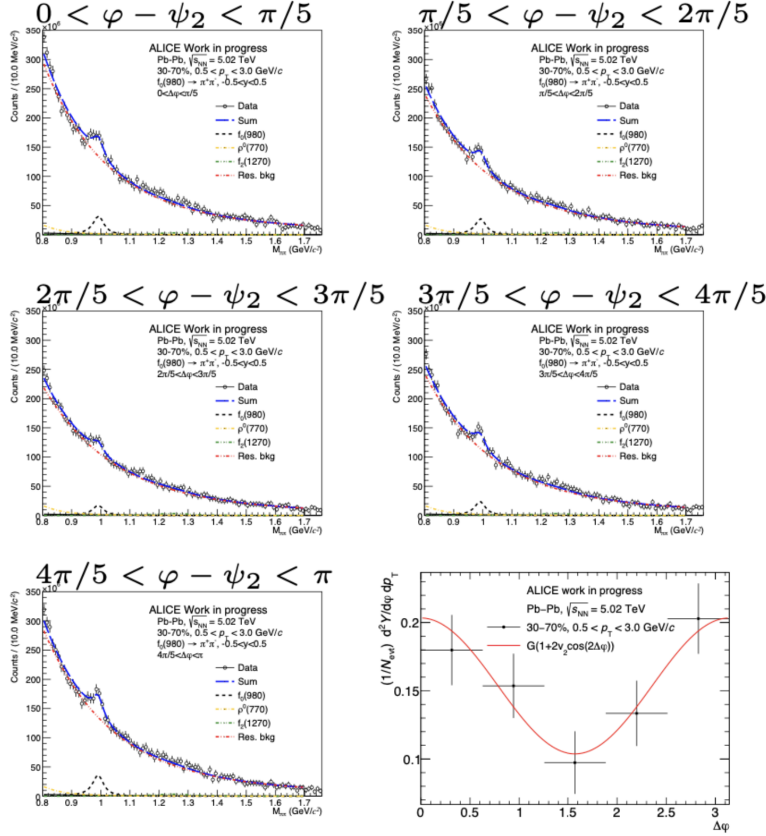


Figure 7.13: Invariant mass distributions of $\pi^+\pi^-$ for different event planes in 30–70% Pb–Pb collisions $\sqrt{s_{\text{NN}}} = 5.02$ TeV for $0.5 < p_T < 3$, which are fitted with the same function described in section ?? and the event-plane dependent invariant yield of $f_0(980)$ fitted with the second harmonic Fourier function (right lowermost).

Figure 7.13 shows the invariant mass distribution of $\pi^+\pi^-$ for different event planes in 30–70% Pb–Pb collisions, where the event plane is reconstructed using the V0A and V0C detectors (see section 4.3.1). The right lowermost panel of figure 7.13 shows the fully corrected invariant yield as a function of relative azimuthal angle to the reconstructed second harmonic event plane for the $0.5 < p_T < 3$ GeV/c range. The yield

distribution is fitted with the Fourier series to extract the anisotropic flow coefficients.

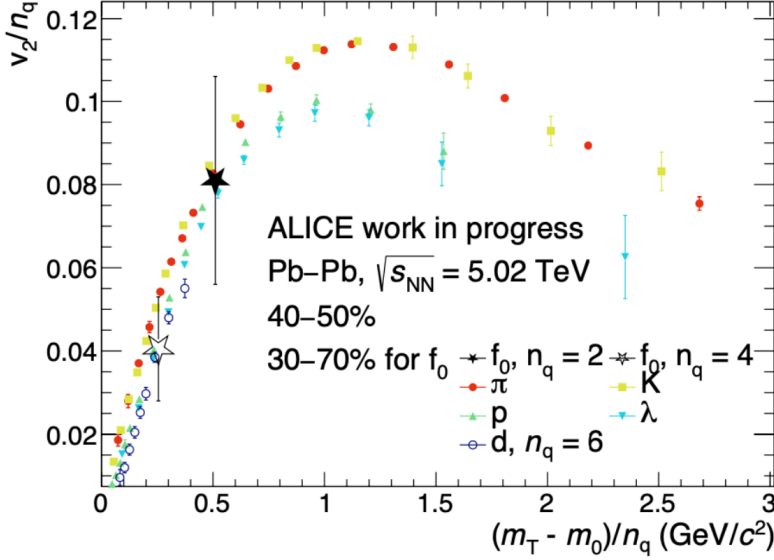


Figure 7.14: NCQ-scaled anisotropic flow coefficients (v_2/n_q) for different particle species in Pb–Pb collisions. The v_2 for $f_0(980)$ is scaled with two assumptions, $n_q = 2$ and 4.

Figure 7.14 shows the anisotropic flow coefficients scaled with the number of constituent quarks (NCQ) as a function of transverse kinetic energy, which is expressed as $m_T - m_0 = \sqrt{m_0^2 + p_T^2} - m_0$, scaled with NCQ for different particles such as mesons, baryons, and deuteron. Their naive agreements in low kinetic energy region indicate the collectivities for different particle species are originating from the anisotropic modulation of constituents. However, there is a splitting between baryons and mesons in a high kinetic energy region owing to the radial flow effects. Such a phenomenon can be applied to the v_2 of $f_0(980)$ to discuss the internal structure of $f_0(980)$ with possible two assumptions of $n_q = 2$ and 4. The v_2 of $f_0(980)$ scaled with two NCQ assumptions is also consistent with other particles at low kinetic energy where the v_2 is just proportional to the kinetic energy, still worth showing this as a verification of the approach. Currently, due to the insufficient statistics collected during Pb–Pb collisions in Run 2, experimental significance at high kinetic

energy is hardly achievable.

The author of the present thesis plans to analyze much more Pb–Pb collisions data collected during Run 3 thanks to the significant upgrade of the TPC to finally approve the internal structure of $f_0(980)$.

Chapter 8

Conclusions

The nature of $f_0(980)$ is still far from the concrete agreement since its first observation, and the measurement of the $f_0(980)$ production and the modification of the $f_0(980)$ yield in the relativistic nucleus–nucleus collisions is expected to largely contribute to the better understanding the nature of the $f_0(980)$ resonance throughout studies described in the present thesis.

The multi-dimensional measurements have been conducted using data collected during Run 2 period. The multiplicity dependence of the $f_0(980)$ production is measured with different multiplicity or centrality estimators. The γ_s -canonical statistical model (CSM) is utilized to predict the particle yield ratio of the $f_0(980)$ with different assumptions of the hidden strangeness for the $f_0(980)$.

The comparison of the particle production can be performed with the measurement of the particle yield ratio, which exhibits the relative enhancement or suppression of the particle production. The production of charged pions is not affected by the strangeness enhancement due to the absence of strange quarks in the charged pion and by the interaction with the hadron gas. The production of $K^{*0}(892)$ resonances is enhanced due to the strangeness enhancement because the K^{*0} includes one strange quark. Conversely, the lifetime of K^{*0} is short enough for its

decay products to interact with the hadron gas, suppressing measured K^{*0} yields.

The p_T -integrated particle yield ratio of $f_0(980)$ to charged pions decreases with increasing charged-particle multiplicity from pp to Pb–Pb collisions, which represents the relative suppression for the $f_0(980)$ production. With the consideration of the short lifetime for the $f_0(980)$, the suppression is mostly due to the rescattering effects. The p_T -differential double ratio of f_0/π in high-multiplicity to low-multiplicity p–Pb collisions exhibits significant suppression only at low $p_T < 4$ GeV/c and no suppression at high $p_T > 4$ GeV/c, indicating similar p_T dependence of rescattering effects, which can be observed from $\rho^0(770)/\pi$, K^{*0}/K , and so on.

The p_T -integrated particle yield ratio of $f_0(980)$ to K^{*0} also monotonously decreases with increasing charged-particle multiplicity from pp to p–Pb collisions and stops to decrease in Pb–Pb collisions. The relative suppression for the $f_0(980)$ production can be understood as the relative enhancement for the K^{*0} production, and the multiplicity dependence of the enhancement for the K^{*0} production is just the same as the multiplicity dependence of the strangeness enhancement. In the situation that both particle yields have similar lifetimes and are suppressed due to the rescattering effects, the suppression for the $f_0(980)$ production can be explained with no strangeness enhancement for the $f_0(980)$. Little effect from the rescattering process is suggested in the p_T -double differential ratio of f_0/K^{*0} in high-multiplicity to low-multiplicity p–Pb collisions, which exhibits the suppression in the measured entire p_T range. The p_T dependence of the double ratio is different from the dependence of rescattering effects, suggesting a different source for the difference.

The canonical statistical model (CSM) also helps to build up the interpretations of experimental results by calculating the particle yield ratio of the $f_0(980)$. While no rescattering effects are implemented in CSM, the particle yield ratio of $f_0(980)$ to charged pions is expected to increase with increasing multiplicity for the assumption of $|S|_{f_0} = 2$, which is the opposite trend of experimental data. CSM with the assumption of $|S|_{f_0} = 0$ provides better description for f_0/π , still overestimating it due to no rescattering effects. CSM qualitatively predicts f_0/K^{*0} to be decreasing with multiplicity for the assumption of $|S|_{f_0} = 0$ while esti-

inating f_0/K^{*0} to be increasing with multiplicity for the assumption of $|S|_{f_0} = 2$. All comparisons with CSM predictions support the absence of strange quarks in the $f_0(980)$.

Furthermore, the nuclear modification factor for $f_0(980)$ is measured for different multiplicity classes in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The modification exhibits stronger suppression for $f_0(980)$ than that for charged particles only at $p_T < 4$ GeV/c and no suppression at high $p_T > 4$ GeV/c. In addition to the p_T dependence of the suppression, the suppression becomes more significant with increasing charged-particle multiplicity, which are the same dependencies as rescattering effects. The nuclear modification factor in p-Pb collisions exhibits another valuable information, the Cronin-like enhancement. Such an enhancement is only observed for baryons, while no enhancement is observed for mesons. Based on no Cronin-like enhancement for $f_0(980)$, the number of constituent quarks is estimated to be smaller than three, favoring the conventional diquark structure.

Finally, the elliptic flow of $f_0(980)$ is scaled with the number of constituent quarks to explore the internal structure of $f_0(980)$. The scaled elliptic flow is in good agreement with those of other particles, proving the validity of the method. However, the current scaling does not provide information on the internal structure of $f_0(980)$ because the statistics are not enough to reach a higher kinetic energy region. This will be further improved with Run 3 data, where the collected data in Pb-Pb collisions are expected to be much larger.

Appendix A

Ridge and flow coefficients in small collision systems

The two-particle correlation function is measured as a function of the relative pseudorapidity ($\Delta\eta$) and the azimuthal angle difference ($\Delta\varphi$) between the trigger and the associated particles, which is expressed as

$$\frac{1}{N_{\text{trig}}} \frac{d^2 N_{\text{pair}}}{d\Delta\eta d\Delta\varphi} = B(0,0) \frac{S(\Delta\eta, \Delta\varphi)}{B(\Delta\eta, \Delta\varphi)} \Big|_{p_{T,\text{trig}}, p_{T,\text{assoc}}}, \quad (\text{A.1})$$

where $p_{T,\text{trig}}$ and $p_{T,\text{assoc}}$ are the transverse momenta of the trigger and associated particles, respectively, N_{trig} is the number of trigger particles, and N_{pair} is the number of trigger-associated particle pairs. The average number of two-particle pairs in the same event and in mixed events are denoted as $S(\Delta\eta, \Delta\varphi)$ and $B(\Delta\eta, \Delta\varphi)$, respectively. The normalization of $B(\Delta\eta, \Delta\varphi)$ is performed with its value at $\Delta\eta$ and $\Delta\varphi = 0$, represented as $B(0,0)$. Acceptance effects are corrected by dividing $S(\Delta\eta, \Delta\varphi)$ with $B(\Delta\eta, \Delta\varphi)/B(0,0)$. The right-hand side of equation A.1 is corrected for the track reconstruction efficiency, which is mainly relevant for the associated particles as a function of p_T and pseudorapidity. Primary vertices of events to be mixed are required to be within the same, 2 cm wide, z_{VTX} interval [139, 140] for each multiplicity class. The final per-trigger

yield is constructed by averaging correlation functions over these primary vertex bins.

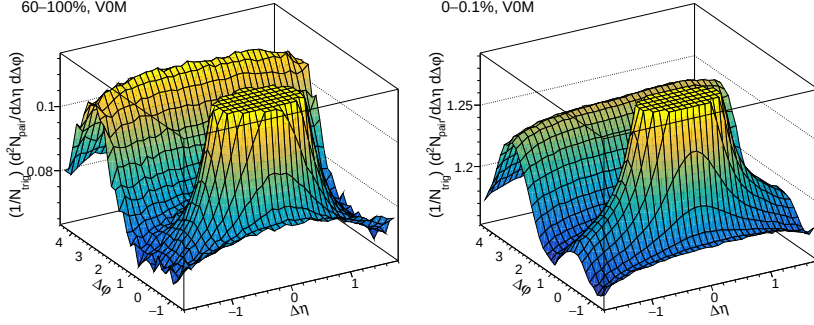


Figure A.1: Two-particle correlation functions as functions of $\Delta\eta$ and $\Delta\varphi$ in low-multiplicity (60–100%, left) and high-multiplicity (0–0.1%, right) events. Note that the near-side jet peaks exceed the chosen range of the z-axis. The intervals of $p_{T,\text{trig}}$ and $p_{T,\text{assoc}}$ are $1 < p_T < 4$ GeV/c in both cases.

Figure A.1 shows the two-particle correlation functions in pp collisions at $\sqrt{s} = 13$ TeV for low-multiplicity (left) and high-multiplicity (right) events for the p_T range of $1 < p_{T,\text{trig(assoc)}} < 4$ GeV/c. In high-multiplicity events, a clear long-range near-side peak, a so-called ridge, is observed, while no ridge structure is in low-multiplicity events.

Ridge yields at large $\Delta\eta$ gap are extracted for different multiplicity classes and p_T intervals. The large $\Delta\eta$ range is required to be $1.6 < |\Delta\eta| < 1.8$, which is the range where the tracking quality – efficiency and precision – is the best. The ridge yield is only measured for $p_T > 1$ GeV/c. Below 1 GeV/c, the jet-like contribution to the correlations function extends into the region where the ridge yield is measured, $1.6 < |\Delta\eta| < 1.8$. In this region, the one dimensional $\Delta\varphi$ correlations, or the so-called per-trigger yield, are expressed as

$$\frac{1}{N_{\text{trig}}} \frac{dN_{\text{pair}}}{d\Delta\varphi} = \int_{1.6 < |\Delta\eta| < 1.8} \left(\frac{1}{N_{\text{trig}}} \frac{d^2N_{\text{pair}}}{d\Delta\eta d\Delta\varphi} \right) \frac{1}{\delta_{\Delta\eta}} d\Delta\eta - C_{\text{ZYAM}}, \quad (\text{A.2})$$

where $\delta_{\Delta\eta} = 0.4$ is the normalization factor to get the per-trigger yield per unit pseudorapidity.

The pedestal of the correlation is subtracted by means of the Zero-Yield-At-Minimum (ZYAM) procedure [141]. The minimum yield (C_{ZYAM}) at $\Delta\varphi = \Delta\varphi_{\text{min}}$ in the $\Delta\varphi$ projection (note that the value of $\Delta\varphi_{\text{min}}$ can be different in data and in models) is obtained from a fit function, which fits the data with a Fourier series up to the third harmonic. The per-trigger yield at $\Delta\varphi_{\text{min}}$ is zero after subtracting C_{ZYAM} from the $\Delta\varphi$ projection.

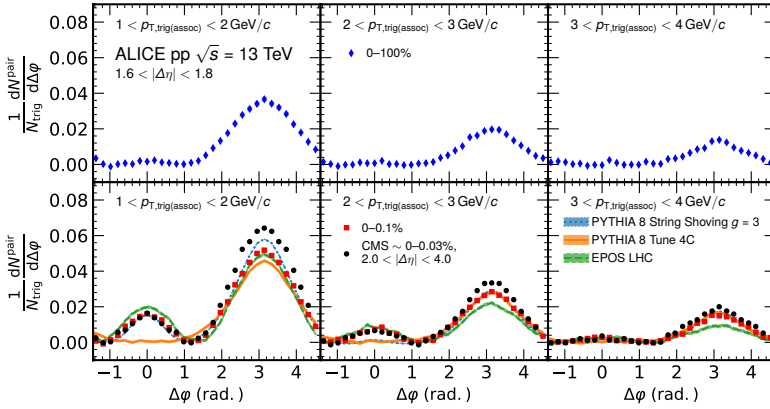


Figure A.2: One-dimensional $\Delta\varphi$ correlations in the large $\Delta\eta$ projection for different p_T intervals in minimum bias (upper panels) and high-multiplicity (lower panels) events after the ZYAM subtraction. Transverse momentum intervals of the trigger particles and associated particles are $1 < p_T < 2$ (left), $2 < p_T < 3$ (middle) and, $3 < p_T < 4$ GeV/c (right), respectively. The model predictions shown in lower panels are calculated using PYTHIA 8 String Shoving, PYTHIA 8 Tune 4C, and EPOS LHC [20].

Figure A.2 shows the projected one dimensional $\Delta\varphi$ correlations at large pseudorapidity of $|\Delta\eta| > 1.6$ using equation A.2 for different p_T intervals after applying the ZYAM process. The near-side ($\Delta\varphi \sim 0$) ridges are clearly observed in high-multiplicity events, while there is no significant signal in minimum-bias events.

The measurements in high-multiplicity events are compared with the results reported by the CMS experiment [142]. CMS experiment obtains charged particle multiplicity by counting the number of charged particles satisfying $p_T > 0.4$ GeV/c in $|\eta| < 2.4$. In the present analysis, the event

multiplicity is determined using the VOM. The difference in the multiplicity selection between ALICE (forward) and CMS (midrapidity) is studied using PYTHIA 8 simulations, and it is found that the calculated multiplicity with the CMS definition is 20% larger than the one from ALICE when compared in the acceptance region of the measurements, $|\eta| < 0.9$. Near-side ridges in all transverse momentum intervals are comparable. The larger away-side yields observed in figure A.2 for the CMS results can be attributed to the overlap in η acceptance between the multiplicity selection and the measurement of the correlations function, possibly inducing so-called autocorrelations [10].

In figure A.2, the ALICE results are also compared with model predictions where a comparable high-multiplicity selection and the same $\Delta\eta$ projection range are applied. The selection of high-multiplicity events in the models is performed by requiring a minimal number of charged particles emitted within the VOM detector acceptance. In the case of PYTHIA8 Tune4C, the 0–0.1% multiplicity class threshold is 105 in terms of charged particles in the V0A and V0C. The threshold for EPOS LHC and PYTHIA8 String Shoving are estimated to be 110 and 108, respectively. The magnitude of string shoving (g) is set to 3.0, where the produced particle multiplicity and correlations are comparable. The statistical uncertainties owing to the limited number of events for the model calculations are shown as bands in each figure. The PYTHIA8 String Shoving reproduces the near-side ridge yield and slightly overestimates the away-side yield for the interval of $1 < p_T < 2$ GeV/c. However, PYTHIA8 String Shoving model underestimates the near-side ridge yield for $p_T > 2$ GeV/c, which is expected from the setup of PYTHIA8 String Shoving. The PYTHIA8 Tune4C model does not exhibit any near-side ridge, which is expected and nice to guarantee no contamination from non-flow effects into the measurement of the ridge. It slightly underestimates the away-side peak for $1 < p_T < 2$ GeV/c and provides good estimates for $p_T > 2$ GeV/c. On the other hand, EPOS LHC describes the trend of the ridge yield quantitatively better in the $2 < p_{T, \text{trig}(\text{assoc})} < 4$ GeV/c range while overestimating the near-side ridge yield for $p_{T, \text{trig}(\text{assoc})} < 2$ GeV/c range.

The ridge yield (Y^{ridge}) is calculated by integrating the near-side peak of the $\Delta\varphi$ projection over $|\Delta\varphi| < |\Delta\varphi_{\text{min}}|$ after the ZYAM pro-

cedure,

$$Y^{\text{ridge}} = \int_{|\Delta\varphi| < |\Delta\varphi_{\min}|} \frac{1}{N_{\text{trig}}} \frac{dN_{\text{pair}}}{d\Delta\varphi} d\Delta\varphi. \quad (\text{A.3})$$

The ridge yield is further studied in events imposing a hard jet or a high- p_T leading particle in the midrapidity region. This event scale is set by requiring a minimum p_T of the leading track ($p_{T,\text{LP}}$) or the reconstructed jet ($p_{T,\text{jet}}^{\text{ch}}$) at the midrapidity region. The leading track is selected within $|\eta| < 0.9$ and the full azimuthal angle. Jets are reconstructed with charged particles only (track-based jets) with the anti- k_T algorithm [143] and the resolution parameter $R = 0.4$. The recombination scheme used in this study is the p_T scheme. Jets are required to be in $|\eta_{\text{jet}}| < 0.4$ and in the full azimuthal angle. The p_T of jets $p_{T,\text{jet}}^{\text{ch}}$ is corrected for the underlying event density that is measured using the k_T algorithm with $R = 0.2$ [144].

The near-side jet-like peak yield is measured from the near-side $\Delta\eta$ correlations to quantify the variation of the near-side jet-like peak with event-scale selections with a minimum $p_{T,\text{LP}}$ or $p_{T,\text{jet}}^{\text{ch}}$. The near-side is defined as $|\Delta\varphi| < 1.28$, where the correlation function is projected onto the $\Delta\eta$ axis. The projection range, ± 1.28 , is chosen to fully cover $\Delta\varphi_{\min}$. The near-side $\Delta\eta$ correlations are then constructed as

$$\frac{1}{N_{\text{trig}}} \frac{dN_{\text{pair}}}{d\Delta\eta} = \int_{|\Delta\varphi| < 1.28} \left(\frac{1}{N_{\text{trig}}} \frac{d^2N_{\text{pair}}}{d\Delta\eta d\Delta\varphi} \right) \frac{1}{\delta_{\Delta\varphi}} d\Delta\varphi - D_{\text{ZYAM}}, \quad (\text{A.4})$$

where $\delta_{\Delta\varphi} = 2.56$ is the normalization factor to obtain per-trigger yield per unit azimuthal angle. The minimum yield, (D_{ZYAM}) of the $\Delta\eta$ correlations is found within $|\Delta\eta| < 1.6$ and used for the subtraction from the $\Delta\eta$ correlations, which results in zero-yield at the minimum. The near-side jet-like peak yield (Y^{near}) is calculated by integrating the $\Delta\eta$ correlations over $|\Delta\eta| < 1.6$,

$$Y^{\text{near}} = \int_{|\Delta\eta| < 1.6} \left(\frac{1}{N_{\text{trig}}} \frac{dN_{\text{pair}}}{d\Delta\eta} \right) d\Delta\eta \quad (\text{A.5})$$

The ridge yields as function of the minimum $p_{T,\text{LP}}$ ($p_{T,\min}^{\text{LP}}$) and $p_{T,\text{jet}}^{\text{ch}}$ ($p_{T,\min}^{\text{jet}}$) selections are shown in figure A.3. High-multiplicity events with

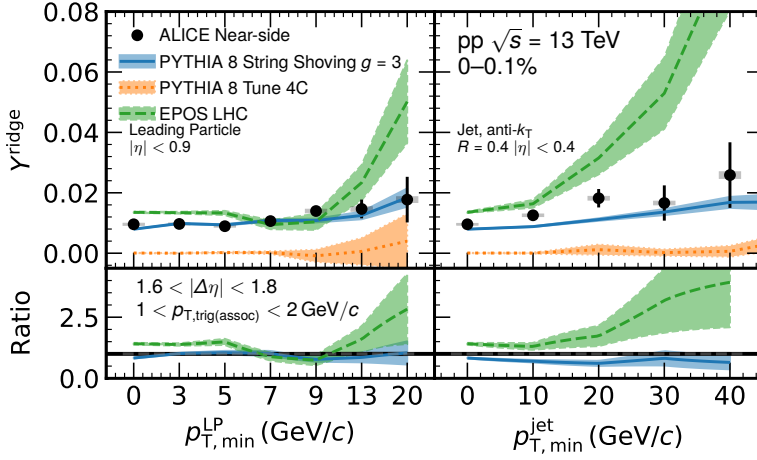


Figure A.3: Near-side ridge yield as a function of the $p_{T,\min}^{\text{LP}}$ (left) and $p_{T,\min}^{\text{Jet}}$ (right). Data points (filled circles) show the ALICE measurement. The statistical and systematic uncertainties are shown as vertical bars and boxes, respectively. As the ridge yield is obtained in the same operational way for data and models, the upper limit of the systematic uncertainty due to jet contamination, which is 18.9%, is not included in the figure. The data are compared with predictions of models, which are represented by colored bands. The bottom panel shows a ratio of the models to the data. The uncertainty of the data is represented by the gray band centered around unity.

imposed event-scale bias exhibit increased ridge yields when compared to unbiased HM events. A small increase of the ridge yields as a function of $p_{T,\text{LP}}$ or $p_{T,\text{jet}}^{\text{ch}}$ is observed, and there is no difference between the two event-scale selections within the uncertainties. Comparisons to model calculations show that PYTHIA8 String Shoving provides a comparable trend with data but underestimates the ridge yield. On the other hand, EPOS LHC overestimates the ridge yield while providing a comparable trend with the data. The origin of the enhanced ridge yields for higher momentum jet-tagged events is not clear to date, but it might be related to the expected smaller impact parameters for dijet or multi-jet production events as studied in [19].

The near-side jet fragmentation yield is measured as a function of

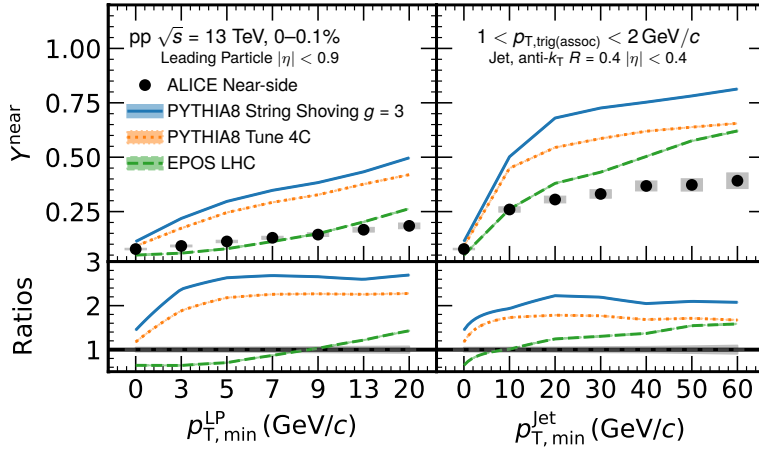


Figure A.4: Near-side jet fragmentation yield as a function of the $p_{T,min}^{LP}$ (left) and $p_{T,min}^{jet}$ (right). The filled circles show measurements with ALICE. The statistical and systematic uncertainties are shown as vertical bars and boxes, respectively. The measurements are compared with model descriptions from PYTHIA8 Tune 4C, PYTHIA8 String Shoving, and EPOS LHC for both selections. The total uncertainty of the ratio is represented by the gray band centered around unity.

minimum $p_{T,LP}$ and $p_{T,jet}^{ch}$ in figure A.4 to further test the models that aim to describe the near-side ridge. EPOS LHC provides comparable estimates of the near-side jet fragmentation yield, while PYTHIA8 Tune 4C and PYTHIA8 String Shoving overestimate the near-side yields for both event selections. In all models, if the ridge is due to final-state interactions, e.g., EPOS LHC and PYTHIA8 String Shoving, one also expects the near-side jet fragmentation yield to be affected. This can be observed when comparing the measured near-side jet fragmentation yields with PYTHIA8 calculations with and without String Shoving. Those results are expected to provide constraints beyond traditional ridge measurements that challenge existing models.

The flow coefficients can be extracted from the long-range $\Delta\varphi$ correlations with the template fit method, which is firstly discussed by ATLAS [145]. In figure A.2, every $\Delta\varphi$ correlations exhibit the peak on the away side, which is mostly attributed to the correlations of leading

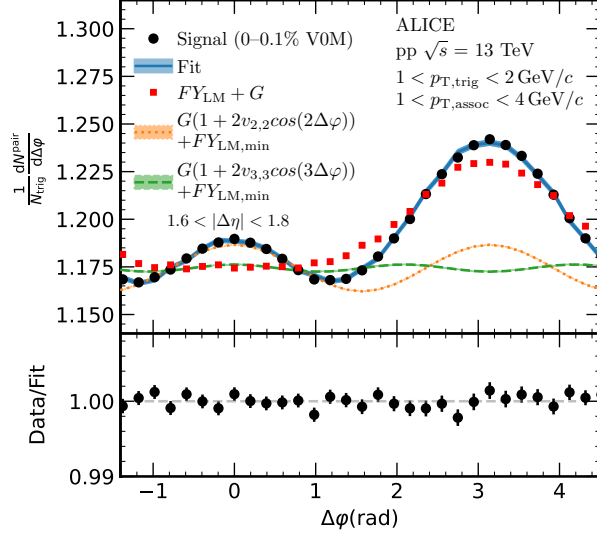


Figure A.5: The template fit results with the low-multiplicity templates. The black markers show the signal for the 0 – 0.1% multiplicity percentile and its fit as a blue band. The red squares correspond to the low-multiplicity signal. The orange and green curves correspond to the extracted v_2 and v_3 signals, respectively. The χ^2 divided by the number of degrees of freedom is 0.894

jets and balancing jets. Therefore, such an away-side peak should be subtracted to exclusively extract the flow coefficients. The template fit method is performed with two assumptions to achieve the proper subtraction of non-flow effects. One assumption is that no flowing particles are there in low-multiplicity events. The other assumption is that the away-side jet fragmentation shape is not changing with the multiplicity class. With those assumptions, the template fit is expressed as

$$Y_{\text{HM}}(\Delta\varphi) = G (1 + 2v_{2,2} \cos(2\Delta\varphi) + 2v_{3,3} \cos(3\Delta\varphi)) + F Y_{\text{LM}}(\Delta\varphi) \quad (\text{A.6})$$

where $Y_{\text{LM}}(\Delta\varphi)$ ($Y_{\text{HM}}(\Delta\varphi)$) is the long-range $\Delta\varphi$ correlations in low-(high-) multiplicity events, G is the normalization factor, $v_{n,n}$ is the flow coefficients of two-particles correlations, and F is the ratio of away-side jet fragmentation in high-multiplicity to low-multiplicity events. Especially, $Y_{\text{LM}}(\Delta\varphi)$ is named the template that is assumed to describe the

away-side jet fragmentations. The validity of the subtraction of the away-side peak is mostly up to the estimation of F . The direct measurement of away-side jet fragmentations is not possible due to the flow modulation along $\Delta\varphi$ direction in high-multiplicity events. Meanwhile, the ratio of the away-side to the near-side jet fragmentations indicates the acceptance of particle pairs in jets, where all pairs of two-particle correlations in the same jet can be captured near $(\Delta\eta, \Delta\varphi) = (0, 0)$, and some of the pairs of two-particle correlations in the leading jet and the balancing jet can not be captured due to the limited η acceptance [146].

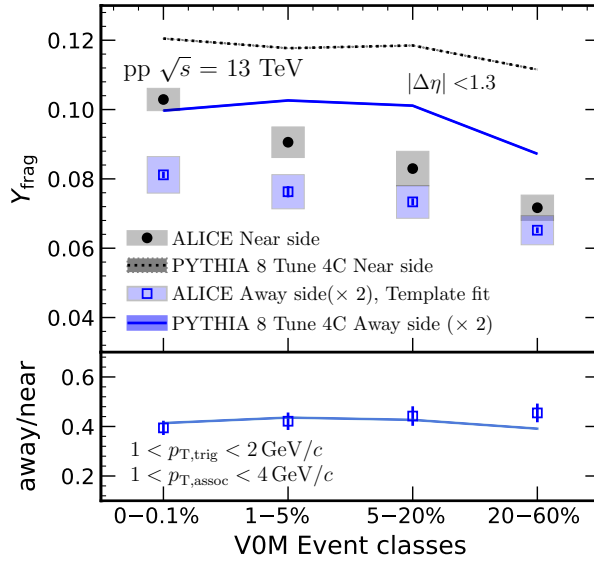


Figure A.6: The Y_{frag} for the near- and away-side as a function of multiplicity percentiles with both ALICE and PYTHIA data. The ratio includes the combined statistical and systematic errors in quadrature.

To evaluate the acceptance of two-particle pairs, away-side jet fragmentation in high-multiplicity events can be indirectly measured using the F , which is expressed as

$$Y_{\text{frag}}^{\text{away,HM}} = Y_{\text{frag}}^{\text{away,LM}} \times F \quad , \quad (\text{A.7})$$

where $Y_{\text{frag}}^{\text{away,LM}}$ can be directly measured with the assumption of no flowing particle in low-multiplicity events. With the direct measurement of

the near-side jet fragmentations using equation A.5, the acceptance is evaluated as a function of multiplicity class as shown in figure A.6. The jet fragmentation in the away- and near-side from experimental data is compared with that obtained from the PYTHIA 8 event generator with Tune 4C, showing that the jet fragmentation is not well reproduced by the PYTHIA 8. However, the ratio, indicating the acceptance and irrelevant with detailed descriptions of the jet fragmentations, is well reproduced. Because there is no flow modulation in PYTHIA 8 tune 4C, the agreement of acceptance between data and PYTHIA 8 confirms the indirect measurement of the away-side jet fragmentation. Finally, the scale factor, F is proven to be valid for the subtraction.

After the subtraction of away-side jet fragmentations, v_n can be measured using $v_{n,n}$ with the factorization to single harmonics based on equation 1.13, which is expressed as

$$v_n(p_{T,\text{trig}}) = v_{n,n}(p_{T,\text{trig}}, p_{T,\text{assoc}}) / \sqrt{v_{n,n}(p_{T,\text{assoc}}, p_{T,\text{assoc}})}. \quad (\text{A.8})$$

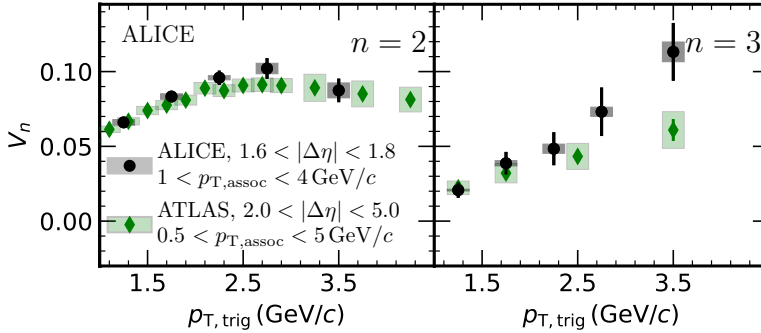


Figure A.7: v_2 and v_3 as a function of $p_{T,\text{trig}}$ in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV. Measured v_2 and v_3 are comparable with those measured by ATLAS

Figure A.7 shows the v_2 and v_3 as a function of $p_{T,\text{trig}}$ in pp collisions at $\sqrt{s} = 13$ TeV. The measured results are compared with the results from ATLAS [145], showing the agreement between them from the same flow extraction method.

Figure A.8 shows the v_2 and v_3 as a function of event-scale selection in high-multiplicity pp collisions at $\sqrt{s} = 13$ TeV. The flow coefficients are

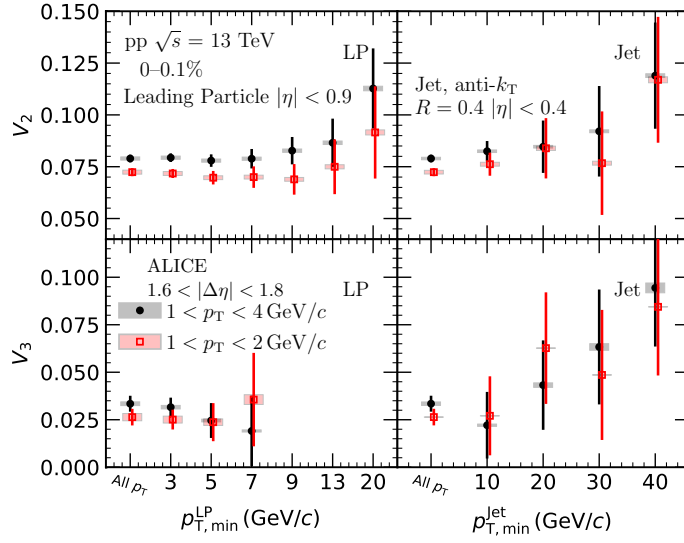


Figure A.8: v_2 and v_3 as a function of event-scale selection using leading hadrons and reconstructed jets

consistent with increasing event scale with poor statistical uncertainties. The event classification with the hard probes can be further improved with the di-jet mass or four jets events, as mentioned in figure A.3.

Appendix B

Machine learning for the angular reconstruction of γ

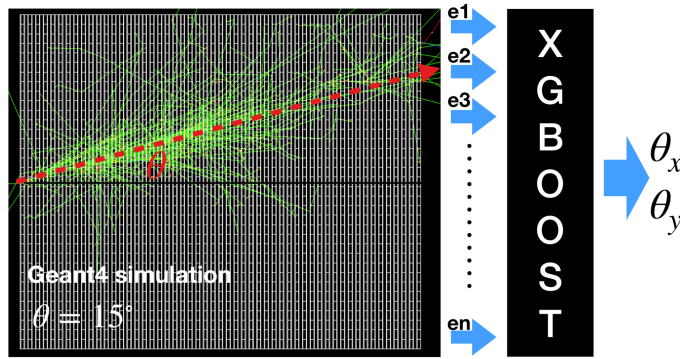


Figure B.1: Schematic of machine learning process

The incident angle of the photons is reconstructed using the XGBoost (XGB) [147], which maps a feature dataset onto a target variable as shown in figure B.1. Here, the feature dataset and the target variable correspond to the energy deposit in each scintillator strip and the incident angle of a photon, respectively, while dimensions of the simu-

lated sampling calorimeter can be found in Ref. [104]. Because the input dataset for the training is representative of the detector response of the sampling calorimeter, the electromagnetic showers are carefully generated in the wide range of $0 < \theta < 50^\circ$ and $0 < \varphi < 360^\circ$. The model, which is trained with the above dataset, is limited to reconstructing the polar angle up to 50° due to the limited input dataset.

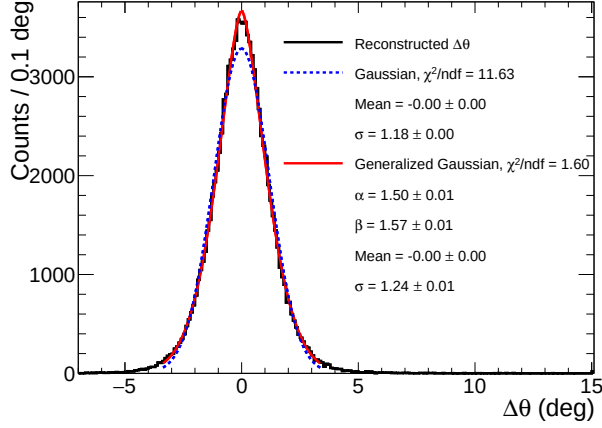


Figure B.2: The relative incident angle for 1-GeV photons generated at $\theta = 10^\circ$. The distribution is fitted with the Gaussian and generalized Gaussian functions. The generalized Gaussian function describes the obtained distribution better.

Figure B.2 shows the distribution of the relative incident angle ($\Delta\theta$) for 1-GeV photons with the true incident angle of $\theta = 10^\circ$. The $\Delta\theta$ is calculated with the radial displacement of the reconstructed incident angle relative to the true incident direction. The central 98% of the distribution is fitted with the Gaussian and generalized Gaussian (GG) functions [148]. The GG function shows better agreement with the obtained distribution, which is expressed

$$f(x; \mu, \alpha, \beta) = \frac{\beta}{2\alpha\Gamma(1/\beta)} e^{-(|x-\mu|/\alpha)^\beta}, \quad (\text{B.1})$$

where μ is the mean value. The parameters α and β determine the scale and shape of the distribution, respectively. The variance in the GG function is given by $\sigma^2 \equiv \alpha^2\Gamma(3/\beta)/\Gamma(1/\beta)$, and the angular resolution is defined as σ .

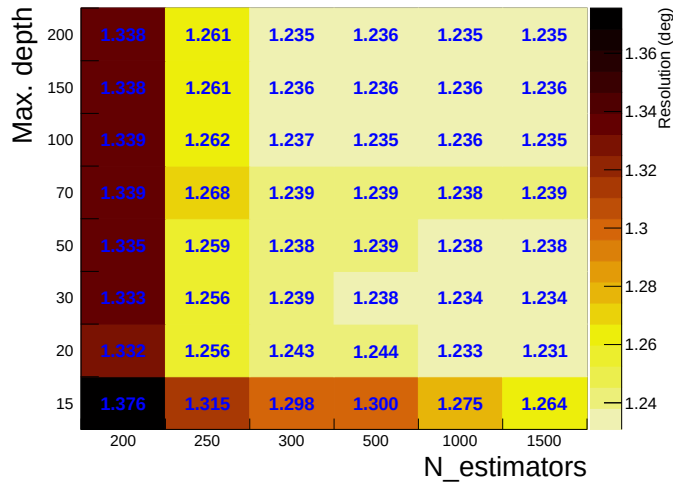


Figure B.3: Angular resolutions are displayed in terms of combination of $N_estimators$ and $Max. depth$. The uncertainty for each value is 0.01 deg.

The hyperparameter in machine learning arranges how the model is trained with the input dataset. As an example, figure B.3 shows the test results for $N_estimators$ and $Max. depth$. The angular resolution improves when both $Max. depth$ and $N_estimators$ increase. Considering the CPU time for the training, the $N_estimators$ and $Max. depth$ are set to 300 and 100, respectively. Similar processes are also performed for different hyperparameters, $subsample$, $learning rates$, and $gamma$. Table B.1 lists the optimized set of hyperparameters.

Table B.1: Hyperparameters of the XGBoost model

Parameter	Function	Default value	Used value
$N_estimators$	The number of decision trees	N.A.	300
$Max. depth$	Possible maximum depth of tree structure	6	100
$Subsample$	Fraction of total data used for a single decision	1	0.8
$Learning rate$	Step length for calculation	0.3	0.02
$Gamma$	Requirement on minimum loss function	0	0

The angular resolution is evaluated by varying the strip width from 5 to 35 mm, as shown in figure B.4, where each point is obtained through the same fit procedure to inspect the dependency of the angular resolution on the strip width. The narrower the strip width, the larger the fea-

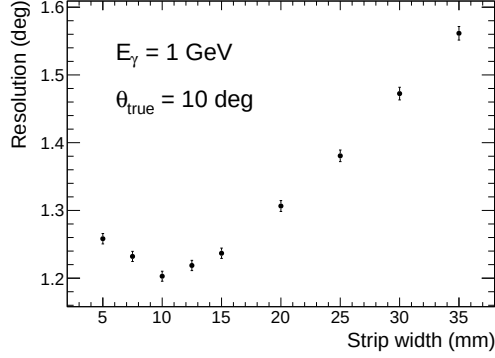


Figure B.4: Angular resolutions as a function of the strip width for 1-GeV photons at $\theta = 10^\circ$

ture dataset size. Consequently, the angle-reconstruction performance is better with 10-mm-wide strips than with 5-mm-wide strips. However, the wider strips hardly provide detailed information on the EM shower. The 15-mm-wide strips yield an angular resolution of $1.24(1) \pm 0.01^\circ$ for photons impacting the center of the strips and are selected considering the actual construction of a prototype detector, although 10-mm-wide strips provide the best resolution.

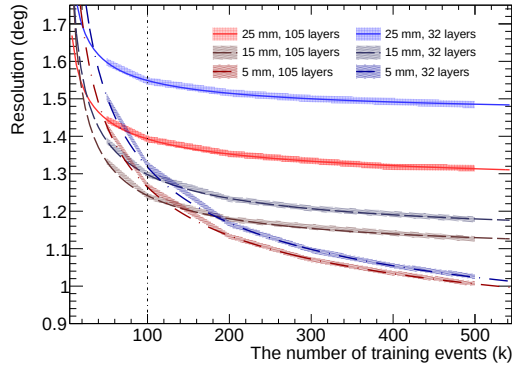


Figure B.5: The angular resolution for 1-GeV photons at $\theta = 10^\circ$ as a function of the number of training samples with different strip widths obtained using the first 32 layers or the full layers. The spectra are fitted with $\sqrt{I^2 + (D/N_{\text{sample}})^k}$. The vertical dashed-dotted line indicates the 100k events.

Figure B.5 shows the angular resolution for different numbers of training samples and detector configurations. The width of each curve represents the statistical uncertainty. All configurations show improvement in the angular resolution with the increasing number of training samples. With an increase in the number of training events, the angular resolution improves most rapidly for the configuration with 5-mm-wide scintillator strips. In other words, the required number of training events should be large enough to fully train the model from the correlations between many features.

The spectra in figure B.5 are fitted with the empirical function, which can be expressed as $\sqrt{I^2 + (D/N_{\text{sample}})^k}$, where I denotes the intrinsic resolution and D and k describe the decreasing trend. The numbers of training samples reaching the angular resolution of $1.1I$ with 25-, 15-, and 5-mm-wide strips are estimated to be 60k, 300k, and 8M, respectively, indicating that the number of training samples is insufficient, especially for 5-mm-wide strips.

Appendix C

SHINCHON project

SHINCHON, which stands for Simulation for Heavy IoN Collision with Heavy-quark and ONia, is a newly developing toolkit for Monte-Carlo simulation for heavy quarks and heavy quarkonia in nucleus-nucleus collisions, enabling complex calculations in dynamical environments. By analogy with JETSCAPE [149], SHINCHON successfully implements the hydrodynamic evolution of medium created from nucleus-nucleus collisions and the dissociation of heavy quarkonia, which is calculated based on the hard thermal loop perturbation theory [150]. A recent research for a detailed study of the in-medium effects for upilon states in small collision systems has been conducted by implementing the hydrodynamic evolution using SONIC and the dissociation process using thermal widths for upilons.

C.1 Framework

The simulation process reported in Ref. [18] can be divided into two parts. One is the hydrodynamic evolution of the initial collision geometry, which is generated by the Glauber model, and the other is the interactions between the evolving medium and upilons. Simulated initial collision geometry is converted into energy density with a scale factor,

which yields the same charged particle multiplicity with experiments in p–Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV [151]. The obtained scale factor in p–Pb collisions is taken for the conversion in pO and OO collisions owing to no experimental data yet. The scale factor is dependent on the centrality class in Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV while no dependency is observed in p–Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV. After the conversion, the pre-equilibrium dynamics [152] is applied by simply approximating the temperature would be effectively 20% higher.

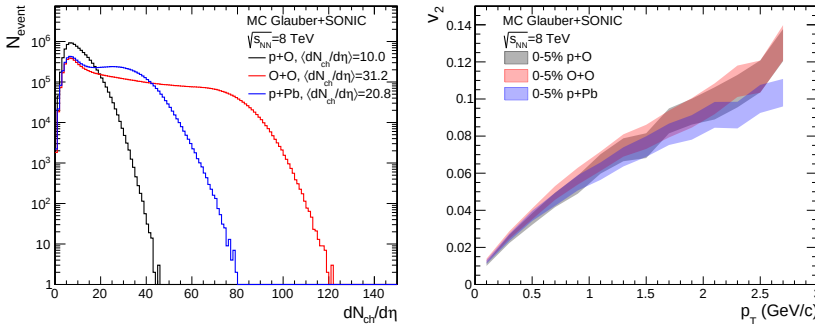


Figure C.1: The charged particle multiplicity and the second harmonic flow coefficients calculated in SONIC for different collision systems

Figure C.1 shows the charged particle multiplicity (left) and the second harmonic flow coefficients (right) in p–Pb, pO, and OO collisions at $\sqrt{s_{NN}} = 8.16$ TeV. They are calculated at the end of the hydrodynamic evolution, which is performed by SONIC [153]. As the multiplicity is tuned with one in p–Pb collisions, the calculated multiplicity is the same as the experimental data in p–Pb collisions, and the same flow coefficients are also expected with this configuration.

Upsilons are generated based on the initial collision (temperature) distribution. The initial position of the upsilons is determined to be proportional to the initial temperature distribution. The transverse momentum (p_T) is generated based on a Tsallis function [16] that fits experimental data. The azimuth is randomly generated. The medium response is simulated using the snapshots of temperature distributions, which are saved during the hydrodynamic evolution. The medium response is ex-

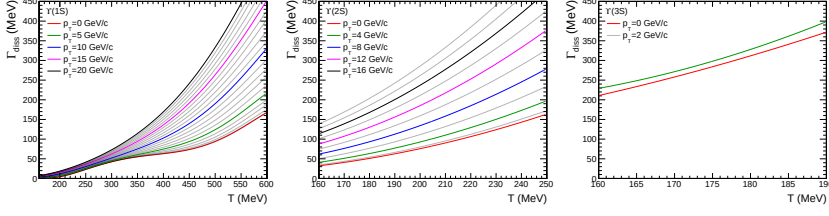


Figure C.2: The thermal widths of $\Upsilon(nS)$ as a function of the temperature for different p_T

pressed as

$$\frac{N(t + \Delta t, p_T)}{N(t, p_T)} = e^{-\int_t^{t+\Delta t} dt' \Gamma_{\text{diss}}(t', p_T)}, \quad (\text{C.1})$$

where Γ_{diss} and Δt is the thermal width and the time step of the snapshot, which is set to 0.02 fm/c, respectively. The dissociation temperature, no binding of upsilons allowed, is set to 600 MeV, 240 MeV, and 190 MeV for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively [154]. Figure C.2 shows the thermal width as a function of temperature for different p_T of upsilons. The thermal widths gradually increase with increasing p_T and the temperature. The thermal width of $\Upsilon(3S)$ is calculated up to very limited p_T range in theory and is extrapolated with the thermal width of $\Upsilon(2S)$ in the present simulation. The extrapolation can be expressed as

$$\Gamma_{\text{diss}}^{\Upsilon(3S)}(p_T) = \Gamma_{\text{diss}}^{\Upsilon(3S)}(p_T = 2 \text{ GeV}/c) \frac{\Gamma_{\text{diss}}^{\Upsilon(2S)}(p_T)}{\Gamma_{\text{diss}}^{\Upsilon(2S)}(p_T = 2 \text{ GeV}/c)}. \quad (\text{C.2})$$

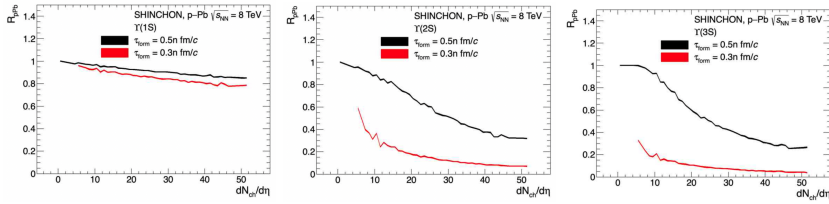


Figure C.3: The nuclear modification factor of $\Upsilon(nS)$ for different formation times

The upsilons are simulated to be formed after the timescale of the formation time, and the formation times are set to 0.5 fm/c, 1.0 fm/c,

and 1.5 fm/c for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively [155]. The medium responses are simulated after the formation of upsilons, hence, some upsilons are produced afterward the hydrodynamic evolution. Figure C.3 shows the nuclear modification factor for different formation times, indicating that the formation time is a key parameter that largely affects simulation results because the formation time is comparable with the time span of the hydrodynamic evolution, especially for excited states.

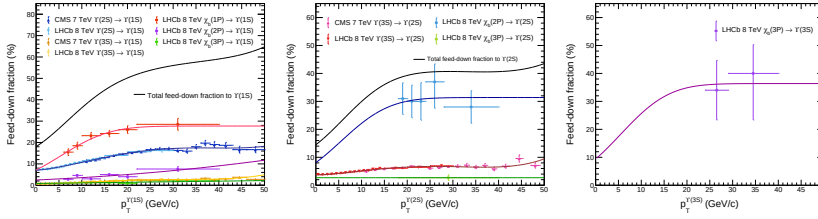


Figure C.4: Feed-down fractions of $\Upsilon(nS)$ obtained in CMS and LHCb experiments

The feed-down correction, which considers products from the decay of higher excited states, is applied to calculate the inclusive bottomonium production. Weak decay modes are not considered as not significant. The feed-down correction for a higher state (Q_m) to a lower state (Q_n) can be expressed as

$$\mathcal{F}_{Q_n}^{Q_m} = \mathcal{B}(Q_m \rightarrow Q_n) \frac{\sigma_{Q_m}}{\sigma_{Q_n}}, \quad (\text{C.3})$$

where $\mathcal{B}$ and σ_Q denote the branching ratio of Q_m into Q_n and the cross section of the corresponding state, respectively. Figure C.4 shows the feed-down fractions which are obtained from experimental data measured by CMS [156] and LHCb [157, 158] in pp collisions. The fractions are fitted with empirical functions for extrapolation. For $\Upsilon(2S)$ and $\Upsilon(3S)$, low p_T region of the fit function is assumed to be the same shape with the fraction of P-wave state to $\Upsilon(1S)$.

The feed-down correction for the nuclear modification factor is performed by calculating the weighted average of the nuclear modification

factor (R), which can be expressed as

$$R_n(p_T) = \sum_i R_i(p_T) \mathcal{F}_{Q_n}^{Q_i}(p_T), \quad (\text{C.4})$$

where R_n is the weighted and averaged value for a certain $\Upsilon(nS)$ state, R_i is the value for a certain state contributing the $\Upsilon(nS)$ state, and $\mathcal{F}_{Q_n}^{Q_i}$ is the feed-down fraction. The nuclear modification factor for χ_b states is assumed to be $R_{\Upsilon(2S)} \approx R_{\chi_b(1P)}$ and $R_{\Upsilon(3S)} \approx R_{\chi_b(2P)} \approx R_{\chi_b(3P)}$ [150].

C.2 Results

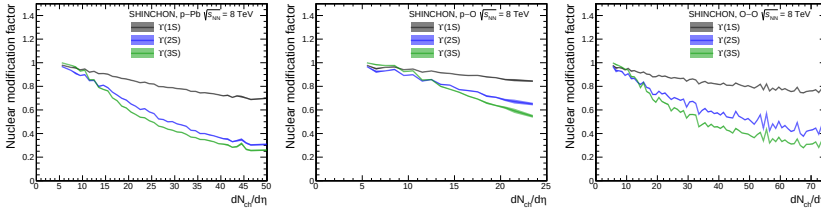


Figure C.5: Nuclear modification factors for $\Upsilon(nS)$ as a function of $dN_{ch}/d\eta$ in p-Pb, pO, and OO collisions at $\sqrt{s_{NN}} = 8$ TeV. The uncertainties represent the statistical uncertainty to the corresponding simulated events.

Figure C.5 shows the nuclear modification factor as a function of multiplicity for $\Upsilon(nS)$ in p-Pb, pO, and OO collisions at $\sqrt{s_{NN}} = 8$ TeV. As the system size increases, the nuclear modification factor decreases, indicating stronger suppression in all three systems. Stronger suppression is visible for higher excited states.

Figure C.6 compares the nuclear modification factor between different collision systems for each upsilon state. The suppression is comparable between p-Pb and pO collisions. The nuclear modification factor in OO collisions is higher than that in p-Pb collisions because the energy density is smaller in OO collisions.

Figure C.7 shows the nuclear modification factor as a function of charged particle density, which can be expressed as

$$S_T = \pi \sqrt{\langle x^2 \rangle \langle y^2 \rangle - \langle xy \rangle^2}, \quad (\text{C.5})$$

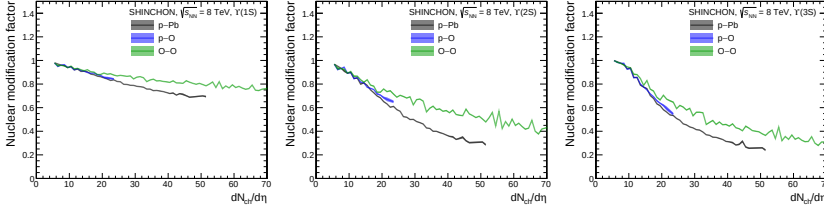


Figure C.6: Nuclear modification factors for $\Upsilon(nS)$ as a function of $dN_{ch}/d\eta$ for a different representation.

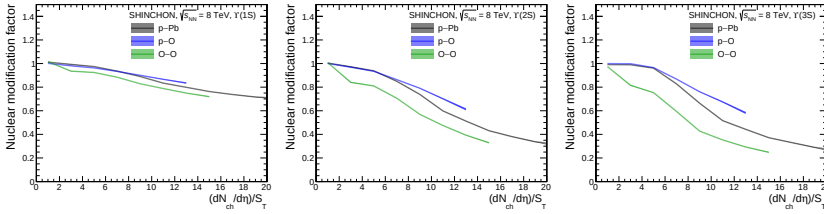


Figure C.7: Nuclear modification factor as a function of $(dN_{ch}/d\eta)/S_T$ for $\Upsilon(nS)$ in p-Pb, pO, and OO collisions at $\sqrt{s_{NN}} = 8$ TeV.

where x and y are the spatial coordinates of the participating nucleons from the MC-Glauber. In the comparison between pO and p-Pb collisions, it shows an agreement at low charged particle density, but a slightly stronger suppression is seen in higher density. In comparison with OO collisions, the nuclear modification factor in OO collisions is smaller than one in pA systems at the same charge particle density because the overall system size and the duration of hydrodynamic evolution are larger in O-O collisions.

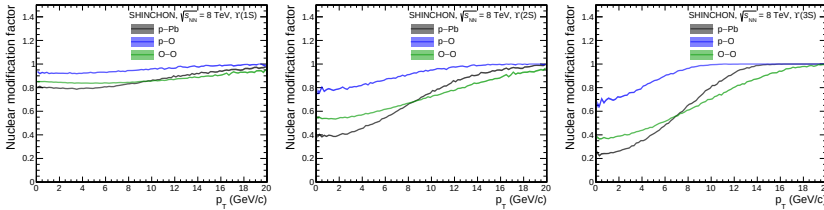


Figure C.8: Nuclear modification factors for $\Upsilon(nS)$ as a function of p_T in p-Pb, pO, and OO collisions at $\sqrt{s_{NN}} = 8$ TeV.

Figure C.8 shows the nuclear modification factor as a function of p_T for $\Upsilon(nS)$ in p-Pb, pO, and OO collisions at $\sqrt{s_{NN}} = 8$ TeV. The nuclear modification factor increases with increasing p_T because the formation is much delayed at high p_T due to the time dilation. There is a crossing point between p-Pb and OO collisions around $p_T = 8$ GeV/c. In low p_T , Υ is more likely to fully interact with medium, not escaping. On the other hand, the duration time of the medium evolution is longer in OO collisions, and upsilons are more likely to interact with the medium in OO collisions in high p_T . These two competing effects result in the crossing point.

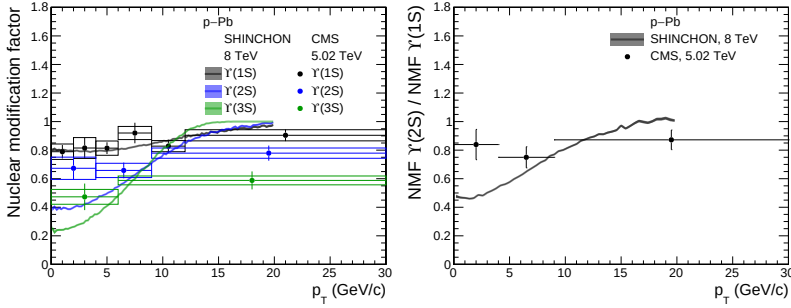


Figure C.9: Nuclear modification factor for $\Upsilon(nS)$ as a function of p_T in p-Pb collisions at $\sqrt{s_{NN}} = 8$ TeV compared with the CMS results at $\sqrt{s_{NN}} = 5.02$ TeV, and the ratio of nuclear modification factors between $\Upsilon(1S)$ and $\Upsilon(2S)$

The left panel of figure C.9 shows the comparison of the nuclear modification factors between the simulated calculations and experimental data measured by CMS [159]. SHINCHON nicely describes the nuclear modification factor for $\Upsilon(1S)$ while not reproducing it for higher excited states. Note that the model calculates the nuclear modification factor in the energy of 8 TeV, in which the average multiplicity is about 15% higher than that at 5.02 TeV [151]. This may affect the comparison, especially in low p_T . The right panel of figure C.9 shows the ratio of the nuclear modification factor for $\Upsilon(2S)$ to $\Upsilon(1S)$ and the comparison of it between the calculation and experimental data to understand the disagreement. Taking a ratio helps cancel out the initial-state effects as one can expect the initial-state effects would be similar between $\Upsilon(2S)$ and

$\Upsilon(1S)$. With the cancellation of initial-state effects, the disagreement is still visible, indicating a different source, rather than initial-state effects, may attribute to the disagreement.

Appendix D

Length-scale calibration in luminosity measurement

The measurement of luminosity provides the baseline for the measurement of cross sections for particle productions. To measure the luminosity in the environment of Run 3 at the LHC, a dedicated framework has been developed, which is implemented to the O2Physics package. The present appendix describes the usage of the `lumi.cxx` in the O2Physics and the reproduction of luminosity measurements, which had been conducted in Run 2, using converted Run 2 data and the application of O2Physics to them.

D.1 Converted Run 2 data

Converted Run 2 data are used to develop a new framework and to reproduce the previous results, where the data are acquired in 2018 for the Van der Meer (vdM) scan. The run numbers of corresponding data are 288866 and 297325 for pp and Pb–Pb collisions, respectively. The converted data, used in the present appendix, are located at (in ALICE legotrain system):

- pp: `PWGZZ/Run3_Conversion/218_20210928–2022`

- Pb–Pb: PWGZZ/Run3_Conversion/217_20210928–2012

It is worth mentioning that Pb–Pb data have no reconstructed tracks from TPC. Considering such an absence, only ITS tracks are used to reconstruct the primary vertex (PV).

D.2 Usage of `lumi.cxx`

The main framework to reconstruct the PV and plot the PV in terms of the timestamp can be executed with the following command:

- `o2-analysis-timestamp -b -aod-file AO2D.root`
`-aod-writer-keep="AOD/EventInfo/0"`
`| o2-analysis-trackextension -b`
`| o2-analysis-trackselection -b`
`| o2-analysis-mm-lumi -b -configuration json://./config.json`
 (for Converted Run 2 data)
- `o2-analysis-timestamp -b -aod-file AO2D.root`
`-aod-writer-keep="AOD/EventInfo/0"`
`| o2-analysis-collision-converter -b`
`| o2-analysis-track-propagation -b -processCovariance`
`| o2-analysis-trackselection -b`
`| o2-analysis-mm-lumi -b -configuration json://./config.json`
 (for Run 3 data)

The “timestamp” task provides the timestamp of a given event. The “track extraction” and “track selection” (or “track propagation”) tasks make it possible to select the proper track and to extend reconstructed tracks with the primary vertex. After generating relevant tables, the “lumi” task (`lumi`) reconstructs the PV. To improve the quality of the refitted PV, specific selections are additionally applied to each track and the PV. Tracks are required to have p_T larger than 0.8 GeV/c and place within $|\eta| < 0.8$. The reduced chi-square of the refitted PV is required to be smaller than 4. The option “-aod-writer-keep” creates the derived table from `lumi`, which is located in

`O2Physics/PWGMM/Tasks/lumi.cxx`.

D.2.1 Inputs

There are a few inputs for the `lumi`, which are described in the table D.1. The `lumi` requires the timestamp for the first event to calculate relative time from the first event. The `lumi` can exclusively select the PV by requiring minimal and maximal numbers of PV contributors. These values can be controlled within a “json” input.

Table D.1: Input parameters for `lumi`

Parameter	function	default
<code>fts</code>	Timestamp of the first event	1530319778000 (pp) 1543488980000 (Pb-Pb)
<code>nContribMin</code>	The minimal number of PV contributors	10
<code>nContribMax</code>	The maximal number of PV contributors	2500

D.2.2 Outputs

As a result of the code execution, the outputs are obtained in the “lumi” directory. There are a few histograms for Length-Scale calibration (LSC). Histograms named “`vertexx(y)_Refitted_timestamp`” include $x(y)_{\text{vtx}}$ in terms of the timestamp. One can select the specific range of the timestamp and project along $x(y)_{\text{vtx}}$ direction. After the projection with the proper timestamp selection, Comparison can be performed for the displacement of the PV with the LHC step length for beam position arrangement. Derived table, “EventInfo” is stored in “AnalysisResults_trees.root”. The tree includes vertex information for each event. This is also input for the non-factorization analysis.

D.3 Length-Scale calibration

D.3.1 pp

Figure D.1 shows the x (left) and y (right) components of the refitted primary vertex in terms of the timestamp in pp collisions. Red lines in each figure represent the range of the timestamp to be projected. One can see that the position of the beam is moving in x -direction while fixing y -direction, and moving in y -direction while fixing x -direction.

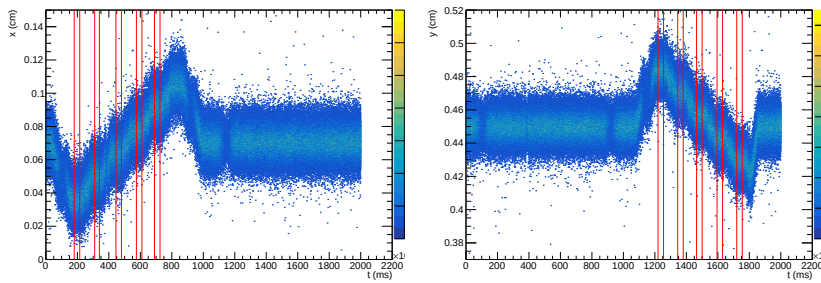


Figure D.1: The x (left) and y (right) components of the refitted primary vertex in terms of the timestamp in pp collisions.

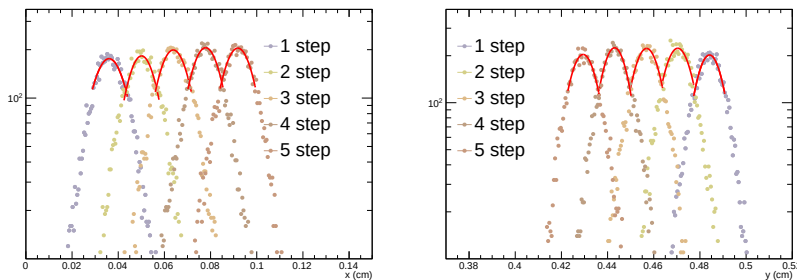


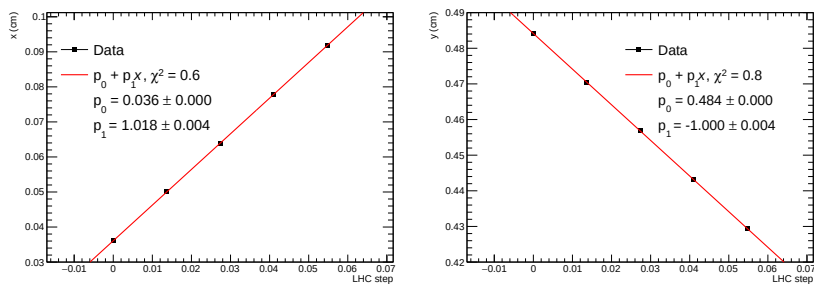
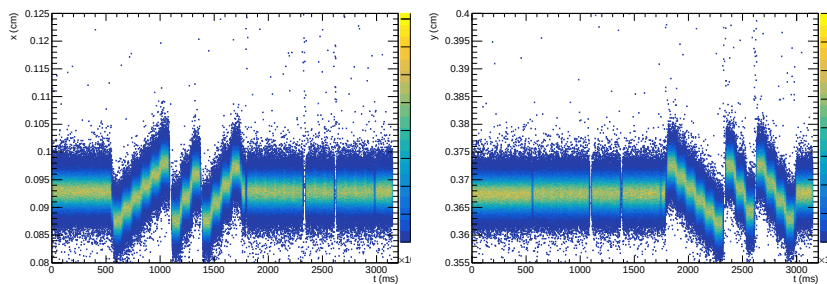
Figure D.2: x (left) and y (right) components of the refitted primary vertex over the specific timestamp range.

Figure D.2 shows the integrated x (left) and y (right) components of the refitted primary vertex over the specific timestamp range. The ranges are summarized in table D.2. As shown in figure D.2, the vertex is moving gradually. Each projected distribution is fitted with the Gaussian function to determine the center of the distribution, which are represented as red lines.

Figure D.3 shows the comparison between the displacement operated by LHC and the refitted PV for x (left) and y (right) directions. The data were fitted with a linear function. The slope of the linear function is estimated to be near unity, guaranteeing good agreement.

Table D.2: The ranges of selected timestamps

Step	Range for x (ms)	Range for y (ms)
1	180–215	1220–1255
2	310–350	1345–1380
3	445–480	1465–1500
4	575–610	1595–1630
5	690–725	1720–1755

Figure D.3: Comparison between the displacement operated by LHC and the refitted PV for x (left) and y (right) directionsFigure D.4: The x (left) and y (right) components of the refitted primary vertex in terms of the timestamp in Pb–Pb collisions.

D.3.2 Pb–Pb

Figure D.4 shows the x (left) and y (right) components of the refitted primary vertex in terms of the timestamp in Pb–Pb collisions. The beam position is moving three times in each direction.

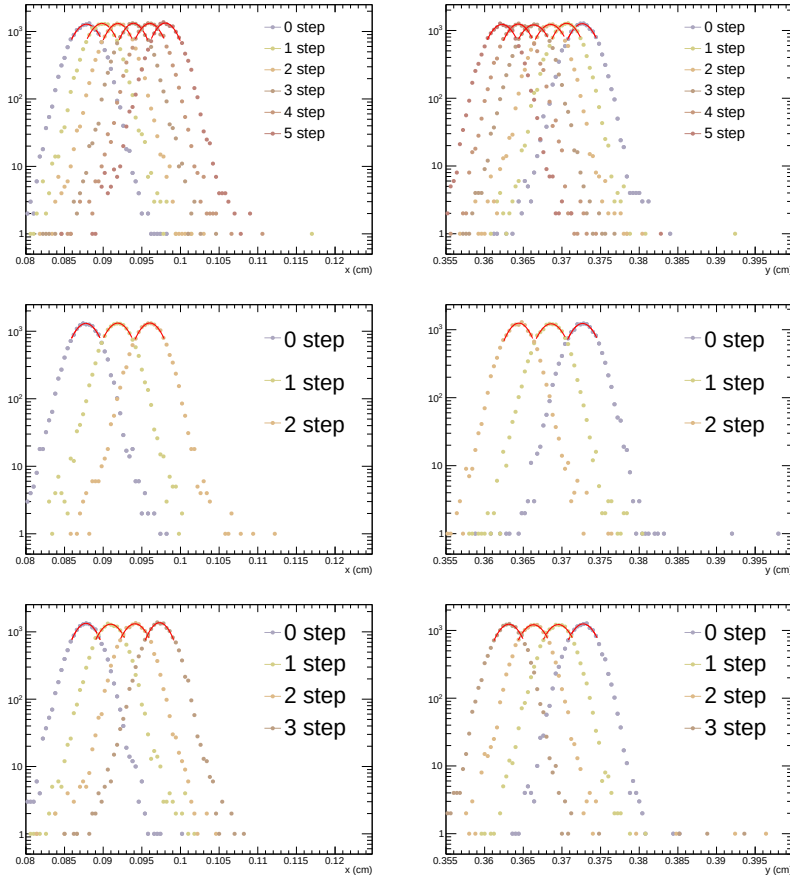


Figure D.5: x (left) and y (right) components of the refitted primary vertex over the specific timestamp range.

Figure D.5 shows the integrated x (left) and y (right) components of the refitted primary vertex over the specific timestamp range. The ranges are summarized in table D.3. As shown in figure D.5, the vertex is moving gradually. Each projected distribution is fitted with the Gaussian function to determine the center of the distribution, which are represented as red

lines.

Table D.3: The ranges of selected timestamps in the unit of millisecond

Stage	1st stage		2nd stage		3rd stage	
Step	Range for x	Range for y	Range for x	Range for y	Range for x	Range for y
1	600–620	1840–1860	1140–1160	2380–2400	1420–1440	2660–2680
2	700–720	1940–1960	1240–1260	2480–2500	1520–1540	2760–2780
3	780–800	2020–2040	1320–1340	2560–2580	1620–1640	2860–2880
4	860–880	2120–2140	-	-	1700–1720	2940–2960
5	940–960	-	-	-	-	-
6	1030–1050	-	-	-	-	-

Figure D.6 shows the comparison between the displacement operated by LHC and the refitted PV for x (left) and y (right) directions. The data are fitted with a linear function. The slope of the linear function is estimated to be 0.98, which is bit lower than unity.

Figure D.7 shows the comparison of Length-scale Calibration results. The slopes for x and y are multiplied. For different three steps, The results are consistent within the uncertainty.

D.4 Non-factorization analysis

Non-factorization analysis considers possible displacement of central beam position for the transverse direction with respect to the beam scan direction. For each step, the representative beam position can be estimated with the log-likelihood function (see Ref. [160]). The function is constructed for all events in each beam step. The input, vertex information, can be obtained from “AnalysisResults_trees.root”. The macro for the fitting can be found in

- PWGMM/Lumi/macros/non_fac.C

The code can be run with the following command, `root -l non_fac.C`. In line 95 to 97, the number of beam steps and ranges for each step should be controlled by an user. Figure D.8 shows the distributions of input parameter, central vertex position and their covariance for each beam step in Pb–Pb collisions. Figure D.9 shows the fitted information for representative vertices for each beam step for x -axis scan. The central x position

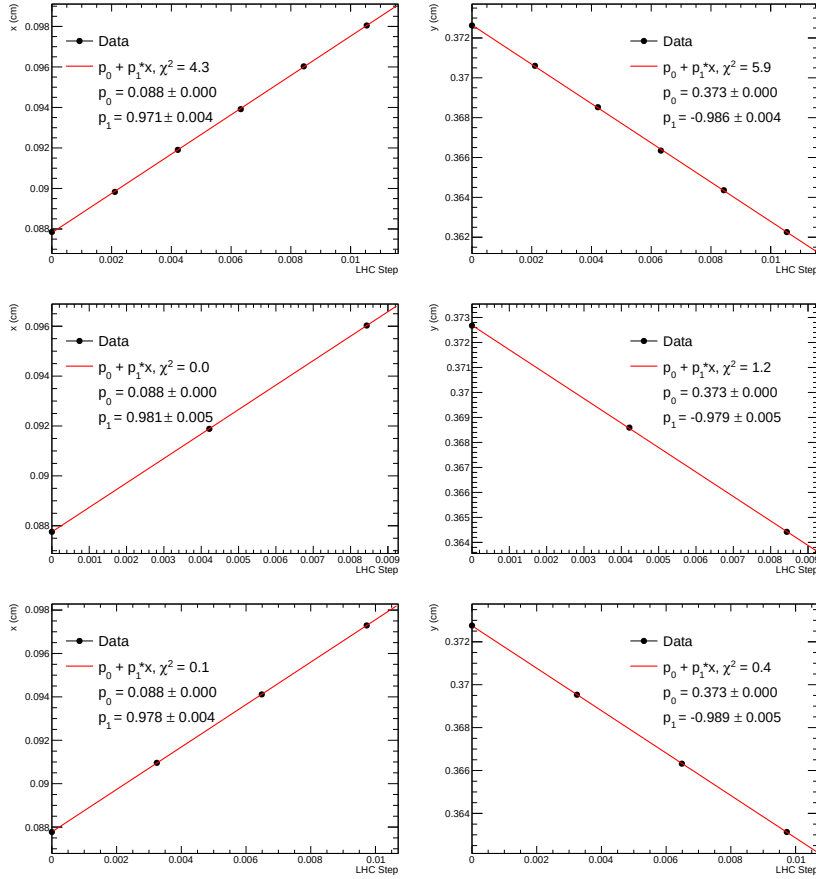


Figure D.6: Comparison between the displacement operated by LHC and the refitted PV for x (left) and y (right) directions

is moving while the central y position is fixed. The same figures for pp collisions are prepared in figure D.11 and figure D.11, showing the inputs and outputs, respectively.

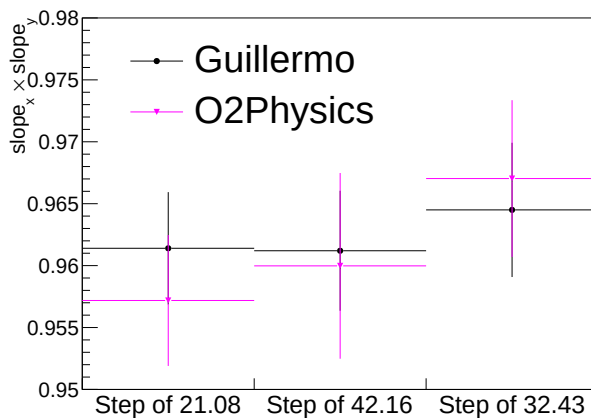


Figure D.7: Comparison of Length-scale Calibration results between Run2 framework made by Guillermo and Run3 framework (O2Physics)

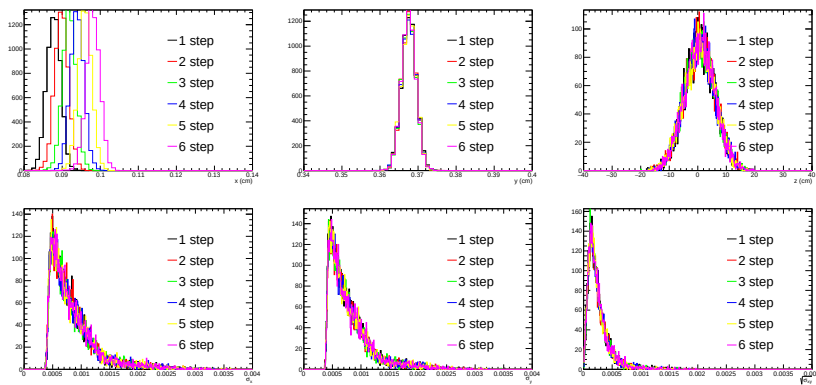


Figure D.8: input parameters for non-factorization analysis in Pb–Pb collisions

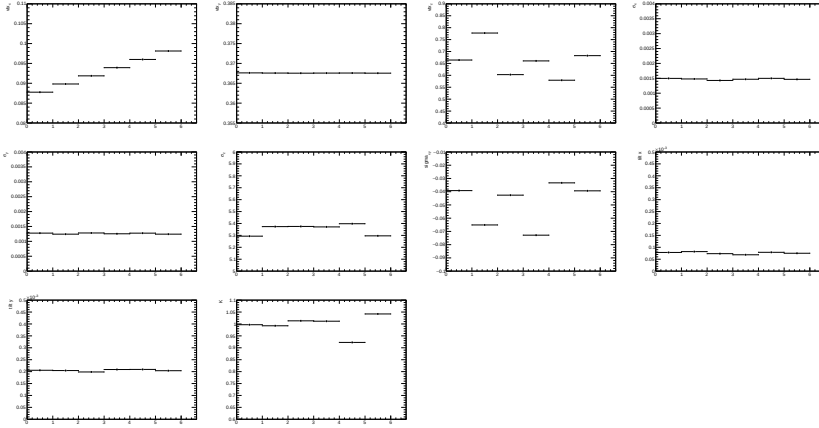


Figure D.9: output parameters for non-factorization analysis in Pb-Pb collisions

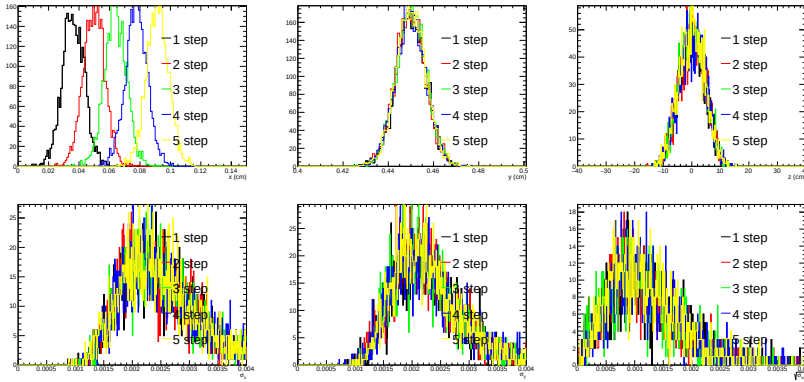


Figure D.10: input parameters for non-factorization analysis in pp collisions

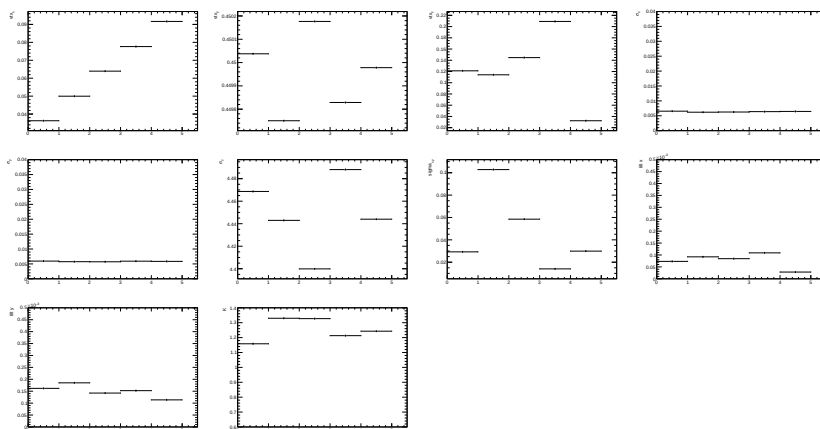


Figure D.11: output parameters for non-factorization analysis in pp collisions

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요약

$f_0(980)$ 공명입자는 예전부터 $\pi\pi$ 산란 실험에서 관측되었으며, 스칼라 경입자 중 하나로 예상되었다. 1970년대 처음으로 관측된 이후 가벼운 스칼라 경입자의 성질은 아직까지 이해되지 않았으며, $f_0(980)$ 의 내부구조에 대한 협의도 이루어지지 않았으며, 네쿼크, 경입자적 분자 혹은 전통적인 쌍쿼크 구조 등 $f_0(980)$ 의 구조에 대한 다양한 제안이 제기되고 있다.

고온 및 고밀도와 같은 극한 환경은 쿼크와 글루온을 자유롭게 하여 쿼크-글루온 플라즈마 (QGP) 를 형성시킨다. QGP는 빅뱅 직후에 생성된다고 여겨진다. 이러한 초기 우주의 상태는 상대론적 중이온 충돌로부터 재현될 수 있으며, 이를 이용하여 QGP의 특성을 연구할 수 있다. QGP의 존재에 대한 많은 과학적 수용과 함께, QGP의 기하학적 구조, 전개 그리고 입자 생성의 광범위한 연구가 진행되었다. 놀랍게도, QGP의 존재를 나타내는 현상들이 양성자-양성자 및 양성자-이온 충돌에서도 관찰되어, QGP의 형성에 대한 추가적인 탐구가 진행되고 있다.

이러한 맥락에서, 상대론적 핵-핵 충돌로부터 생성된 $f_0(980)$ 공명입자를 ALICE 검출기를 이용하여 $f_0(980) \rightarrow \pi^+\pi^-$ 붕괴 채널로 측정하였다. 이 논문은 중간신속도에서 다양한 다중도 등급의 불변산출량 측정 방법론을 기술한다. $f_0(980)$ 공명입자의 입자 산출량 비율은 후기의 강입자 가스 특성을 의논하고 $f_0(980)$ 공명입자의 내부 구조를 탐구하기 위하여 측정하였다. 더욱이, 핵 변형 인자와 타원형 흐름계수의 측정 및 $f_0(980)$ 에 대한 다양한 가정 하에서의 모델 예측은 현재 논문의 물리학적 내용을 강화한다.

주요어 : ALICE, resonance, $f_0(980)$, scalar meson